PROCEEDINGS
OF THE
SYMPOSIUM/WORKSHOP ON SPIN AND SYMMETRIES

TRIUMF, VANCOUVER
JUNE 30-JULY 2, 1989

Editors: W.D. Ramsay and W.T.H. van Oers
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PREFACE

The Symposium/Workshop on Spin and Symmetries was held at TRIUMF, June 30 - July 2, 1989 in conjunction with the 12th International Conference on Few Body Problems in Physics. The Symposium was attended by 76 people including 25 speakers and 3 discussion leaders. The first day of the three day symposium was devoted to Parity Violation, the second to CP and Time Reversal and the final day to Charge Symmetry Breaking. These proceedings have been prepared directly from camera ready copy supplied by the speakers. Each day finished with a discussion session aimed at arriving at a set of conclusions and recommendations. The scientific secretaries for each day, in collaboration with the discussion leader, prepared written versions of the day's conclusions and recommendations. These reports are grouped at the end of this proceedings.

On behalf of the local organizing committee of Charles Davis, Des Ramsay, and Wim van Oers, I would like to thank the many people who contributed to the success of the symposium: the discussion leaders Willy Haeberli (Parity), Ernie Henley (CP & T), and Tony Thomas (CSB), the scientific secretaries Peter Blunden and Jim Birchall (Parity), Gerry Roy and Byron Jennings (CP & T), Norm Davison and Steven Bass (CSB) and, of course, all the speakers. Special thanks to Pat Stewart for attending to the administrative details of registration and meals, Margaret Lear for her help with registration and Alice Hamian and Jeff Schachter for running the audio-visual equipment. Finally, I would like to acknowledge the financial support received from the Natural Sciences and Engineering Research Council of Canada, The International Committee for Symposia on High Energy Spin Physics, the TRIUMF Users Organization, and TRIUMF.

W.D. Ramsay
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INTRODUCTION

This three-day Symposium/Workshop on Spin and Symmetries concerned studies of parity violation, tests of CP and T conservation, and determinations of charge-symmetry breaking. With complex experimental efforts at many intermediate-energy physics laboratories currently under way and with intense theoretical efforts ongoing, it appeared most timely to make an assessment of what has been learned to date, and to chart a course for further experimental and theoretical work. In order to have ample time for discussion, which is most important to the success of such a Symposium/Workshop, the time for the presentations was restricted. Also, at the end of each of the three afternoons, a discussion period was held during which the designated discussion leader formulated a set of conclusions and arrived at a set of recommendations for future work. It is clear that such recommendations are of great importance to the current intermediate-energy physics facilities and also relate to the advent of an advanced hadron facility. With participants at this Symposium/Workshop representing most ongoing efforts in these areas of nuclear science, very fruitful discussions took place leading to the success of this Symposium/Workshop.

Willem T.H. van Oers
A MEASUREMENT OF THE PARITY VIOLATING WEAK ASYMMETRY IN \( \vec{p} - p \) SCATTERING AT 222 MeV

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ABSTRACT

An experiment is underway at TRIUMF\(^1\) to measure the parity violating asymmetry \( A_z \) in the elastic scattering cross section when a beam of longitudinally polarized protons at 222 MeV impinges on a liquid hydrogen target. The beam energy and detector geometries are selected to ensure that only parity mixing in the \( ^3P_2 - ^1D_2 \) partial wave amplitude contributes to the measured asymmetry, which allows the cleanest possible interpretation of \( A_z \) at intermediate energy. Independent measurements will be performed in transmission and scattering geometries to ensure maximum confidence in the experimental results. Two new theoretical predictions of \( A_z \) have been reported recently\(^2,3\), which predict measured asymmetries in our experiment of \( A_z = +0.4 \times 10^{-7} \) and \(+0.7 \times 10^{-7}\), confirming that the proposed accuracy of \( \pm 2 \times 10^{-8} \) in each measurement will provide a significant measurement of parity violation in \( \vec{p} - p \) scattering.

INTRODUCTION

Hadronic weak interactions are of fundamental importance, yet elusive to experiment, since the ratio of weak to strong coupling constants is of order \( 10^{-7} \). Experimental studies are restricted to situations in which the competing strong interaction is blocked by a symmetry principle, such as parity violation or quark flavor changing. At low and intermediate energies, the parity-violating weak nucleon-nucleon (N-N) interaction is described in terms of a meson exchange model involving one weak and one strong vertex. The strong meson-nucleon couplings are taken from a conventional meson exchange description of the strong N-N interaction (e.g. the Bonn potential). The weak couplings are predicted\(^4\) from the Weinberg-Salam model involving W and Z exchanges between quark constituents of the nucleon and meson.

Below the pion production threshold, the weak N-N interaction is described in terms of the exchange of \( \pi \), \( \rho \) and \( \omega \) mesons, neutral scalar mesons being excluded by CP conservation. The couplings are denoted \( (f_\pi^\rho, h_\rho^\rho, h_\rho^\omega, h_\omega^\rho, h_\omega^\omega) \) where the superscripts refer to isospin changes. The isovector pion coupling \( f_\pi \) is unique in that 95% of its predicted value is attributed to neutral current (Z exchange) diagrams in the quark model. Hadronic weak neutral currents can only be observed in quark flavor conserving processes, which restricts experiments to the domain of parity violation in nuclear systems. Measurements\(^5\) of gamma ray circular polarization in \(^{18}\)F have shown \( f_\pi \) to be anomalously small: 4 standard deviations smaller than the 'best' theoretical value\(^4\) from quark model calculations. A recent, novel approach using the chiral effective Lagrangian method\(^6\) gives a prediction for \( f_\pi \) which is a factor of 20 smaller than the quark model value, and in agreement with experiment, while finding less dramatic suppressions of the \( \rho \) and \( \omega \) couplings \( (x_{\frac{1}{5}} \) and \( x_{\frac{2}{3}} \) respectively). While the suppression of \( f_\pi \) in these calculations is encouraging, the \( \rho \) and \( \omega \) suppressions are difficult to reconcile with existing high precision data from \( \vec{p} - p \) scattering at low energy.
In a recent review, Adelberger and Haxton fitted the most significant nuclear parity violation data to a 2 parameter expression based on the quark model formalism of Desplanques et al. Unfortunately, the resulting values of the weak meson-nucleon couplings were only marginally better constrained than the 'reasonable range' estimates from quark model predictions, due to the lack of independent, high-precision data. One of the most significant constraints at present is obtained from parity violation measurements in $\bar{p} - p$ scattering at low energy. In principle, there are two constraints that can be obtained from $\bar{p} - p$ scattering experiments, which determine effective $\rho$ and $\omega$ couplings summed over isospin:

$$h_{\rho}^{pp} = (h_0^\rho + h^\rho + h^\rho / \sqrt{6})$$ and $$h_{\omega}^{pp} = (h_0^\omega + h^\omega);$$
direct pion exchange is suppressed by CP conservation. To obtain a second constraint from $\bar{p} - p$ scattering it is necessary to measure the longitudinal analyzing power, $A_z$, at intermediate energy, where

$$\frac{d\sigma}{d\Omega} = \frac{d\sigma}{d\Omega} (1 + P_2 A_z(\theta))$$ and $$A_z = \frac{1}{\sigma} \int A_z(\theta) \frac{d\sigma}{d\Omega} d\Omega.$$  

The longitudinal analyzing power $A_z$ arises from the interference of opposite parity amplitudes in $\bar{p} - p$ scattering. In general, $A_z(\theta)$ is expressed as a linear combination of partial wave amplitudes: \((S-P),(P-D),(D-F)...\), where the angular distribution of each term is governed by the strong interaction. The relative strengths of each term are determined by matrix elements of the weak meson exchange interaction and must be calculated using appropriate wave functions. The \(1S_0-3P_0\) amplitude is well determined from precision measurements of $A_z$ at 15 and 45 MeV ($A_z = (-1.5\pm0.2) \times 10^{-7}$). In contrast, the $3P_2-1D_2$ amplitude, which contributes significantly to $A_z$ above 100 MeV, has not been measured. In the conventional single meson exchange model of the weak N-N interaction, these two amplitudes have complementary dependences on $h_{\rho}^{pp}$ and $h_{\omega}^{pp}$. Simonius has shown that the $3P_2-1D_2$ amplitude depends only on weak $\rho$ exchange in this model, whereas $\rho$ and $\omega$ contribute to the $1S_0-3P_0$ amplitude with approximately equal weight. The energy dependence of the two lowest partial wave contributions to $A_z$ as predicted in the recent meson exchange calculation of Simonius, is shown in Figure 1. In a separate contribution to this workshop, Silbar proposes a new mechanism whereby weak charged pion exchange $(2\pi)$ can contribute to $A_z$ in $\bar{p} - p$ scattering, via a $(n-A^{++})$ intermediate state. The mechanism is predominantly $J=2$, and therefore competes directly with $\rho$ exchange in the $3P_2-1D_2$ amplitude.

The unique feature of the present experiment is that the $1S_0-3P_0$ amplitude can be eliminated from $A_z$ at 230 MeV, since its angular distribution integrates to zero at this energy. This results directly from the cancellation of the strong $1S_0$ and $3P_0$ phase shifts, and is completely independent of the values of the weak meson-nucleon coupling strengths. With the exception of a small contribution from the $1D_2-3F_2$ amplitude ($\sim 5\%$), the longitudinal analyzing power in $\bar{p} - p$ scattering measures the $3P_2-1D_2$ term alone. Two new, independent calculations have been performed using the weak meson exchange model, finding $A_z(230 \text{ MeV}) = +0.4 \times 10^{-7}$ and $+0.7 \times 10^{-7}$ for $\rho$ exchange only. Thus, our proposed measurement of $A_z$ to $\pm 2 \times 10^{-8}$ in two independent geometries will obtain an accuracy of $\pm 25\%$ or better.

THE TRIUMF EXPERIMENT

The aim of the TRIUMF experiment is to measure the $3P_2-1D_2$ partial wave contribution to the longitudinal analyzing power $A_z$ in $\bar{p} - p$ scattering to an accuracy of $\pm 2 \times 10^{-8}$ or better. This will be achieved in two separate measurements by transmission and scattering methods. It is planned to move rapidly to complete the first measurements in transmission
geometry, and to add the scattering detector at a later stage. In the first measurements, the incident and transmitted protons from a LH$_2$ target will be detected by parallel plate ionization chambers operated in current mode; at a later stage, a large, cylindrically symmetric parallel plate ionization chamber will be added downstream of the target to collect the scattered particles.

![Figure 1: Partial wave contributions to $A_z$ in $\vec{p} - p$ scattering.](image)

In addition to the main ionization detectors and the LH$_2$ target, the apparatus includes high precision beam monitoring devices which measure the beam intensity distribution and transverse polarization distribution in $x$ and $y$ at two points, plus a split plate intensity monitor which provides the error signal to a fast analog feedback system which stabilizes the beam position. These devices are necessary to provide a means of monitoring false parity violating asymmetries which will be produced if the beam is not identical in the $+$ and $-$ helicity states, apart from exact reversal of the polarization direction. The approach used to design the experiment is to account for systematic effects in the following manner:

$$\Delta A_z = \frac{\partial A_z}{\partial x_i} \Delta x_i$$

where $x_i$ is any beam property (e.g. current, position, transverse polarization...) and $\Delta x_i$ is the helicity correlated change in that property. The main detectors are designed to minimize the sensitivities $\frac{\partial A_z}{\partial x_i}$, based on Monte Carlo calculations accounting for realistic beam properties. The beam monitoring devices are designed to achieve measurements of $\Delta x_i$ to sufficient accuracy that the resulting false asymmetries are independently...
determined to an accuracy which exceeds that of the parity violating asymmetry measured in the main detectors on a comparable timescale. The main detectors will be calibrated to determine the sensitivities $\frac{\partial A_s}{\partial x}$ by introducing known modulations of the beam properties $x_i$ in ancillary experiments. The planned layout of the parity violation apparatus on beamline 4A at TRIUMF is shown below in Figure 2.

Figure 2: Layout of parity violation apparatus showing existing shielding configuration in Beamline 4A.

Beam Energy

Monte Carlo calculations have been performed to determine the optimum beam energy and geometry of the scattering and transmission detectors to ensure cancellation of the $^1S_0 - ^3P_0$ partial wave contribution to the measured asymmetries in the two experiments. The input to the Monte Carlo program included a beam profile whose size and emittance were taken from realistic beam transport calculations. The simulated beam scattered from a 20 cm LH$_2$ target in which multiple scattering and energy loss were accounted for. The shapes of the $^1S_0 - ^3P_0$ and $^3P_2 - ^1D_2$ angular distributions were calculated using helicity amplitudes and phase shifts taken from SAID$^{11}$ as discussed by Simonius$^2$; multiple scattering and energy loss in the entrance windows and H$_2$ gas of the ionization chambers were also accounted for in order to determine detector response functions for the scattering and transmission detectors.
The requirements for the transmission experiment were found to be least flexible: for a symmetric detector that accepts scattered protons from $0^\circ$ to $6^\circ$ (lab.), the incident beam energy which cancels the $^1S_0 - ^3P_0$ contribution ranges from 223.2 to 222.0 MeV, as illustrated in Figure 3. In this range of beam energies, a contour of the small and large angle boundaries ($\theta_1, \theta_2$) of the scattering detector which cancel the $^1S_0 - ^3P_0$ contribution ranges from $(2.8^\circ, 40.4^\circ)$ to $(11.0^\circ, 38.8^\circ)$. When the calculations were repeated using different phase shift solutions (Arndt SP88, Saclay S260, Bonn potential, Paris potential), the $^1S_0 - ^3P_0$ contribution was found to be less than 4% of the $^3P_2 - ^1D_2$ in all cases considered. Furthermore, it was found that the ratio of asymmetries in the two geometries was $(A_x^T/A_T^F) = 0.97$ (for $\theta_1 = 6^\circ$, 222 MeV); this constraint will be used to check the consistency of the two measurements. It should be emphasized that both the cancellation of the $^1S_0 - ^3P_0$ contribution and the ratio of measured asymmetries are independent of the parity violating interaction and its coupling strengths.

Figure 3: Monte Carlo calculation of the $^1S_0 - ^3P_0$ contribution $\times 10^7$ (labelled $K_0+$) to $A_z$ measured in transmission geometry, versus incident beam energy, for various angular acceptances of the transmission detector.

**Ion Source**

The parity violation measurements will require approximately 300 hr of running with a 500 nA beam of longitudinally polarized protons, $P_z \geq 0.7$. The beam will be provided by the
optically pumped polarized ion source (OPPIS) at TRIUMF, which is in principle ideal for these measurements, since polarization reversal is achieved by changing the frequency of laser light with the minimum possible changes to beam emittance and intensity. The collaboration has been working closely with the TRIUMF Ion Source Group as the new polarized source is being commissioned, in order to optimize its performance for the parity measurements. A mechanism for creating a small \((\sim 10\%)\) source of pure intensity modulation in the beam from OPPIS will be developed in order to calibrate the linearity of the in-beam ionization chambers used in the parity violation measurements.

It has been possible to carry out a few preliminary tests of the stability of the beam intensity from the new source, and these results have been very encouraging. A search for changes in beam intensity correlated with spin state was carried out using a (prototype) low noise transverse field ionization chamber in beamline 4A. The variation of beam intensity upon spin reversal, which sets a limit on the differential nonlinearity of the ion chambers and associated electronics, was found to be \(dI/I \leq 0.001\). This is to be compared with the helicity correlated \(dI/I\) from the Lamb Shift polarized Ion source used in the Los Alamos 800 MeV experiment, which was \(\sim 10^{-4}\). The TRIUMF OPPIS is expected to exhibit even smaller changes of beam intensity upon spin flip, once the source parameters have been optimized.

**Transverse Field Ionization Chambers**

To perform the transmission mode measurement, the beam current will be measured upstream and downstream of a liquid hydrogen target with a pair of transverse field parallel plate ionization chambers (TRIC's) filled with hydrogen gas. Because the chambers must be operated in current mode to meet the high statistical accuracy required for the experiment, the gas gain \((20 P_{\text{stg}} \text{ cm}^{-1}\) at 222 MeV) must be constant as a function of time and beam current. This puts stringent requirements on the gas purity, since even trace amounts of electronegative impurities (at the ppm level) can compromise the linearity of the ion chamber response at levels that are intolerable for the experiment. In the final configuration, it will be necessary to bake the chambers under high vacuum for several days before filling them with ultra-high purity hydrogen gas for the parity violation measurements.

Test measurements have been performed with prototype TRIC’s on loan from Los Alamos National Laboratory, that were used in the 800 MeV \(A_2\) measurements\(^{12}\). Optimal gas purity conditions have been difficult to attain due to the necessity of moving the apparatus in and out of beamline 4A for short test runs, with limited access to the beamline in advance of the test periods. However, it has been possible to set upper limits for the intrinsic noise and nonlinearity of the TRIC’s under these conditions, and to establish important design criteria for the final ionization chambers. Numerical simulations have been undertaken to predict the behavior of the downstream TRIC when exposed to scattered beam from the LH\(_2\) target. A detailed model has been developed, accounting for \(\delta\) ray production and spallation processes occurring at the entrance windows, electrode plates and walls of the chamber. Parameters were varied to determine the optimum geometry of the planned detectors. In all cases, the predicted noise is below the upper limit imposed by the condition of \(\leq 300\) hours counting time to achieve the planned statistical error in \(A_2\) provided that the beam is free of halo at the \(10^{-4}\text{cm}^{-2}\) level. When the calculations are adapted to a beam energy of 800 MeV, agreement is obtained with the measured noise figures from the Los Alamos experiment. The design
of the final TRIC's, as shown in Figure 4, incorporates improved field shaping electrodes to ensure that the neutral axes for position and transverse polarization fluctuations coincide.

Figure 4: Layout of the new transverse field ionization chambers.

Scattering Detector

The geometry of the planned scattering detector is shown in Figure 5. It is a planar, cylindrically symmetric ionization chamber filled with hydrogen gas at 1 Atm. There is a central high voltage plane, in front of and behind which is a symmetric arrangement of sense planes at ground potential. The primary proton beam passes through a central hole in the chamber to avoid both heating the gas and contributions to the signal from primary protons scattering from entrance and exit windows. Based on tests in beamline 1B of a small, axial field ionization chamber, a noise figure of $\alpha \approx 1$ is expected for this detector, where $\alpha$ is averaged over scattering angle, weighted with the detector response function. The counting time in hours for the scattering experiment, assuming 500 nA primary proton beam with $P_z=0.7$ is given by: $t \approx 22 \left(1 + \alpha_s^2 + \alpha_t^2\right)$, where $\alpha_s$ is the noise factor for the scattering detector, $\alpha_t$ is the noise factor for the upstream transmission detector used for normalization, $S = 0.02 \frac{L}{20}$ is the scattering probability in the liquid hydrogen target with $L$ the target length in cm, and we have made the approximation that $S \ll 1$. Sensitivity to transverse polarization is the major concern for this device, and will be determined in a series of calibration measurements with transversely polarized beam. By segmenting one of the sense planes into 4 quadrants, the scattering detector itself can be used as an extremely
sensitive polarimeter, referred to the liquid hydrogen target. This feature will be extremely useful in tuning the beamline for minimum transverse polarization components.

Figure 5: Geometry of the planned scattering detector (side view).

Transverse Polarization Components

The major systematic error which must be overcome is an apparent parity-violating effect generated by residual transverse polarization components in the beam that couple to a relatively large parity-allowed transverse analyzing power $A_y \sim 0.3$. The false asymmetry is proportional to the first moment of transverse polarization of the beam, which has two sources: (i) a net transverse polarization coupled with an offset of the beam from the symmetry axis of the detection apparatus; (ii) an intrinsic first moment of transverse polarization, which can arise from a finite energy spread in the beam and/or nonuniform magnetic fields in the beamline elements. To minimize the false asymmetry, beam properties must be identical in the two helicity states and the detection apparatus must be highly symmetric.

The vertical transverse polarization of the beam from the cyclotron will be transformed to longitudinal polarization by a combination of dipole and superconducting solenoid magnets, which are presently installed on beamline 4 at TRIUMF. The net transverse polarization will be minimized by careful design of beam optics and control of the magnets used to rotate the polarization from transverse to longitudinal. Possibilities of a slow feedback loop to stabilize these beamline elements are being explored, with the aim of suppressing transverse
polarization components below $10^{-3}$.

Initial measurements of transverse polarization profiles have been performed for a longitudinally polarized beam at the T1 target location in beamline 4B. Thin strip targets were mounted in the target ladder of a conventional 4-branch polarimeter, and were scanned slowly through the beam by a stepping motor under computer control. Horizontal and vertical polarization components were measured for a beam from the Lamb Shift polarized ion source, since the optically pumped ion source was not sufficiently stable at the time of the test run. Based on this experience, and on results from improved systematic error calculations, a design has been developed for a 4-branch ‘conventional’ counting polarimeter with rotating blade targets, similar to polarimeters used in the S.I.N. parity violation experiments.

Systematic Error Calculations

Initial calculations of the sensitivity to transverse polarization components were performed for idealized scattering and transmission detectors having perfect cylindrical symmetry. When edge effects were carefully considered in the scattering case, it was found that a crossover occurs in the sign of the false asymmetry as a function of the small angle boundary of the detector acceptance between 5° and 6° (lab), as seen in Figure 6. This means that it should be possible to choose a nominal small angle cutoff to design the scattering detector and ‘tune’ the position of the detector along the beamline using transversely polarized beam to find a position where the false asymmetry is minimal, corresponding to the angle cutoff for zero false asymmetry predicted by the calculations. Work is now focusing on the transverse field ionization chambers to determine whether a similar feature is found when the exact geometry of the detectors is carefully accounted for.

Fast Feedback System

Error estimates based on a net transverse polarization of 0.01 have shown that the beam centroid must be stabilized at two points on the symmetry axis of the apparatus to within $10 \mu$m. Initial measurements of centroid motion using a split plate ionization chamber (SPIC) indicated the need for a beam position feedback system. A prototype single loop feedback system with x and y position signals from a hydrogen-filled SPIC successfully reduced beam excursions to a few $\mu$m at frequencies up to 1 kHz. The correction system is based on low inductance air core steering magnets driven by high power operational amplifiers. The feedback system is automatically disabled when the beam current drops below a preset value, and a logic signal indicating the status of the loops is provided by the analog dividers. The electronics components, including custom built high power steering magnet drivers, analog dividers and low noise preamplifiers for the detector signals, are now essentially in their final form.

Although the prototype feedback system has been extremely successful at controlling beam position excursions, the relatively thick windows and sense planes of the split plate ionization chambers cause a large amount of beam broadening which cannot be tolerated for the parity violation measurements. In the past year, one of the two hydrogen-filled SPIC’s borrowed from Los Alamos for prototyping and development work has been replaced by a split foil SEM device constructed at TRIUMF which works in vacuum, reducing the amount
Figure 6: Monte Carlo simulation of the false asymmetry $\delta A_z$ in scattering geometry, induced by a first moment of transverse polarization ($xP_y$) correlated with spin flip, for various inner cutoff angles of the scattering detector. Outer cutoff: 45°(lab).

Intensity Profile Monitor

A number of tests have been performed to determine the quality of signal obtainable with various wire and foil strip configurations to measure the intensity profile of the beam. Design criteria for the final monitor are: low mass, high linearity, high signal to noise ratio, and high accuracy/reproducibility of wire/strip positions. The first two criteria strongly favour a multi-strip SEM foil device operating in vacuum. Systematic error calculations set the requirements of measuring the position centroid of the beam to ± 3 µm in 1 hour, translating to a maximum tolerable detector noise of 50 pA per channel for an SEM foil strip device; the foil strip positions must be known to an accuracy of ±10 µm, and the amplifier gains must be matched and stable to ± 0.3%. After considerable development work, success of material in the beam by roughly an order of magnitude. The signals from the SEM device are much smaller than from the SPIC, and necessitated the construction of new, sensitive electrometer-based preamplifiers with two orders of magnitude higher gain than the original low noise preamplifiers used with the SPIC's. The feedback system has run reliably in two loop mode with one SPIC and one SEM monitor and exhibited adequate beam position control. A second SEM monitor will be built this winter to complete the system.
was recently achieved at manufacturing thin multi-strip foil arrays that approach the geometrical tolerances set by the systematic error calculations. A 32-channel electrometer preamp unit has been constructed and successfully tested together with several foil strip arrays.

During a recent test run, a 30 nA beam was systematically deflected \( \pm 5 \mu m \) at 0.7 Hz about the central axis of the profile monitor, and beam intensity profiles were accumulated in the two position states. The measured centroid shifts of the beam profile in three consecutive 30 min. runs were 10\( \mu m \), 7\( \mu m \) and 12\( \mu m \), demonstrating that the sensitivity criterion established by systematic error calculations has already been met for this device. The statistical errors of each measurement as deduced from the measured noise on the individual foil strips with beam off are less than 1\( \mu m \), but the variation in the three profile measurements is roughly twice this, possibly due to slow drift in the beam position at the \( \mu m \) level coupled with imperfectly matched gains in the individual preamps and ADC channels. Work is continuing on an algorithm for gain matching to the required accuracy, guided by numerical simulations, and on improving the foil strip manufacturing technique. A fast electronic system for readout of the 128 foil strip signals (from 2 monitors) for each spin state is also being developed.

**FUTURE PLANS**

The TRIUMF parity violation experiment is presently in the second year of an instrumentation development programme. During the next year, the final apparatus required to perform the transmission measurements will be constructed and tested. Work is underway to modify beamline 4A at TRIUMF in order to accommodate the beam transport requirements of the parity violation measurements, and a design for the liquid hydrogen target is underway. Data taking in transmission mode is currently scheduled for 1991-2.

**REFERENCES**

11. R.A. Arndt, Interactive Dial-in Program SAID
PARITY VIOLATION IN THE NUCLEON-NUCLEON INTERACTION

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ABSTRACT

I discuss the present status of our understanding of parity nonconservation (PNC) in the nucleon-nucleon interaction, and some of the difficulties inherent in nuclear tests of PNC. I also discuss the nucleon/nuclear anapole moment, the parity violating coupling of the photon, and its relation to the PNC NN interaction.

INTRODUCTION

In the oral version of this talk I reviewed NN and nuclear tests of PNC, discussed new efforts to improve calculations of nuclear PNC matrix elements, and described a recent calculation of nucleon and nuclear anapole moments. The review portion of my talk covered familiar themes: the importance of the PNC NN interaction as a test of neutral weak currents, the single-meson-exchange description of the PNC NN interaction, possible observables, and PNC mixing in $^{18}$F (and the associated $^{18}$Ne $\beta$ decay) as an example of the virtues and limitations of nuclear PNC tests. These topics are discussed in the literature (see, for example, Ref. 1) and will not be repeated here. One concludes from such a review that there is clear evidence for an isoscalar PNC NN interaction with a strength roughly in accord with the "best value" estimate of Desplanques, Donoghue, and Holstein (DDH). However there is no evidence for the expected neutral current enhancement of the isovector pion-exchange PNC interaction.

A group of us at Seattle (C. Johnson, V Zeps, W. Haxton, and E. Adelberger) have been trying to better understand the limitations of shell model calculations of PNC nuclear matrix elements. For instance in $^{19}$F it can be shown that calculated matrix elements depend sensitively on nuclear deformation, the radial dependence of single-particle wave functions, and suppression effects associated with the pairing force. Thus if the nuclear model is inaccurate in any of these respects, the resulting calculation may be unreliable. These considerations lead to the conclusion that our best hope for reliable theoretical results may be large-basis shell model calculations for light nuclei employing realistic single-particle bases. Our current efforts focus on a full 4$\hbar\omega$ calculation of $^{16}$O, and the use of the resulting "tuned" Hamiltonian to predict PNC matrix elements in $^{14}$N, $^{18}$F, $^{19}$F, and $^{21}$Ne. We interpret the resulting shell model density matrices in terms of a Woods-Saxon basis.

Although this program is incomplete, some interesting results are already in hand. In $^{16}$O there exists a well-studied $0^-1 \leftrightarrow 0^+0$ axial-charge transition similar to those studied in $^{18}$F and $^{19}$F. The nuclear structure considerations that determine the strength of axial-charge $\beta$-decay transitions are identical to those governing the mixing of PNC doublets: this argument is exploited in the $^{18}$Ne $\beta$-decay calibration of the $^{18}$F PNC mixing matrix element. Thus the $^{16}$O transition can be used to test nuclear structure issues important to PNC calculations.

The effect of the pairing force on operator matrix elements depends on the behavior of the operator under particle-hole conjugation. Axial charge multipoles

Axial charge multipoles
and the more familiar transverse electric multipoles are both odd under particle-hole conjugation. The pairing force suppresses matrix elements of odd operators. In terms of density matrix elements, the pairing suppression is embodied in the coefficients \( \psi_{\beta \alpha} \) that, having the same sign as \( \psi_{\alpha \beta} \), generate cancellations in \( \psi_{\alpha \beta} = \psi_{\alpha \beta} - \psi_{\beta \alpha} \). Odd operators depend on this linear combination \( \psi_{\alpha \beta} \). In the shell model the \( \psi_{\beta \alpha} \) enter only in very large-basis calculations: for \(^{16}\text{O}\) (or \(^{18}\text{F}\)) \( \psi_{\alpha 1/2,2s_{1/2}} \) interferes with the dominant \( \psi_{2s_{1/2}1p_{1/2}} \) contribution only if \( 2\hbar\omega \) pairing excitations are included in the wave functions.

Full \( 2\hbar\omega \) shell model calculations for \(^{18}\text{F}/^{18}\text{Ne} \) produced a sizeable suppression in the \(^{18}\text{Ne} \) \( \beta \)-decay matrix element, in accord with experiment. No such calculations have been performed for \(^{19}\text{F} \) and \(^{21}\text{Ne} \). The naive explanation for the importance of \( 2\hbar\omega \) configurations is that they couple strongly to the dominant \( 1\hbar\omega \) configurations in the \( 0^- \) state through the axial charge operator.

However, if \( 2\hbar\omega \) components in the \( 0^+ \) wave function are important, one might conclude that \( 3\hbar\omega \) corrections to the \( 0^- \) wave functions could also play a role. That is, does the inclusion of \( n\hbar\omega \) configurations quickly converge to give a stable result for axial charge matrix elements, or is the convergence so slow that reliable calculations of PNC matrix elements become prohibitively difficult? The \(^{16}\text{O} \) calculations Johnson and I have performed provide a first test of this convergence.\(^5\)

The results shown in Table 1 are encouraging. As in \(^{18}\text{F} \), the inclusion of \( 2\hbar\omega \) configurations sharply reduces the axial charge matrix elements mediating \( \beta \)-decay \((q \approx 0 \text{ MeV}) \) and \( \mu \)-capture \((q \approx 95 \text{ MeV}) \). After \( 3\hbar\omega \) and \( 4\hbar\omega \) corrections are added to the \( 0^- \) and \( 0^+ \) states, respectively, the results change only by a few percent.

<table>
<thead>
<tr>
<th>(&lt; J_0^A ) &gt;</th>
<th>( 0 + 1\hbar\omega )</th>
<th>( 1 + 2\hbar\omega )</th>
<th>( 3 + 4\hbar\omega )</th>
</tr>
</thead>
<tbody>
<tr>
<td>( \beta )-decay</td>
<td>.042</td>
<td>.027</td>
<td>.029</td>
</tr>
<tr>
<td>( \mu )-capture</td>
<td>.041</td>
<td>.028</td>
<td>.029</td>
</tr>
</tbody>
</table>

A topic related to PNC in nuclei is the nuclear anapole moment, a coupling E. Henley, M. Musolf, and I have recently investigated.\(^6\) For spin-1/2 fermions two PNC static couplings to the electromagnetic field can arise, the electric dipole moment, which violates both parity- (P) and time-reversal (T) invariance, and the anapole moment, which violates P but conserves T. Although first discussed by Zel'dovich 30 years ago\(^7\), the anapole moment has not been studied to the same extent as the EDM. Naively, one expects the experimental effects of the anapole moment (generated by weak radiative corrections like those of Fig. 1a) to be swamped by tree-level neutral current processes, the former being suppressed by roughly a factor of \( \alpha \) over the latter. However, nuclear or atomic many-body effects may enhance the size of the anapole moment to the point where it competes effectively with \( Z_0 \) exchange.\(^8,9\) Thus the anapole moment could be important in experimental tests of the standard model that use charged particles (e.g. electrons) as probes.
Figure 1. Examples of one-body (a) and exchange-current (b) contributions to the nuclear anapole moment.

For an on-shell spin-1/2 fermion (and therefore a virtual photon), current conservation and Lorentz invariance requires PNC corrections to matrix elements of the electromagnetic current to have the form

$$< p' | j^{em}_\mu(0) | p >_{\text{PNC}} = \frac{a(q^2)}{m_N^2} \bar{u}(p')(\gamma_\mu q^\mu - q^2 \gamma_5)\gamma_5 u(p)$$

where $q_\mu = p' - p$ is the momentum transfer to the fermion. The form factor $a(q^2)$ evaluated at $q^2 = 0$ defines the fermion anapole moment. For interactions with on-shell external particles, only the $\gamma_\mu \gamma_5$ term will contribute to PNC amplitudes. The explicit $q^2$ factor in this term cancels the $q^2$ from the photon propagator, so that a contact interaction results. Thus, the anapole moment produces the same type of coordinate-space contact interaction as low-$q^2$ neutral current processes.

The electron\textsuperscript{10} (and constituent quark\textsuperscript{11}) anapole moment calculated in the standard model is gauge-dependent; in physical processes like $e-e$ scattering radiative corrections (e.g., two-boson exchange) must be combined with $a(q^2)$ to produce a gauge-independent result.\textsuperscript{11} This complication does not arise here, where we restrict our attention to the effects generated by an on-shell PNC $\pi NN$ vertex. The pion one-loop corrections (Fig. 1a) to the $\gamma NN$ vertex then provide a gauge-independent estimate of the meson cloud contribution to the nucleon anapole moment. The results for pseudovector and pseudoscalar strong couplings agree up to an ambiguity...
associated with the linear divergence of the pseudovector loop integral,

\[ a(0)_{\pi-cloud} = e \left( \frac{f_\pi g_\pi N N}{8\sqrt{2\pi^2}} \right) (\alpha_s + \alpha_v \tau_3) = a_s(0) + a_v(0)\tau_3, \tag{2} \]

where \( g_{\pi N N} \) is the usual strong coupling and \( f_\pi \) is the weak PNC coupling (but defined as minus that of Ref. 2). The terms \( \alpha_{s,v} \) contain logarithms \( \log(m_\pi/m_N) \); their numerical values are \( \alpha_s \simeq 1.6 \) and \( \alpha_v \simeq 0.4 \). Similar logarithms suppress the contributions from heavier mesons. While, as noted above, other terms contribute to nucleon anapole amplitudes, we will use Eq. (2) to estimate the scale of \( a(0) \). The PNC electron-proton potential generated by \( a(0) \) is smaller than the isovector tree level \( V(\text{electron}) - A(\text{proton}) \) neutral current interaction by a factor of 3.8 \( \alpha(f_\pi/f_\pi^{DDH}) \), where \( f_\pi^{DDH} = 4.5 \times 10^{-7} \) is the best value coupling of Ref. 2. Despite the distinctive isoscalar contribution, the nucleon anapole moment may be impossible to isolate experimentally.

A more tractable task may be the observation of anapole moments in nuclei, where many-body effects enhance the anapole coupling. Two distinct effects are of interest: the "polarization" contribution due to the PNC mixing of the nuclear ground state with nearby excited states with the same angular momentum but opposite parity, and two-body (and, in principle, higher order) currents arising from the interaction of the photon with \( NN \) pairs and virtual mesons. Long-range pion exchange contributions should dominate these nuclear amplitudes.

Time reversal symmetry restricts the electromagnetic static moments of a nucleus to even multipoles of the electromagnetic charge operator (\( C_0, C_2 \ldots \)) and odd multipoles of the magnetic and electric current operators (\( M_1, M_3, \ldots, E_1, E_3 \ldots \)). [Note that the \( T \)-odd electric dipole moment is a \( C_1 \) coupling.] If we consider only those PNC photon couplings arising from the weak interaction in first order, parity further restricts the nonzero matrix elements to the odd electric projections of the axial (anapole) current \( <gs|\hat{E}_A|^gs> \) and to the odd electric projections of the ordinary vector current that can arise because of wave function polarization

\[ \sum_n <gs|\hat{E}_A|^gs> \frac{n^-}{E_g - E_n} H_{PNC} |n^-> + <gs|\hat{H}_{PNC} \frac{n^-}{E_g - E_n} \hat{E}_A^V |gs^+>. \tag{3} \]

Here \( |gs^+> \) is the ground state, \( |n^-> \) denotes an excited state having the same spin but opposite parity, and \( H_{PNC} \) is the parity nonconserving NN interaction. For dipole transitions, the form of the static \( E_1^A \) and \( E_1^V \) operators is determined by an extended Siegert's theorem.\(^{12}\)

\[ <gs||E_1||gs> = \frac{-i\alpha^2}{9\sqrt{6\pi}} \int d\tau^2 \frac{d\tau^2}{2\pi} <gs||j_{em}(\tau) + \sqrt{2}Y_2(\Omega_\tau)\otimes j_{em}(\tau)|gs>. \tag{4} \]

where \( \otimes \) denotes a spherical tensor product and \( || \) denotes a reduced matrix element. Thus current conservation requires the ground state \( E_1 \) moment to have the same leading \( q^2 \) behavior as the transverse part of the single nucleon current of Eq. (1). Like its free nucleon counterpart, the nuclear anapole moment generates a contact interaction between the nucleus and an on-shell external particle. The explicit form
of Eq. (4) depends on properly removing those $0(q^2)$ terms in the E1 matrix element that vanish because of current conservation.

The one-body contribution to Eq. (4) is derived by reducing the axial single-nucleon current operator (Eq. (1)) nonrelativistically and transforming into coordinate space,

$$<gs|E1|gs>_{1\text{-body}} = \frac{i}{q^2 \to 0} \frac{\vec{q}^2}{\sqrt{6\pi} m^2_N} <gs| \sum_{i=1}^{A} (a_s(0) + a_v(0)\tau_3(i))\vec{J}(i)||gs>.$$  (5)

The anapole operator within the matrix element, which we will denote $\vec{A}_{\text{nuc}}(1)$, is the direct analog of $a(0)\vec{r}\gamma_5$ appearing in Eq. (1). In a naive nuclear model with a single nucleon outside a spin-paired core, the one-body contribution to the nuclear anapole moment is just the anapole moment of the valence nucleon.

The nuclear current operator appearing in Eq. (4) also contains important two-body currents. The operators generated by the $NN$ and pionic diagrams (Fig. 1b) are given in Ref. 6. One can determine the importance of the two-body currents by reducing them to an approximate one-body form

$$<a|\vec{A}_{\text{nuc}}^{\text{eff}}(1)|\beta> = \sum_{\delta < F} <a|\vec{A}_{\text{nuc}}^{\text{eff}}(2)|\beta \delta - \delta \beta>,$$

where the sum is taken over the nuclear core. The Fermi gas model, with a spin symmetric but isospin asymmetric core, yields for the $NN$ contribution

$$\vec{A}_{\text{NN}}^{\text{eff}}(1) = 2.74 a_s(0) m_N^2 m^2 \sum_{i=1}^{A} \rho(\vec{r}_i) \tau(i)^2 \left[ \vec{\sigma}(i) + \sqrt{2\pi} (Y_2(\vec{r}_i) \otimes \vec{\sigma}(i)) \right]$$

$$\left[ \frac{Z}{A} \omega_\tau^Z \left( 1 - \frac{2}{3} \tau_3(i) \right) + \frac{N}{A} \omega_\nu^N \left( 1 + \frac{2}{3} \tau_3(i) \right) \right].$$  (6)

with $\rho(\vec{r}_i)$ the nuclear density operator and $\omega_\tau^Z(\omega_\nu^N)$ a proton (neutron) Fermi gas response function that depends on $k(i)/k_F$, the nucleon momentum as a fraction of the Fermi momentum. The $\omega$'s vary only gently, ranging from 0.33 to 0.19 as $k(i)/k_F$ increases from 0 to 1. Thus we can approximate $\omega \sim 0.25$.

Using a nuclear density of 0.195/fm$^3$, we conclude that nuclei will exhibit enhanced isoscalar anapole moments $\sim 0.9 A^{2/3} a_s(0)$ due to the $NN$ exchange current. The inclusion of short range correlations reduces this estimate by 25%. Thus we expect the net isoscalar $NN$ anapole moment for $^{133}$Cs (an example discussed below) to be 17 times the single nucleon value. The isovector anapole moment is smaller by a factor of $2(Z-N)/3A$ due to a cancellation between contributions from core neutrons and protons.

Finally, we turn to the third piece of the anapole moment, the polarization contribution of Eq. (3). This contribution is similar to the exchange currents in that the net effect of interactions with core nucleons is a PNC polarization of orbits of valence nucleons. It differs in that the energy denominators governing the mixing are determined by the spectrum of excited nuclear states. Thus, if one selects a nucleus where an excited state of the same spin but opposite parity appears very
near the ground state, a substantial enhancement can result. In nuclei where no such enhancement occurs, it is argued in Ref. 8 that the polarization contribution also scales like $A^{2/3}$. To remain consistent with our treatment of other terms, we retain only the pionic contribution to $H_{PNC}$ in evaluating polarization sums.

Calculations were performed for two nuclei, $^{19}$F and $^{133}$Cs. The former is an example of a nucleus where a ground-state parity doublet (the $1/2^+ - 1/2^-$ splitting is 110 keV) could lead to an enhanced polarization contribution to the anapole moment. The latter was the subject of a recent atomic experiment on anapole moments. Details of these shell model calculations, including an elegant Lanczos procedure for evaluating polarization sums, are given in Ref. 6. The results are summarized in Table 2. The total anapole moment is dominated by the polarization and exchange current terms, with these contributing constructively in the ratio of about 4 to 1 in $^{133}$Cs. The contribution of the $^{19}$F 110 keV doublet (with the mixing matrix element taken from the measured $^{19}$Ne $1/2^+ ightarrow 1/2^- \beta$-decay rate, not the shell model) to the polarization sum is significant but not extraordinary, accounting for 53% of the total.

Table 2. Shell model estimates of the one-body, polarization, and exchange current contributions to the anapole matrix element $<gs|A_1|gs>$ in units of $ef_\pi$. The last column gives the ratio of the anapole interaction with an on-shell electron to that generated by $Z_0$ exchange, assuming $f_\pi = f_\pi^{DDH}$ and $\sin^2 \theta_w = 0.23$.

<table>
<thead>
<tr>
<th>nucleus</th>
<th>one-body</th>
<th>polarization</th>
<th>$N\bar{N}$</th>
<th>pionic</th>
<th>total</th>
<th>$V^{AN}/V^{Z-0}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{19}$F</td>
<td>0.55</td>
<td>20.03</td>
<td>1.79</td>
<td>-0.62</td>
<td>21.8</td>
<td>1.07</td>
</tr>
<tr>
<td>$^{133}$Cs</td>
<td>-0.58</td>
<td>-41.97</td>
<td>-9.90</td>
<td>0.76</td>
<td>-51.7</td>
<td>2.72</td>
</tr>
</tbody>
</table>

In the last column of Table 2 we compare the strength of the interaction between the nuclear anapole moment and an on-shell electron with that arising from $V$(electron)-$A$(nucleus) $Z_0$ exchange. For $f_\pi = f_\pi^{DDH}$ the former exceeds the latter by about a factor of three for $^{133}$Cs, a ratio smaller than that found in Ref. 8. However we expect that the inclusion of heavy meson contributions in $\tilde{H}_{PNC}$ will increase this ratio. The comparison would be more favorable for nuclei with very degenerate doublets (e.g., $^{229}$Pa) and, of course, for $T = 0$ nuclei. This suggests that radiative corrections could generate the dominant weak axial coupling in selected nuclei. Clever atomic or electron scattering PNC experiments thus might provide new measures of the effects of $H_{PNC}$ in nuclei.

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REFERENCES
STATUS OF PARITY VIOLATION IN $\bar{p} - N$ SCATTERING AT HIGHER ENERGIES

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ABSTRACT

A brief overview is given of parity violation in proton-nucleon scattering at intermediate and high energies and the status of its analysis.

DATA

Let me summarize first the relevant data according to the process and Lab kinetic energy $T_{\text{Lab}}$ or momentum $p_{\text{Lab}}$.

1. $\bar{p} - \text{H}_2\text{O}$ , $p_{\text{Lab}} = 6$ GeV/c , Argonne$^1$:
   
   \[ A_L = (26.5 \pm 6.0 \pm 3.6) \times 10^{-7} \]  

   According to ref. 2 $A_L$ in p-$\text{H}_2\text{O}$ scattering is reduced to $A_L$ in p-N scattering, appropriately averaged over p-p and p-n, by a factor ~0.6 which means that $A_L^{\text{p-H}_2\text{O}}$ in (1) translates to
   
   \[ A_L^{\text{pN}} = (45 \pm 12) \times 10^{-7} \]  

   in average p-nucleon scattering.

2. $\bar{p} - p$ , $T_{\text{Lab}} = 800$ MeV , LAMPF$^3$:

   \[ A_L = (2.4 \pm 1.1 \pm 0.1) \times 10^{-7} \]  

3. $\bar{p} - ^2\text{H}$ , $T_{\text{Lab}} = 800$ MeV , LAMPF$^4$:

   \[ A_L = (1.7 \pm 0.8 \pm 1.0) \times 10^{-7} \]  

In all these results the first error is statistical and the second systematic.

Note the large value at 6 GeV/c compared to the 800 MeV results or low energy data in p-p scattering ($A_L^{\text{pp}}(45$ MeV$)=-(1.5\pm0.2) \times 10^{-7}$) as well as in nuclei.

In the following I discuss the theoretical analysis of the first two measurements, leaving aside the p-d experiment at 800 MeV.

GENERALITIES

The starting point is

\[ A_L = \frac{\sigma_+ - \sigma_-}{\sigma_+ + \sigma_-} = \frac{\sigma_{\text{PNC}}}{\sigma_{\text{strong}}} = \frac{\int \text{Re} F_{\text{PNC}} F_{\text{strong}} d\rho_f}{\int |F_{\text{strong}}|^2 d\rho_f} \]  

where $\sigma_+$ ($\sigma_-$) refers to the cross section for incident protons with positive (negative) helicity.

The third (somewhat symbolical) representation refers to the general form of the interference measured between parity non conserving (PNC) and strong amplitudes with appropriate integration over final states usually extending over large part of phase space.
All measurements cited above are done by absorption. Thus, up to (neglected) corrections due to forward scattered particles accepted by the transmission detector\textsuperscript{6}, $\sigma_{\pm}^\text{tot}$ refers to the total cross sections $\sigma_{\pm}^\text{tot}$. Using the optical theorem one can thus write $A_L$ in the form

$$A_L^\text{tot} = \frac{\text{Im} \ F_{\text{PNC}}(0)}{\text{Im} \ F_{\text{strong}}(0)} .$$

The consequence is that $A_L = A_L^\text{tot} = 0$ if $F_{\text{PNC}}$ is calculated in Born approximation since this leads to a real amplitude. For this as well as other reasons the inclusion of strong rescattering and/or absorption in the calculation of $F_{\text{PNC}}$ is crucial.

At low energy (below pion threshold), the phases of $F_{\text{PNC}}$ are for each angular momentum state determined by elastic unitarity in terms of the known strong nucleon-nucleon scattering phase shifts and well established potential models can be used to calculate the amplitude\textsuperscript{6,7}, though many open problems remain in establishing direct contact with weak interactions at the quark level\textsuperscript{8–10}.

Contrary to this no established method exists yet for the analysis at high energy in particular for the inclusion of strong absorption and rescattering. Clearly, the problems involved concern primarily the strong and not the weak interaction underlying the PNC effects. It is emphasized that the standard techniques of perturbative QCD are not expected to work (even at the highest energies) since forward scattering and total cross sections are dominated by low momentum transfer and diffraction type processes.

### 800 MeV p-p SCATTERING

Earlier calculations, based on PNC meson exchange\textsuperscript{11–14} with PNC coupling at one MNN vertex as well as $Z^0$ exchange\textsuperscript{11} between the protons, all use more or less crude methods to include strong interaction effects. While this certainly served to estimate the order of magnitude at which an effect could be expected these calculations cannot be used anymore for quantitative analysis. They give rather widely varying results which, however, partially is due also to different choice of coupling constants.

A new calculation based on PNC meson exchange has come out last year\textsuperscript{15}. It is a non relativistic DWBA calculation using strong wavefunctions obtained from N-N potentials extended to 1000 MeV. The phases of the PNC amplitude and absorption is determined directly from the experimentally determined strong partial wave scattering amplitudes. For the PNC interaction the same strong coupling parameters and short range cutoff-formfactors are used as in the Bonn potential. The result obtained with DDH\textsuperscript{8} “best value” PNC coupling constants is $A_L^{pp} = (0.3 \to 0.4) \times 10^{-7}$. Different values in the range indicated are due to different strong potentials; the smallness of the sensitivity is due to the short range cutoffs applied.

Remarks:

- The use of formfactors reduces the effective strength of the interaction which has to be compensated by the use of enhanced effective coupling constants as exemplified by the Bonn potential analysis. For consistency correspondingly enhanced PNC effective coupling constants should be used.

- Breakup intermediate states or intermediate N-\Delta states are not included in the DWBA amplitudes. Their inclusion might be important since they admit explicit PNC one pion
exchange otherwise excluded by CP conservation in the p-p system. For a consistent analysis a three body treatment might be needed.

Nevertheless, these calculations are in the direction to go for quantitative and not merely qualitative analysis. It must be emphasized, however, that for quantitative analysis considerably more accurate data are needed.

p-N AT 6 GeV AND HIGHER

Almost all calculations and in particular all those\(^{11-14,16-18}\) based on meson or weak vector boson (\(W^\pm, Z\)) exchange between nucleons yield \(A_{\text{L}}^{pN} \lesssim 2 \times 10^{-7}\) at high energies, though results vary considerably due to different treatment of strong interaction effects and choice of parameters. There are two exceptions, however, where effects ten times larger were found: Ref. 19 based on PNC wavefunction admixture to the nucleon and ref. 17 which uses the same idea in a diquark-quark scattering model. Similar general ideas have also been considered in ref. 20 but again with much smaller results \(A_{\text{L}} \lesssim 10^{-7}\). In the following I discuss the two approaches leading to \(A_{\text{L}} > 10^{-6}\) in more detail.

PNC wave function admixture\(^{19}\)

The procedure used in ref. 19 may be summarized as follows\(^{21}\):

1. The free nucleon is assumed to be an eigenstate of an effective PNC hamiltonian

   \[ H_{\text{eff}} = \not{p} - m - c\gamma_5 \]  

   By a chiral gauge transformation \(\psi \rightarrow \exp\{\frac{i}{2}\gamma_5 \sinh c\} \psi\) parity violation is then eliminated from \(H_{\text{eff}}\) and rotated into the vertices transforming any \(\gamma_\mu\) into \(\gamma_\mu (1 + c\gamma_5)\). (This procedure was discussed already in ref. 22.)

2. An effective N-N interaction containing a \(\gamma_\mu\) (+\(\epsilon_{\mu\nu}\)) at each vertex is chosen such that its on-shell amplitude reproduces the required Regge behaviour (Pomeron exchange). Using the optical theorem it is then found\(^{19}\) that

   \[ A_{\text{L}} = \frac{|F_{\text{Lab}}|}{E_{\text{Lab}}} c \approx c \text{ at 6 GeV/c} \]  

3. The last step is to determine \(c\) which is due to weak interaction within the nucleon. The value obtained in ref. 19 is \(c \approx 2 \times 10^{-6}\) for protons and neutrons which leads to the average

   \[ A_{\text{L}}^{pN} (6 \text{ GeV/c}) \approx 2 \times 10^{-6} \]  

This could look like a resolution of the problem posed by the large experimental result. However, as pointed out in\(^{21}\), the same procedure applied to the usual parity conserving \(\rho\) and \(\omega\) exchanges at low energies yields PNC coupling constants of the form \(g \cdot c\), where \(g\) is the strong \(\rho\) or \(\omega\) coupling, and to a predicted effect at 45 MeV of \(A_{\text{L}}^{\rho\rho}(45 \text{ MeV}) \approx 15 \times 10^{-7}\) an order of magnitude larger than experiment\(^5\). It must be emphasized, however, that this approach neglects direct PNC vertex contributions. But the PNC admixture leads in the same way to a PNC pp\(\gamma\) vertex which violates current conservation and gauge invariance. Thus if there is any PNC admixture of this kind to the nucleons, an other, similarly strong vertex contribution must exist which, at least in the \(\gamma NN\) case, is able to cancel the effect of the admixture.
Quark-diquark scattering

In refs. 17 and 23 a quark diquark scattering model is presented in which weak interaction acts within the diquark. An absolute prediction of the order \((1 \rightarrow 2) \times 10^{-6}\) is given in ref. 17, which, however, is proportional to a cutoff parameter which has to be introduced and which is rather arbitrary. In a second paper\textsuperscript{23}, therefore, no absolute prediction is made but the overall normalization fixed to reproduce \(A_{L}^{pN}(6 \text{ GeV}/c)=2.6 \times 10^{-6}\). Besides accommodating a large effect at 6 GeV/c these calculations have the prominent feature that in going from 6 to 500 GeV/c (FNAL) they predict a further increase by one to two orders of magnitude.

To discuss these calculations consider Fig. 1 and denote by \(F_a\), \(F_b\) and \(F_c\) the three amplitudes corresponding to the graphs a, b and c, respectively. \(F_c \equiv F_{PNC}\) contains parity violation within the diquark while \(F_{\text{strong}} = F_a + F_b\).

![Fig. 1. Quark-Quark (a) and Quark-Diquark (b,c) contributions to nucleon-nucleon scattering amplitudes. Wavy lines represent gluon exchange and the dot an effective parity violating (weak) four fermion interaction.](image)

Then

\[
\sigma_{PNC} = \int d\rho_f \, Re \, F_{PNC} \, F_{\text{strong}}^{*} = \int d\rho_f \, Re \, F_c F_b^{*}
\]  

where integration goes over all final states and the contribution due to \(F_c F_a^{*}\) vanishes. Similarly

\[
\sigma_{\text{strong}} = \int d\rho_f \, |F_{\text{strong}}|^2 = \int d\rho_f \, \{ |F_a|^2 + |F_b|^2 \} = \sigma_a + \sigma_b
\]

since the interference between \(F_a\) and \(F_b\) vanishes in the zero quark mass limit\textsuperscript{24}. Thus in this model

\[
A_L = \frac{\sigma_{PNC}}{\sigma_a + \sigma_b}.
\]

However, in refs. 17 and 23 \(\sigma_b\) was omitted from the denominator of eq. (12) and only \(\sigma_a\) was used where \(\sigma_a=40 \text{ mb energy independent up to logarithmic terms in agreement with experiment.}\)

But this is clearly not consistent in view of the fact that \(F_b\) is the amplitude interfering with \(F_{PNC}\) in \(\sigma_{PNC}\). In addition, since \(|F_b|^2\) has the same topology as \(F_b F_c^{*}\) one could expect that \(\sigma_b\) has similar behaviour as \(\sigma_{PNC}\) and in particular increases similarly with energy. This has been verified\textsuperscript{24} in a calculation performed with the same method and parameters as in
ref. 17. Including $\sigma_b$ in eq. (12) it was found that $A_{LN}^{PN} \leq 2 \times 10^{-7}$ in the whole energy range from 6 to 500 GeV/c, independent of the cutoff factor mentioned above since it enters both in $\sigma_b$ and $\sigma_{PN}$ in the same way. It must be emphasized, that no actual prediction is claimed with this calculation since the perturbative methods used are not expected to be valid at the low momentum transfers dominating the total cross section. The calculation shows, however, that the large and rising result of the original calculation is an artifact due to inconsistent normalization of $A_L$ and that results larger than $\sim 2 \times 10^{-7}$ are hardly to be expected. It should perhaps also be mentioned that the calculations are very sensitive the choice of parameters as can be seen from the fact that for a not too drastically different choice of parameters in ref. 23 vs. ref. 17 the absolute predictions differ by about a factor of 20.

CONCLUSIONS AND OUTLOOK

1. The large 6 GeV/c result (1) is presently not understood and one should not expect from theory a further increase of $A_L$ with energy. There is a serious qualitative problem which should be settled. Whether it is experimental or theoretical is open. I think serious new theoretical as well as experimental efforts should be undertaken.

2. At intermediate energies further work must aim at quantitative comparison between theory and experiment. Besides refinements of the theory more accurate data is needed for this, however.

Measurements addressing the problem at high energy should preferably be done above 6 GeV/c, for instance at Brookhaven or KAON. At FNAL one could, due to the restricted polarized beam intensity, search only for very enhanced effects. Measurements at SATURNE (or COSY) are of course also of interest though their impact on the high energy problem is less straightforward (unless large values, are found there also). They may, however, eventually be used to extend the quantitative intermediate energy regime to higher energies. Their experimental advantage is that they can be performed below hyperon threshold thus avoiding problems with background from parity violating hyperon decays.
REFERENCES

A MEASUREMENT OF PARITY VIOLATION IN PROTON-PROTON SCATTERING AT 14.5 MeV

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ABSTRACT

At the Bonn Isochronous Cyclotron an experiment is in progress to measure the longitudinal analyzing power \( A_z = (\sigma^+ - \sigma^-)/(\sigma^+ + \sigma^-) \) in proton-proton scattering at 14.5 MeV. \( A_z \) is found to be: \( A_z = -3.1 \pm 1.6 \cdot 10^{-7} \). The phase space of the proton beam is controlled and monitored with high precision by means of secondary electron monitors, beam profile scanners and with the target detector system. A multilinear regression analyzes the noise of this system and determines at the operating point of the experiment most of the corrections.

INTRODUCTION

Parity violation is one of the signatures of the weak interaction. Since for the pure hadronic \( \Delta S=0 \) part of the weak interaction the existence of neutral currents has not been proven unambiguously, more independent and precise experiments should be performed. From an experiment determining the longitudinal analyzing power \( A_z = (\sigma^+ - \sigma^-)/(\sigma^+ + \sigma^-) \) at 45 MeV, \( A_z \) at 14.5 MeV is predicted to be \( A_z(14.5 \text{ MeV}) = -0.85 \cdot 10^{-7} \). This is just consistent with a measurement at 15 MeV with \( A_z(15 \text{ MeV}) = -1.7 \pm 0.8 \cdot 10^{-7} \).

Since the quantity to be measured is of the order of \( 10^{-7} \), it is a challenge to find all relevant systematic errors and correct for them. Usually these error contributions are determined by taking into account the amplitudes of different kinds of errors, measured during the data taking runs and the sensitivity of \( A_z \) with respect to these errors. The sensitivities are calculated from additional experiments with artificially enhanced error amplitudes, being derived from scattering experiments.

Calculating these error contributions whenever possible from measured currents with the underlying statistic of the incident proton beam, or the currents of the highly efficient detector, increases the accuracy noticeably. Thus, a multilinear regression analysis becomes feasible and allows for a correction of the measured longitudinal analyzing power. This analysis uses exclusively the statistic which is accumulated in the data taking runs and the error amplitudes that were actually present.
EXPERIMENTAL SET-UP

The polarized beam is prepared with the polarized ion source (cf. fig. 1) and injected in the center of the Bonn Isochronous Cyclotron, which accelerates the protons to 14.5 MeV. The particles are transferred to the experiment by means of the bending magnets $A_1$, $A_2$ and $A_4$. The magnet $A_4$ together with the solenoid and the wienfilter prepares the longitudinal polarization.

By means of the two secondary electron monitors (SEM) the position and angle of the proton beam is measured. The protons penetrating the horizontally or vertically split thin foils (cf. fig. 2) set free electrons at the surface, which are collected by an electrode at a potential of $+100V$. The normalized difference signal of the split foils is fed back via power amplifiers to horizontal and vertical fast steering magnets. The quality of this control system can be estimated from fig. 3.

The intensity- and polarization-distribution is measured by beam profile scanners at two positions in the beam-line. Thin (<0.1 mm) carbon strings are moved continuously in both transversal directions through the beam (cf. fig. 2). The currents of the strings are measured according to the SEM principle and protons scattered from the strings are counted by means of surface barrier detectors, being part of a polarimeter. A polarimeter between the two SEM/BPS combinations
Fig. 2 Schematic of a secondary electron monitor (SEM) and a beam profile scanner (BPS).

Fig. 3 The quality of the fast ($\tau < 100 \, \mu s$) position- and angle-control. The measured quantities refer to the SEM system of reference only.
allows to minimize quickly the transversal polarization components of the proton beam. Having passed the target system the protons are stopped in the faraday-cup.

The target detector system is shown in fig. 4. The beam passes the target filled with $\text{H}_2$ gas at 12 bar and is stopped in the faraday-cup. A special electrode suppresses electron currents from the cup. The $\text{H}_2$ gas of the detector serves as target medium as well as gas of the ionization chamber. This chamber consists of a 5 $\mu$m thick aluminum foil at 1.25 kV and four guarded collector plates connected to fast ($\tau < 10\mu$s) current to voltage converters.

Fig. 4 The target-detector-system IGEL (cf. fig. 1)

A problem arises from the finite collection time of the ions in the ionization chamber. With the aid of a simple network the time response $\tau$ of the cup current to voltage converter is slowed down appropriately to give the best normalization as shown in fig. 5.

Since the transversal analyzing power in p-p scattering at 14.5 MeV is about an order of magnitude smaller than at 45 MeV the detector could be given a four-fold symmetry. This concept facilitates to find the central axis for the beam in the detector, where the sensitivities due to position modulations are smallest. In addition the detector can be used as an on-line polarimeter for transversal polarization components of the beam, once the beam axis is fixed by means of the fast angle- and position-control. Finally further consistency tests of the data can be
performed, since $A_z$ is expected to be independent from the number and combination of ionization chamber currents taken into account.

The currents of the SEM foils, the BPS carbon strings, the ionization chamber of the detector and the cup signal are first converted to voltages for control purposes and then converted to frequencies by means of fast ($5\, \text{MHz}$) voltage to frequency (VFC) converters. The pulses of the VFC are sent via a transformer to the on-line computer. Thus the whole target detector system and the beam control is isolated from the rest of the facility.

The currents are integrated by means of two sets of 100 MHz scalers, which either count the pulses of the VFC or transfer their accumulated data to the on-line computer alternately. Since, depending on the configuration of the scalers, a systematic error of $A_z$ of up to $1 \times 10^{-7}$ is introduced, the polarization state of the beam is chosen randomly every second integrating interval. In the integrating intervals following these, the beam has the opposite polarization. An integrating interval lasts 20 ms. Thus a value for $A_z$ can be calculated every 40 ms. The digitizing accuracy achieved with this system is equivalent to the accuracy of an integrating 16 bit ADC.

In order to suppress all error contributions that are not correlated with the helicity state of the beam, the solenoid field is reversed frequently.

**DATA ANALYSIS**

The 6 dimensional phase space of the beam is controlled by means of the SEM, BPS, detector and cup signals. The high accuracy of the information derived from the SEM, detector and cup allow for a multilinear regression analysis. The error contributions, the devices are sensitive to, are for instance:

- The difference of the left/right or up/down currents of the SEM at two positions in front of the detector measure the position ($<x>,<y>$) or angle ($<\theta_x>,<\theta_y>$).  

![Fig. 5](image-url) The normalized standard deviation of $I_{\text{Det}}$ depends on the time response of the cup current amplifier.
- The sum of the left/right or up/down currents of the SEM is sensitive to correlated modulations of the beam width \(<x^2>,<y^2>\) because of the finite slit width \(<100 \mu m\) between the split SEM foils.
- The sum of the SEM currents is sensitive to coherent modulations of the proton energy \(<E>\) since the production rate of the electrons is about inversely proportional to the energy of the incident protons.
- The difference of the left/right or up/down detector currents measures the transversal polarization components \(<P_x>,<P_y>\) of the beam.
- The cup current measures beam intensity modulations \(<I>\).

A multilinear regression of these quantities determines the value of the longitudinal analyzing power for which all coherent modulations taken into account are zero.

Since relations between the different kinds of phase space modulations are taken into account, double counting of error contributions is avoided and an error reduction for the average value of \(A_z\) results (cf. fig. 6). If the multilinear regression corrects this averaged value of \(A_z\) also for zero coherent modulations, the error of \(A_z\) increases. Depending upon the magnitude of the correction, still an error reduction for the multilinear corrected value of \(A_z\) may result.

![Graph](image)

**Fig. 6** An example of an error reduction due to multilinear regression as discussed in the text.
In fig. 6 the measured standard deviation of $\Delta A_Z^{\text{meas.}}$ of $A_Z$ is normalized to the standard deviation $\Delta A_Z^{\text{theory}}$ assumed of being solely due to the shot-noise of the protons scattered in the target. This normalized quantity is plotted against its regressed value with the error reduction factor $(1 - R^2)^{1/2}$. $R$ is the multilinear regression coefficient with $-1 < R < 1$.

In case the normalization quantity is reasonable, no data point must lie inside the hatched area. Furthermore, since the straight line through the origin and the line fitted to the data points cross near the point $(1,1)$ (cf. fig. 6), the real noise contribution is only slightly greater than $\Delta A_Z^{\text{theory}}$. Therefore, noise contributions from the signal processing are considered to be negligible.

**RESULTS**

The regression quality could not be improved by nonlinear terms like $<x> <y>$ or $<x>^2 \ldots$ Table 1 shows the various error contributions. The error of the multilinear corrected $A_Z$ is smaller than the measured value. The major correction is due to the first position polarization moments $<x_P>, <y_P>, <x> <y_P>, <y> <x_P>$, which are measured by the BPS polarimeters. Therefore, it could not be included in the multilinear regression analysis.

The contributions due to other sources of errors are negligible.

*Table 1*

Summary of data, corrections and limits on systematic error contributions for $A_Z$ in units of $10^{-7}$

<table>
<thead>
<tr>
<th>&quot;Raw&quot; asymmetry $A_Z$ :</th>
<th>-6.4 ± 1.4</th>
</tr>
</thead>
<tbody>
<tr>
<td>Corrections due to correlated modulations of :</td>
<td></td>
</tr>
<tr>
<td>- Position moments$^a$ $&lt;x&gt;, &lt;y&gt;$</td>
<td>0.1</td>
</tr>
<tr>
<td>- Polarization$^a$ $&lt;P_x&gt;, &lt;P_y&gt;, &lt;x&gt; &lt;P_y&gt;, &lt;y&gt; &lt;P_x&gt;$</td>
<td>0.2</td>
</tr>
<tr>
<td>- Intensity$^a$ $&lt;I&gt;$</td>
<td>-1.1</td>
</tr>
<tr>
<td>- Width, Energy$^a$ $&lt;x^2&gt;, &lt;y^2&gt;, &lt;E&gt;$</td>
<td>-0.1</td>
</tr>
<tr>
<td>- Position polarization moments $&lt;x_P&gt;, &lt;y_P&gt;$</td>
<td>4.3 ± 1.1</td>
</tr>
<tr>
<td>Corrected asymmetry $A_Z$ :</td>
<td>-3.0 ± 1.6</td>
</tr>
<tr>
<td>Limits on systematic effects (1σ) :</td>
<td></td>
</tr>
<tr>
<td>- Electronic crosstalk</td>
<td>0.05</td>
</tr>
<tr>
<td>- Gas impurities</td>
<td>0.001</td>
</tr>
<tr>
<td>- Beta decay asymmetry</td>
<td>0.01</td>
</tr>
</tbody>
</table>
- Gasdynamic 0.001
- Background due to beam-target scattering 0.05
- Double scattering 0.037
- Uncertainty in beam polarization 0.03

*Corrected by multilinear regression*

The experiment is in progress and will be continued with an improved polarized ion source, providing an order of magnitude more current.

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POSSIBLE CEBAF MEASUREMENTS OF PARITY VIOLATION IN ELECTRON SCATTERING

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The CEBAF parity-violation program will constitute a major test of the validity of, and extensions to the standard model of electroweak unification. If extensions to the model exist the quark-electron sector, which is accessible via parity violation measurements, has a combination of coupling constants sufficiently distinct from those measurable in other sectors that a very strong model dependent lower mass limit on a higher mass $Z'$ can be obtained. The magnitude of the asymmetry in $p(\vec{e},\vec{e}')p$ or $^4\text{He}(\vec{e},\vec{e}')X$ scattering must be measured to an accuracy $\pm 3\%$ or $\pm 1\%$, respectively, in order to determine the quantity $\Delta \sin^2 \theta_w / \sin^2 \theta_w$ to $\pm 1\%$. Since these measurements will be conducted on different targets and at various $q^2$ the effects of theoretical anomalies in interpreting the results, such as uncertainties in nuclear strange quark content, the magnitude of $G_F$, etc., can be separately determined and extracted with a high degree of confidence. Further, we will test the accuracy of some of the most basic assumptions of nuclear physics, such as CVC, the lack of second class currents, and isospin invariance.

1. INTRODUCTION

A major accomplishment in physics in the last fifteen years has been the success of the electroweak model of Glashow, Weinberg, and Salam in describing the unification of the electromagnetic and weak forces. It is generally felt, however, that the model is incomplete (it contains $>20$ parameters) and new physics could appear at a mass scale as low as a few hundred GeV. The study of parity violation in electron scattering yields high-precision, low-$q^2$ measurements of ratios of coupling constants in the lepton-quark sector. The CEBAF parity-violation program will constitute a major test of the validity of, and tests of extensions to the standard model of electroweak unification. Parity violation yields a measurement of unique combinations of coupling constants predicted by the standard model such that, after radiative corrections, $\sin^2 \theta_w$ can be extracted if the standard model is rigorously correct. If extensions exist, the quark-electron sector, which is accessible via parity violation measurements, has a combination of coupling constants sufficiently distinct from those measurable in other sectors as to allow parity violation measurements to place a strong model-dependent lower mass limit of a higher mass $Z'$. In addition, parity experiments can address other questions of high physics interest. For example, we would learn whether the Standard Model is correct in the low-energy strong-coupling regime of strong interactions. Further, we can test the accuracy of some of the most basic assumptions of nuclear physics, such as CVC, the presence or absence of strange quarks, the lack of second class currents, and isospin invariance.
2. WITHIN AN \( E_6 \) MODEL

Parity violation measurements will significantly contribute to a mass limit on a \( Z' \) when considered within an \( E_6 \) framework, for extending the standard model. Marciano, Lykken, Pollock, Langacker, Lynn, and others have independently made calculations to determine the effect that a high accuracy parity violation measurement in the electron-quark sector would have on a \( Z' \) mass limit, when combined with the anticipated measurements summarized below.

- \( M_Z \) measured at SLC and LEP to \( \pm 50 \) MeV.
- \( M_W/M_Z \) measured at CERN with ACOL or Tevatron combined with the above experiments gives \( M_W \) to \( \pm 200 \) MeV.
- AFB, forward/backward asymmetry in \( e^+e^- \Rightarrow Z^0 \Rightarrow \mu^+\mu^- \) to be measured at LEP with accuracy of \( \pm 2 \times 10^{-3} \).
- \( R \), measured at the Los Alamos LCD to \( \pm 2 \times 10^{-2} \).

To extract the effect of the proposed parity violation measurements one must also assume that the top quark is found and that the Higgs mass is not too big, so standard model radiative corrections can be calculated accurately. For this discussion we will ignore atomic parity violation experiments since their theoretical interpretation is difficult. We will, also, ignore semileptonic experiments which have a theoretical uncertainty of \( \pm 3\% \). Given below are definitions of parameters associated with an \( E_6 \) extension to the Standard Model.

**Definition of Parameters \( M_{Z^0}, \theta_1, \theta_2, \) and \( \lambda \) in \( E_6 \) Model**

\[
\begin{align*}
M_{Z^0} & \equiv \text{Observed } Z^0 \text{ mass} \\
M_{Z'} & \equiv \text{Actual } Z' \text{ mass} \\
M_{Z'} & \equiv \text{Mass of Second } Z \\
\theta_1 & \equiv \text{Mixing angle between hypercharge eigenstates.} \\
\theta_2 & \equiv \text{Mixing angle between } Z^0 \text{ and } Z'. \\
\lambda & \equiv \tilde{g}_2/\tilde{g}_3 \quad \text{(The ratio of coupling constants)} \\
M_{Z^0}^2 - M_{Z'}^2 & = (M_{Z'}^2 - M_{Z^0}^2) \tan^2 \theta_2 \\
Z' \Rightarrow \tilde{Y} & = \cos \theta_1 Y' - \sin \theta_1 Y''
\end{align*}
\]

The Tevatron might produce a \( Z' \) directly if \( \lambda M_Z < 300 \text{ GeV} \). Note, however, that if \( \lambda \ll 1 \), then the, Tevatron limit may be poor at certain hypercharge mixing angles. If any of these experiments shows evidence for another \( Z \), then information concerning its manifestation in each sector becomes of **critical** theoretical importance. Figure 1 shows the results of Lykken's\(^1\) calculation to determine the effect a \( p(e^, e')p \) parity violation measurement has upon the lower mass limit for a \( Z' \), assuming nominal hypercharge mixing angle of \( \theta_2 = 0.2 \).
Figure 1. Plot of calculated Z' lower mass limits versus hypercharge mixing angle when the parity measurement is combined with anticipated experimental measurements in other sectors.

3. A POSSIBLE APPARATUS

The University of Illinois and CEBAF have examined several spectrometer designs that appear to be capable of simultaneously satisfying the design requirements of both the out-of-plane coincident electron scattering and the parity violation program. A common design for the "STAR" spectrometer, a Symmetric Toroidal ARray, has emerged. A sketch of the proposed device, configured to conduct the parity program, is shown in Figure 2.

Figure 2. The STAR spectrometer configured to conduct the parity program.
The instrument has eight azimuthally-symmetric sections, each of which consists of an inward-bending, single-focusing dipole spectrometer with point-to-parallel optics. The focal plane is parallel to the axis, making a cylindrical arrangement of the detectors convenient. The characteristics of this device most relevant to the parity program are summarized in Table I.

**Table I**

**Symmetric Toroidal ARRay Spectrometer Characteristics**

- Large solid angle (25 - 70 msr).
- Symmetric acceptance about the beam direction.
- Good scattering angle range (9° ≤ θ_e ≤ 30°).
- Extended targets (10 to 20 cm).
- Extremely high rates.
- Modest resolution of 5 x 10^{-3} obtainable by a hardware focus.

For an iron-dominated toroid with pole-gaps of 20 to 40 cm, the maximum field is about 18 kG. This is a more conservative design than an iron-free superconducting toroid with a comparable field integral. It has much better field uniformity, and better extended target optics.

### 3.1 Data Collection and Detectors

The symmetric arrangement of the eight spectrometer sectors is ideally suited for parity measurements. However, the very high count rates needed for the parity program will require the installation of specialized focal plane detectors in the STAR. Parity violation experiments require high statistics and two techniques have traditionally been employed to measure the asymmetries. Scattered particles can be counted or the analog currents, which are proportional to the number of particles, can be measured in a ring of water or lead glass Cherenkov detectors. There are advantages [+] and disadvantages [-] to each approach. These are summarized below:

**Counting Scattered Particles:**
- Dead-time effects have to be handled carefully,
+ Kinematic separation of elastic events becomes simple.

**Analog Current Integrating Techniques:**
+ There are no detector dead-time corrections,
+ the detector electronics are simpler,
- The magnetic and geometrical properties of the hardware focusing tune of the spectrometer must select elastic events.
At a continuous-beam accelerator such as CEBAF either or both techniques are suitable for these measurements. In practice, a hybrid system may actually be employed.

3.2 Polarized source and monitoring requirements

A stable, reliable, polarized source is essential for this program of research. The GaAs sources used in previous parity violation measurements have typically produced a beam with about 40% polarization. A recent CEBAF/Illinois/SLAC measurement demonstrated a practical approach to increasing the polarization to 49% by using specially-grown cathodes with a 0.2 mm active area. Therefore, the CEBAF source is now expected to achieve approximately 49% polarization. Several concepts for significantly improving the overall polarization are under serious investigation and prototyping at the University of Illinois as part of a joint Illinois/CEBAF effort to develop a polarized source for CEBAF. Several promising avenues of investigation may lead to source polarizations as high a 90%. Such an advancement would dramatically decrease the running time of these measurements.

In order to conduct a high accuracy parity violation measurement it is necessary to monitor the helicity correlated absolute beam polarization to better than 1%. Laser backscattering is the technique of choice to provide the necessary continuous monitoring for electron beam energies above 1 GeV. Its advantages, in order of importance, include:

- It is non-destructive to the beam,
- It is the only practical option for beam energies above 1 GeV.
- It is valuable resource to the entire CEBAF program,
- The laser technology is currently available.

3.3 Helicity reversal rate

It will, also, be necessary to determine an optimum helicity-reversal rate in order to minimize the effect of the intrinsic noise spectrum carried on the beam. Several philosophies exist on this subject, and the final decision can be made only after studying the properties of the actual beam. A rapid reversal of the beam helicity, in the kHz range, has the advantage that the experimental conditions cannot change significantly between helicity reversals. However, the statistical sampling is relatively poor at high reversal rates. A slower reversal rate (~1 sec) provides the advantage that a much wider bandwidth can be employed for all servo systems; this would reduce systematic errors associated with polarization reversal. The slow reversal technique is readily applicable to non-pulsed machines, such as CEBAF.

3.4 Target Considerations

A liquid H₂ or ⁴He target 10 to 20 cm in length capable of withstanding 100 μA of 1.5 to 4 GeV electrons appears feasible based on the straightforward, but non-trivial, extrapolation of technology developed for SLAC, BATES and LAMPF cryogenic target systems. These targets are limited by peak current local heating. The CEBAF cryogenic target will probably require a vertical flow path for the cryogenic liquid, and scanning of the beam to reduce local
heating effects. If the liquid $^4$He target is held at superfluid temperatures, its thermal conductivity should be extremely good.

4. INTERPRETATION OF THE HYDROGEN EXPERIMENT $p^-(e',e)p$

The parity-violating asymmetry in $p^-(e',e)p$ scattering can be expressed in terms of the weak mixing angle and a combination of weak and electromagnetic form factors for the proton. Using CVC and the assumption that the proton and neutron form a perfect isodoublet, the weak vector proton form factors can be expressed in terms of electromagnetic form factors for the proton and neutron. Consequently, a measurement of the weak mixing angle is ultimately limited in precision by our knowledge of these form factors. The parity-violating asymmetry can be expressed in terms of $q^2$, the electron energy, the electron scattering angle, and the nucleon structure functions by:

$$A_P = \frac{[d\sigma^+ - d\sigma^-]}{[d\sigma^+ + d\sigma^-]} = \frac{-Gq^2}{\sqrt{2}(4\pi\alpha)\xi} \cdot \left\{ \cos^2\theta_e \cdot \left[ (F_1\gamma_P^2 - F_1\gamma_N^2) + q^2(F_2\gamma_P^2 - F_2\gamma_N^2) \right] + \sin^2\theta_e \cdot \left[ G_M\gamma_P^2 - G_M\gamma_N^2 \right] \right\}$$

Where,

$$\xi = \cos^2\theta_e \cdot \left[ (F_1\gamma)^2 + q^2(F_2\gamma)^2 \right] + \sin^2\theta_e \cdot \left[ \frac{q^2}{4M^2} \cdot \left[ F_1\gamma + 2MF_2\gamma \right]^2 \right]$$

$G_M = F_1 + 2M F_2$

= the Electromagnetic Structure Function of the Nucleon

$e = Electron Initial/Final energy$

$q^2 = 4$ Momentum Squared

$\theta_e = Recoil Electron Scattering Angle$

$\gamma_P = Charged Structure Function$

$F_2 = Weak Magnetism Structure Function$

In these expression the superscripts $N$ and $P$ refer to the proton and neutron, respectively. Up and down arrows indicate the direction of the longitudinal polarization of the beam. By proper selection of $\theta_e$, $E'$, and $q^2$, the terms that are strongly dependent upon $\sin^2\theta_w$ can be enhanced, while the axial terms and additional terms containing poorly-determined quantities can be suppressed. This can be seen clearly in Figure 3, taken from Pollock's dissertation. To the extent that strange quarks do not contribute, $F_A^{(0),S} = 0$, so $F_A^{(0),P} = 1/2 F_A^{(0),V}$, and is known to better than 10% from charge-changing neutrino experiments. We therefore know all the form factors in the expression.
Figure 3. Parity violating asymmetry versus recoil electron lab angle for the reaction $p(e,e')p$. Shown separately are the contributions from the terms proportional to $F_1^0$, $F_2^0$, and $F_A^0$.

Clearly, kinematic conditions that allow us to minimize the impact of these uncertainties are desired. Since, $A_p$ is very sensitive to $\theta_e$, $E'$, and $q^2$, the experimental separation of the various contributions is possible, and by varying the $q^2$ at which the measurements are taken, different terms can be isolated and extracted. This provides the experiment with internal cross checks that allow detection of any unexpected physics, such as the effect of strange quarks. The next critical question is how accurately can the combinations of coupling constants $(\sin^2\theta_w)$ be extracted from the observable scattering asymmetry.

Following the work of Donnelly$^4$, et.al., the parity violating asymmetry for elastic scattering of longitudinally polarized electrons from free protons can be rewritten as:

$$A_p = \frac{d\sigma^{\uparrow} - d\sigma^{\downarrow}}{d\sigma^{\uparrow} + d\sigma^{\downarrow}} = A_0^t\left[g_V(1+\delta) - X_E + X_M\right]$$

Here $\tau = -q^2/4M_p^2 \geq 0$, $A_0 = -GM_p^2/\pi\alpha\sqrt{2} = -3.17 \times 10^{-4}$, and $g_V = 1-4\sin^2\theta_w$. The quantities $\delta$, $X_E$, and $X_M$ depend on the neutron and proton electromagnetic form factors and on the proton weak axial vector form factor as follows:

$$\delta = (1-\epsilon)^2 \sqrt{\tau(1+\tau)} \frac{G_{M_p}G_A}{F^2}$$

$$X_E = \frac{\epsilon G_{E_p}G_{E_n}}{F^2}$$

$$X_M = -\frac{\tau G_{M_p}G_{M_n}}{F^2}$$
In these expressions, $\varepsilon=\left[1+2(1+\tau)\tan^2(\theta_e/2)\right]^{-1}$, is the longitudinal polarization of the virtual photon, $\theta_e$ is the electron scattering angle in the lab frame, and $F^2=\varepsilon G_{Ep}^2+\tau G_{Mp}^2$. Typically, $\varepsilon$ is expected to be near unity and $\tau$ small, such that $\delta<<1$. Consequently, for $p(e,e')p$ scattering:

$$
\left(\frac{d\sigma}{d\Omega}\right)_{lab} \equiv \sigma M \frac{E'}{E} \left(\frac{F^2}{\varepsilon (1+\tau)}\right)
$$

where $-q^2 = 4EE'\sin^2(\theta_e/2)$

and

$$E' = \frac{E}{1+2 \frac{E}{M_p} \sin^2(\theta_e/2)}$$

Here $E'$ is the scattered electron energy, and $\sigma M = [\alpha \cos(\theta_e/2)/2E\sin^2(\theta_e/2)]^2$ is the Mott cross section. Following Donnelly, et.al. we parameterize the form factors as follows:

$$
\begin{align*}
G_{Ep}(q^2) &= G_D(q^2) = (1+4.98\tau)^{-2} \\
G_{Mp}(q^2) &= \mu_p G_D(q^2) \\
G_{Mn}(q^2) &= -\mu_n G_D(q^2) \\
G_{En}(q^2) &= +\mu_n (1+5.6\tau)^{-1} G_D(q^2)
\end{align*}
$$

In calculating cross-sections from these expressions we take $\mu_p=2.79$ and $\mu_n=1.91$, that is, we set both positive constants equal to the magnetic moments in terms of nuclear magnetons. Whereas Donnelly, et.al. assumed that $G_A$ follows the same $q^2$ behavior as $G_D$, we instead use the measured $q^2$ dependence from Ahrens et.al.. Therefore,

$$G_A(q^2) = G_A(0) (1+3.15\tau)^{-2} = 1.26 (1+3.15\tau)^{-2}$$

where $G_A(0)$ has been taken from Klemt et.al..

Table II list these various quantities as a function of $q^2$ for a beam energy of 2 GeV. We use a value of $\sin^2\theta_w=0.230$ as deduced by Amaldi et.al.. The last column in each table lists the figure-of-merit, $FOM=A_p^2 \cdot d\sigma/d\Omega$, which is inversely proportional to the running time needed to measure $A_p$ to a specified precision.
4.1 Error analysis for parity violation in p(\vec{e},e')p scattering

The our expression for the parity-violating asymmetry, \( A_p \), can be used to express \( g_V \) in terms of \( A_p \) and the form factors. This expression can be used to relate the contributions of the uncertainties in the nucleon form factors and in the asymmetry to the uncertainty in \( \sin^2 \theta_w \). Assuming that the measurements of the form factors are all uncorrelated, this uncertainty can be expressed as:

\[
\frac{\sigma_{\sin^2 \theta_w}}{\sin^2 \theta_w} = \left( \sum_{i=1}^{6} F_{Q_i}^2 \left( \frac{\sigma_{Q_i}}{Q_i} \right)^2 \right)^{1/2}
\]

where \( Q_i = G_{M_p}, G_{E_p}, G_{M_n}, G_{E_n}, G_A, \) and \( A_p \). The error coefficients, \( F_{Q_i} \), are obtained from the partial derivatives of \( g_V \) with respect to the individual \( Q_i \). Let us assume that the quantity \( F^2 \) is known to arbitrary precision since it is proportional to the cross section and requires no longitudinal-transverse separation to be measured. We also assume that the
kinematic quantities $\varepsilon$ and $\tau$ are known to arbitrary precision. Noting that $\sigma_{\text{sin}^2\theta_w}/\sin^2\theta_w = \sigma_{gV}/4\sin^2\theta_w = \sigma_{gV}$, we find:

\[
F_{G_{M_p}} = \frac{X_M}{(1+\delta)} \quad F_{G_{M_n}} = \frac{X_M}{(1+\delta)} \quad F_{G_{Ep}} = \frac{X_E}{(1+\delta)}
\]

\[
F_{G_{En}} = \frac{X_E}{(1+\delta)} \quad F_{G_A} = g_V \frac{\delta}{(1+\delta)} \quad F_A = \frac{A_p/A_0 \tau}{(1+\delta)}
\]

These coefficients are tabulated as a function of $q^2$, along with the figure-of-merit (FOM) in Table III for a beam energy of 2 GeV. Note that for $0.2 \leq q^2 \leq 0.5$ GeV$^2$/c$^2$, the figure-of-merit is relatively constant, although, the error coefficients are smaller for the lower $q^2$.

### Table III

**Error Coefficients for sin$^2\theta_w$ Measurement**

| Beam Energy 2.0 GeV, Assuming sin$^2\theta_w = 0.230 \Rightarrow g_V = 0.08$ |
|---|---|---|---|---|---|---|---|
| $q^2$ (GeV) | $d\sigma/d\Omega^* A_p^{2*10} FOM$ | $F_{G_{M_p}}$ | $F_{G_{M_n}}$ | $F_{G_{Ep}}$ | $F_{G_{En}}$ | $F_{G_A}$ | $F_A$ |
| 0.05 | 3.415 | 0.068 | 0.068 | 0.0225 | 0.0225 | 0.0004 | 0.126 |
| 0.10 | 4.867 | 0.122 | 0.124 | 0.0377 | 0.0377 | 0.0011 | 0.166 |
| 0.20 | 6.581 | 0.198 | 0.206 | 0.0545 | 0.0545 | 0.0030 | 0.231 |
| 0.40 | 6.647 | 0.271 | 0.300 | 0.0618 | 0.0618 | 0.0077 | 0.318 |
| 0.60 | 5.173 | 0.288 | 0.343 | 0.0566 | 0.0566 | 0.0128 | 0.366 |
| 0.80 | 3.695 | 0.278 | 0.358 | 0.0485 | 0.0485 | 0.0178 | 0.390 |
| 1.00 | 2.554 | 0.258 | 0.360 | 0.0402 | 0.0402 | 0.0226 | 0.400 |
| 1.20 | 1.741 | 0.234 | 0.353 | 0.0327 | 0.0327 | 0.0270 | 0.401 |
| 1.30 | 1.434 | 0.222 | 0.349 | 0.0293 | 0.0293 | 0.0290 | 0.399 |

4.2 Current knowledge of the form factors

The best available data on the various form factors, in the $q^2$ range of interest, seem to be from the following papers:

- Hanson, et.al. Phys Rev D 8, 753 (1973):

We note that the measurement of $G_{En}$ extracts this form factor from forward elastic $d(e,e')d$ scattering and uses existing data on $G_{Ep}$ to extract the result. Consequently, the
errors on these two form factors are correlated. We will neglect this here, but it should be examined in some detail in the future. These papers indicate that for \(0.2 \leq q^2 \leq 0.5\) GeV\(^2/c^2\) the electric and magnetic form factors for the proton are both known to about 3%. The neutron magnetic form factor is known to about 6%. The neutron electric form factor is known most poorly. The error in this region is roughly 40%. The error on the weak axial vector form factor is hard to determine since the authors assume a "dipole" form and fit to find the relevant mass scale. However, because of the smallness of \(\delta\) and \(g_V\), there is only weak dependence on this form factor. Assuming a 10% error on \(G_A\), Table IV tabulates the error contributions in percent for a 2 GeV electron beam:

### Table IV
Summary of Error Contributions from Form Factor Uncertainties
in % of \(\Delta \sin^2\theta_w/\sin^2\theta_w\), for \(p(e^-,e')p\) with an Incident 2 GeV Electron Beam

<table>
<thead>
<tr>
<th>(q^2(\text{GeV}/c)^2)</th>
<th>(G_{Ep})</th>
<th>(G_{Mp})</th>
<th>(G_{En})</th>
<th>(G_{Mn})</th>
<th>(G_A)</th>
<th>(\Sigma)</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.20</td>
<td>0.15</td>
<td>0.60</td>
<td>2.1</td>
<td>1.0</td>
<td>0.03</td>
<td>2.4%</td>
</tr>
<tr>
<td>0.25</td>
<td>0.18</td>
<td>0.66</td>
<td>2.2</td>
<td>1.5</td>
<td>0.04</td>
<td>2.7%</td>
</tr>
<tr>
<td>0.30</td>
<td>0.18</td>
<td>0.75</td>
<td>2.4</td>
<td>1.5</td>
<td>0.05</td>
<td>2.9%</td>
</tr>
<tr>
<td>0.50</td>
<td>0.18</td>
<td>0.85</td>
<td>2.4</td>
<td>2.0</td>
<td>0.10</td>
<td>3.2%</td>
</tr>
</tbody>
</table>

These results are essentially unchanged for a 4 GeV beam. Consequently, the majority of the cumulative error comes from our poor knowledge of \(G_{En}\), although the contributions from \(G_{Mn}\) and \(G_{Mp}\) are not negligible for a 1% measurement of \(\Delta \sin^2\theta_w/\sin^2\theta_w\). Thus, given only our present knowledge of the form factors, the accuracy in \(\Delta \sin^2\theta_w/\sin^2\theta_w\), from \(p(e^-,e')p\), would be limited to ±2-3%.

However, the neutron charge form factor, \(G_{En}\), can be extracted from polarization transfer measurements on the deuteron \(d(e^-,e'n)p\). A proposal submitted to Bates in the fall of 1988 plans to measure \(G_{En}\) at \(q^2 = 0.26\) (GeV/c\(^2\)) to ±0.016, approximately 2 to 3 times better than the current uncertainties. The theoretical errors in extracting \(G_{En}\) from the proposed measurement are very small. Arenhovel\(^8\) finds that the uncertainties due to the dependence of the \(G_{En}\) extraction on different theoretical models of the deuteron are indistinguishable out to opening angles of 30 degrees, between the outgoing neutron or proton and the direction of the momentum transferred by the electron. There is some uncertainty in the overall normalization of the theoretical calculation, possibly as large as 15%, which is likely the result of an imperfect understanding of how to do such calculations relativistically. Even though fully covariant calculations are being developed, and are expected to be available long before any CEBAF measurements are conducted, the correct normalization for this experiment can always be determined empirically by measuring proton knockout in \(d(e^-,e'p)n\) and using the results obtained from free protons to calibrate the theory. Fortunately, one of the major goals of the
CEBAF program is to measure basic form factors to better accuracies. With improvements at CEBAF, one can expect to obtain considerably smaller uncertainties in $G_{En}$ than those expected from the Bates program. If $G_{En}$ is 0.04 at $q^2 = 0.26$, then the single Bates measurement will give $G_{En}$ to about 40%; anticipated improvements at CEBAF, by a factor of five would bring this uncertainty down to 10%. Thus, uncertainties in $G_{En}$ are not expected to be the limiting systematic error in the p(e,V)e'p measurement.

4.4 Beam time requirements for the hydrogen measurement

In this section we calculate the data-taking time required for some "practical" running scenarios. The number of counts $N$ needed to measure an experimental asymmetry $A_p < 1$ to a fractional precision $a$ is $(a^2 \cdot A_p)^{-2}$. Therefore, the time $T$ needed to make this measurement is given by

$$
T = N / [\Delta \Omega \cdot d\sigma/d\Omega \cdot \xi]^{-1} = [a^2 \cdot P^2 \cdot \Delta \Omega \cdot \xi \cdot FOM]^{-1},
$$

where $P$ is the beam polarization, $\Delta \Omega$ is the solid angle subtended in the apparatus, and $\xi$ is the luminosity. Three possible kinematic scenarios for running the Hydrogen experiment have been investigated. To calculate the running time necessary to achieve a precision in $A_p = \pm 3\%$, and therefore, a measurement of $\Delta \sin^2 \theta_w/\sin^2 \theta_w$ equal to $\pm 1\%$; a beam polarization of 49%, a luminosity of $2.5 \times 10^{38} \text{cm}^{-2} \text{sec}^{-1}$ (i.e. a 100 $\mu$A beam on a 10 cm LH$_2$ target), and a solid angle of 75 msr (for the STAR at 16 degrees) is assumed. Since the solid angle of the STAR spectrometer is decreasing as the scattering angle is decreased, the running time benefits of very small angles are somewhat diminished.

**Scenario A:** $E = 4\text{GeV}$, $q^2 = 0.20\text{GeV}^2/c^2$, $LH_2$ Length = 10 cm

$\Rightarrow E' = 3.9 \text{ GeV}, \theta_e = 6.5^\circ, T = 224 \text{ hr}$

This option minimizes errors from form factor uncertainties and has a good rate. Unfortunately, it requires a high momentum and small angle for the scattered electrons. If the scattering angle is too small for a practical spectrometer, then:

**Scenario B:** $E = 4\text{GeV}$, $q^2 = 0.40\text{GeV}^2/c^2$, $LH_2$ Length = 10 cm

$\Rightarrow E' = 3.8 \text{ GeV}, \theta_e = 9.3^\circ, T = 224 \text{ hr}$

The scattering angle is probably achievable and the running time is the same. However, contribution to the final error from form factor uncertainties is somewhat larger, but still acceptable. If the momentum of the scattered electron is proves to be too large to handle in the design of the STAR.

**Scenario C:** $E = 2\text{GeV}$, $q^2 = 0.30\text{GeV}^2/c^2$, $LH_2$ Length = 10 cm

$\Rightarrow E' = 1.8 \text{ GeV}, \theta_e = 16.4^\circ, T = 871 \text{ hr}$
From the spectrometer point of view, these parameters are probably reasonable. Also, the contribution from the form factor uncertainties to our final precision is manageable. These conveniences are paid for in rate, which goes roughly like $E^2$.

5. THE HELIUM MEASUREMENT $^{4}$He($e^-,e'\gamma$)

Following the work of Walecka and Donnelly, et.al. the parity violating asymmetry for elastic scattering of longitudinally polarized electrons from $^{4}$He can be written in the following form:

$$A_{He} = \frac{[d\sigma^\uparrow - d\sigma^\downarrow]}{[d\sigma^\uparrow + d\sigma^\downarrow]} = \frac{-Gq^2}{\sqrt{2}(\pi\alpha)^2 \xi} \cdot \left[ \gamma(0)(q^2) / \gamma(q^2) \right]$$

$$= \frac{-Gq^2}{\sqrt{2}(\pi\alpha)^2 \xi} \cdot \sin^2 \theta_w$$

This is the difference of two helicity cross-sections over the sum. The result is a constant of proportionality and a ratio of two form factors. One describes the distribution of weak neutral charge over the hadronic target, and the other describes the distribution of electromagnetic charge over the target; it is the familiar charge form factor. Since, this is a $T = 0 \rightarrow T = 0$ transition, these two form factors are exactly proportional, because the currents are exactly proportional for isoscalar transitions. The ratio of form factors cancels and, the final result is a constant factor times $\sin^2 \theta_w$. Including strange quarks will result in the breakdown of this simple proportionality of currents at some small level due to extra pieces in the current coming from the strange quarks, which are isoscalar under nuclear hadronic isospin. Dubach, et.al. are calculating small corrections of order less than 1% to the simple formula above.

The elastic scattering asymmetry for $^{4}$He is, therefore, to a very good approximation simply proportional to $\sin^2 \theta_w$, and for a particular $q^2$, the asymmetry is approximately three times larger than it is in $p(e^-,e')p$ scattering. All these are good reasons to use this reaction to measure $\sin^2 \theta_w$. Although, the asymmetry is significantly larger, the cross section is 100 times smaller. So, the figure-of-merit is about 1/10 of that for $p(e^-,e')p$ scattering. This can be compensated partially by increased target length and the higher density of $^{4}$He. Furthermore, the $^{4}$He measurement has minimal resolution requirements ($\approx 5 \times 10^{-3}$), since the first excited state is at 20 MeV and is not very strongly excited.

5.1 Beam Time Requirements for the Helium Measurement

Although, the interpretation of $^{4}$He($e^-,e'\gamma$)X is very straight-forward we pay for this convenience in rate. Again, the time $T$ needed to make this measurement is given by

$$T = N / [\Delta \Omega \cdot d\sigma/d\Omega \cdot \xi]^{-1} = [a^2 \cdot P^2 \cdot \Delta \Omega \cdot \xi \cdot \text{FOM}]^{-1}.$$
Here we assume $a = 0.01$ (required to obtain $\pm 1\%$ accuracy for $\Delta \sin^2 \theta_w/\sin^2 \theta_w$ from $^4\text{He}(e^-,e')X$ scattering), $P=0.49$ (present anticipated performance of a CEBAF polarized source), $\Delta \Omega=50$ msr (approximate solid angle of the STAR at 9 degrees), and 100 $\mu$A on a 10 cm liquid Helium target. The density of liquid Helium is 7.62 lb/cu.ft. and when multiply by Avogadros Number and divided by 4 (for mass equals 4) the luminosity $\mathcal{L} = 1.2 \times 10^{38}$ cm$^2$/sec. This is about $1/2$ of what would be obtained with 100 $\mu$A on a 10 cm LH$_2$ target. Table V is a tabulation of the asymmetry, cross section, and the figure-of-merit $= A_{\text{He}^4}\cdot d\sigma/d\Omega$ in units of $10^{-8}$nb/sr, as a function of $q^2$ for 2 GeV. The $^4\text{He}$ form factor was taken from Eqn. 9 in Frosch$^{10}$, et.al. The cross section falls much faster than $q^4$, due to the $^4\text{He}$ form factor; as a result the figure-of-merit falls rapidly with increasing $q^2$, unlike the case for $p(e^-,e')p$ scattering.

Table V
Beam energy 2.0 GeV, Assuming $\sin^2 \theta_w = 0.230$

<table>
<thead>
<tr>
<th>$q^2$ (GeV)$^2$</th>
<th>$\theta$ (deg)</th>
<th>E' (GeV)</th>
<th>$F_x^2$</th>
<th>$-A_{\text{He}^4}\cdot 10^5$</th>
<th>$d\sigma/d\Omega$</th>
<th>$d\sigma/d\Omega\cdot A_{\text{He}^4}\cdot 10^8$ FOM</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.020</td>
<td>4.1</td>
<td>2.00</td>
<td>0.6209</td>
<td>0.166</td>
<td>0.2049 $\cdot 10^7$</td>
<td>561.162</td>
</tr>
<tr>
<td>0.040</td>
<td>5.7</td>
<td>1.99</td>
<td>0.3856</td>
<td>0.331</td>
<td>0.3163 $\cdot 10^6$</td>
<td>346.604</td>
</tr>
<tr>
<td>0.060</td>
<td>7.0</td>
<td>1.99</td>
<td>0.2394</td>
<td>0.497</td>
<td>0.8683 $\cdot 10^5$</td>
<td>214.074</td>
</tr>
<tr>
<td>0.080</td>
<td>8.1</td>
<td>1.99</td>
<td>0.1486</td>
<td>0.662</td>
<td>0.3016 $\cdot 10^5$</td>
<td>132.205</td>
</tr>
<tr>
<td>0.100</td>
<td>9.1</td>
<td>1.99</td>
<td>0.09225</td>
<td>0.828</td>
<td>0.1192 $\cdot 10^5$</td>
<td>81.620</td>
</tr>
<tr>
<td>0.120</td>
<td>10.0</td>
<td>1.98</td>
<td>0.05721</td>
<td>0.993</td>
<td>0.5106 $\cdot 10^4$</td>
<td>50.354</td>
</tr>
<tr>
<td>0.140</td>
<td>10.8</td>
<td>1.98</td>
<td>0.03543</td>
<td>1.159</td>
<td>0.2311 $\cdot 10^4$</td>
<td>31.019</td>
</tr>
<tr>
<td>0.160</td>
<td>11.5</td>
<td>1.98</td>
<td>0.02189</td>
<td>1.324</td>
<td>0.1087 $\cdot 10^4$</td>
<td>19.057</td>
</tr>
</tbody>
</table>

From Table V the running time required $T$(in hr) = $1.85 \times 10^5$/FOM. For 2 GeV incident electrons and a $q^2=0.1$, the FOM=81.6 and the running time becomes $T=2,300$ hours. This is about twice as long as for 2 GeV and $q^2$ around 0.3 for $p(e^-,e')p$ scattering. If a 20 cm liquid $^4\text{He}$ target is used, then only 1,150 hours are required. Since, the cross section varies so rapidly with angle, it is necessary to know, but only at the helicity reversal frequency and phase, what $q^2$ is being averaged over at the $<1\%$ level, this should be manageable. In summary, the preferred running conditions for the Helium measurement are:

**Scenario A**: $E = 2\text{GeV}, q^2 = 0.10\text{GeV}^2/c^2, LH_e^4 \text{Length} = 20 \text{cm}$

$\Rightarrow E' = 1.9\text{GeV}, \theta = 9.3^\circ, T = 1,150 \text{hr}$

**Scenario B**: $E = 2\text{GeV}, q^2 = 0.10\text{GeV}^2/c^2, LH_e^4 \text{Length} = 10 \text{cm}$

$\Rightarrow E' = 1.9\text{GeV}, \theta = 9.3^\circ, T = 2,300 \text{hr}$
6. SYSTEMATIC ERRORS ASSOCIATED WITH THE INSTRUMENTATION & BEAM

Any characteristics of the electron beam that change when the helicity is reversed may affect the scattering measurement and give rise to a spurious asymmetry that mimics the parity non-conservation asymmetry. It is, therefore, necessary to monitor the beam energy, polarization, position, intensity, and size. Each of these beam properties must be monitored continuously, and, if necessary controlled to a level sufficient to ensure that they make an insignificant contribution to the overall uncertainty relative to the level of the statistical precision of the measurement. The spurious contributions to a parity violation experiment can be parameterized as follows:

\[ A_m = A_{(p \text{ or } \text{He})} + \Delta E \partial A / \partial E + \Delta P \partial A / \partial P + \Delta r \partial A / \partial r + \Delta I \partial A / \partial I + \Delta f \partial A / \partial f + \ldots, \]

where \( A_m \) is the measured asymmetry and \( A_{(p \text{ or } \text{He})} \) is the actual asymmetry. \( \Delta E, \Delta P, \Delta r, \Delta I, \) and \( \Delta f \) are the magnitude of helicity correlated energy, polarization, position, intensity, and phase space (beam size) variations carried by the beam, respectively. The partial derivatives \( \partial A / \partial E, \ldots \) define the detector's sensitivity to the corresponding property of the beam. The goal, therefore, is to minimize all variations in the beam's properties and the detector's sensitivity to each component.

For example, the term \( \Delta r \) is expected to be very small in the CEBAF experiment since its magnitude is primarily determined by the beam diameter, and parity experiments have been conducted, at LAMPF, SIN, and BATES, using beam diameters 10 to 100 times larger than anticipated for CEBAF. The relative size of the corresponding sensitivity term, \( \partial A / \partial r \), is determined primarily by the inverse of the size of the detector times the magnitude of beam motion on the target. Again, previous experiments have employed smaller and less sophisticated apparatus, and have successfully dealt with very jittery beams. Thus, we expect the \( \Delta r \partial A / \partial r \) term to be very small in this third generation parity experiment. Owing to the quality of the detector being proposed and the extremely high quality expected for the CEBAF beam, similar results are anticipated for the other terms. Note, also, that only the helicity correlated components of the above expression contribute to systematic uncertainties.

7. CONCLUSION

A major accomplishment in physics in the last fifteen years has been the success of the electroweak model of Glashow, Weinberg, and Salam in describing the unification of the electromagnetic and weak forces. It is generally felt, however, that the model is incomplete (it contains >20 parameters) and new physics could appear at a mass scale as low as a few hundred GeV. The study of parity violation in electron scattering yields high-precision, low-\( q^2 \) measurements of ratios of coupling constants in the lepton-quark sector. The magnitude of the asymmetry in \( p(\bar{e}, e')p \) or \( 4\text{He}(\bar{e}, e')X \) scattering must be measured to an accuracy ±3% or ±1%, respectively, in order to determine the quantity \( \Delta \sin^2 \theta_w / \sin^2 \theta_w \) to ±1%. At this level, these measurements are sensitive to radiative corrections and higher-order charged and neutral-current corrections. These test assumptions about the renormalizability of the Standard Model.
Together with related measurements in the electromagnetic \((e^+, e^-)\) and lepton \((\nu, e)\) sectors, the CEBAF experiment will provide as complete a constraint as possible on higher mass Z's and other phenomena. Measurements are desirable in all sectors because theoretical extensions to the standard model allow new physics to manifest itself somewhat differently in the three sectors. The CEBAF parity-violation program will constitute a major test of the validity of, and extensions to the standard model of electroweak unification.

References

2. C. Sinclair & L. Cardman, Private Communications.
Pion-Exchange Contribution to the Parity-Violating Asymmetry in $\bar{p}p$ Scattering

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ABSTRACT

Charged pion exchange, resulting in an $n\Delta^{++}$ state, contributes positively to the weak scattering asymmetry in longitudinally-polarized $\bar{p}p$ scattering at medium energy. This pion-exchange contribution combined with the traditional $\rho$ and $\omega$ "elastic" contributions moves the total theoretical prediction into better agreement with the 800 MeV experimental data point. The pion-exchange contribution has both inelastic and elastic scattering components and is sizeable even below the pion production threshold. Strong distortions enhance the magnitude of the effect.

I. INTRODUCTION

The non-strangeness-changing non-leptonic weak interaction is still only poorly known. One reason for this is that the weak interaction effects must be seen in the background of a much larger strong interaction. Hence, experimentalists are forced to look for specific weak signatures, such as parity violation. Another reason is that experiments on nuclei of more than a few nucleons are difficult to interpret; the weak interaction quantities must be extracted amidst the complications of many-body nuclear physics.

For these reasons there has been an intense effort to measure the longitudinal (parity-violating) asymmetry, $A_L$, in proton-proton scattering.$^{1,2,3}$ This quantity is $A_L = [\sigma(+) - \sigma(-)]/[\sigma(+) + \sigma(-)]$, where $\sigma(\pm)$ are the total $\bar{p}p$ scattering cross sections for an initial beam in $\pm$ helicity states. Such experiments are very difficult at low energies,$^{1,2}$ and at medium energies$^3$ the situation is complicated further by inelastic scattering. Thus the available data on $A_L$ is rather limited, but the few points which do exist provide important constraints on the nucleon-nucleon weak interaction. (For a theoretical discussion of NN parity violation in general, we refer to Refs. 4 and 5.)

Previous theoretical work has focused almost entirely on the elastic scattering contribution to $A_L$. The weak interaction is usually described by various one-meson exchanges, which have a weak coupling at one of the nucleon lines and a strong coupling at the other. Since CP-invariance forbids a weak parity-violating $NN\pi$ vertex with a $\pi^0$ line,$^6$ there is no one-pion-exchange (OPE) contribution to the parity-violating weak asymmetry $A_L$ for proton-proton scattering in such models. Thus, most theoretical discussions of $A_L$ for $\bar{p}p$ scattering have considered only the short-range interaction
brought about by $\rho$ and $\omega$ exchange\cite{7,8,9} or, more fundamentally, by $W$ and $Z$ exchange in a quark model.\cite{10,11} There is only a rough general agreement between the models themselves and with the data. Relevant to the present calculations at medium energy, one notes that almost all theoretical values predicted for $A_L$ for elastic $\bar{p}p$ scattering are on or below the 800 MeV data point.

However, there may be contributions to $A_L$ other than $\rho$ and $\omega$ exchange. CP-invariance does allow a parity-violating $NN\pi$ vertex for charged pions. Thus, at intermediate energies (where the $I = \frac{3}{2} \Delta$-resonance can be produced), pion-exchange can be expected to play a role in $A_L$. In fact, since about half the total cross section at 800 MeV is due to pion production (which mostly proceeds through the $\Delta$), this pion-exchange contribution might be very important.

There have already been some attempts\cite{12,13} to consider the effect of the inelastic sector on the asymmetry $A_L$. The prediction by Epstein\cite{12} using a simple OPE Born approximation for the weak and strong $NN \rightarrow N\Delta$ amplitudes, which might be characterized as a long-range, inelastic contribution to $A_L$, is negative. When added to a short-range, elastic-sector prediction, this leads to a distinct disagreement of theory with the 800 MeV experimental point.

The present work\cite{14} discusses the effects of including initial and final state interactions and off-shell effects in the calculation of the $\pi$-exchange contribution to $A_L$. In contrast to Refs. 12 and 13, we find that this contribution, $A_L^{(\pi)}$, is positive. The earlier negative predictions for $A_L^{(\pi)}$ appear to be due to the use of a weak coupling constant with sign opposite to what is presently believed. Our results also differ substantially from Epstein's in other ways. We find that $A_L^{(\pi)}$ has an important elastic as well as an inelastic component. In particular, $A_L^{(\pi)}$ is sizeable even below the pion-production threshold at 290 MeV. This is consistent in spirit with the conclusion reached by Musakhanov and Podgornov\cite{15} that virtual $N\Delta$ intermediate states might be important in the context of the 15 and 45 MeV experiments.\cite{1,2} We also find that the strong-interaction distortion enhances the magnitude of $A_L^{(\pi)}$, rather than reducing it.

II. METHOD OF CALCULATION

We calculate $A_L^{(\pi)}$ using the optical theorem to find the spin-dependent total cross sections from the imaginary part of the amplitude for forward elastic $\bar{p}p$ scattering (at momentum $p$ and in helicity state $\pm$),

\begin{equation}
A_L^{(\pi)} = \frac{\text{Im} \langle p, + | T | p, + \rangle - \text{Im} \langle p, - | T | p, - \rangle}{\text{Im} \langle p, + | T | p, + \rangle + \text{Im} \langle p, - | T | p, - \rangle} (1a)
\end{equation}

\begin{equation}
= C \text{Im} \langle p, + | T_{PV} | p, + \rangle . (1b)
\end{equation}

amplitude, which is the weak transition potential $V_W$ for $NN \rightarrow N\Delta$ distorted by the strong interactions, as shown in Fig. 1. As an equation, Fig. 1 takes the DWBA form

\begin{equation}
T_{PV} = (1 - T_{NN \rightarrow N\Delta} G_\Delta) V_W (1 - G_N T_{NN \rightarrow NN}) + \text{h.c.}
\end{equation}

\begin{equation}
\equiv (1 - T_{\Delta} G_\Delta) V_W (1 - G_N T_N) + \text{h.c.} (2)
\end{equation}
Our calculation of $T_{PV}$ is carried out in momentum space, so the above expression involves two three-dimensional integrals over the relative momenta of the intermediate $N\Delta$ and $NN$ states.

Upon substituting $T = T_S + T_{PV}$ into eq. (2a), the matrix elements of $T_S$ cancel in the numerator since they are helicity independent and those of $T_{PV}$ add since $T_{PV}$ is parity-odd. Similarly, the matrix elements of $T_{PV}$ cancel in the denominator while those of $T_S$ add. The coefficient $C$ is then given by

$$
C = \left. -\frac{4\pi \, km^2 \, \frac{1}{k^2} \, \frac{1}{8\pi^2 W \, \sigma_{tot}}} \right.,
$$

where $k$ is the center of mass momentum, $W$ is the total energy, and $\sigma_{tot}$ is the total (spin-averaged) cross section which we shall take from experiment. The sign and the second factor in $C$ account for the normalization conventions we use for the invariant $T$-matrix, which is defined as in Ref. 16 [see especially Eq. (3.5)].

The use of the optical theorem has many advantages for computing $A_L^{(\pi)}$. However, another parity-violating quantity which could be measured\cite{17} is the angular-dependent asymmetry, $A_L(\theta)$, and for this there is no optical theorem approach. In this paper we restrict our discussion to $A_L$ for total cross sections.

Note that $T_{PV}$ (and hence $A_L^{(\pi)}$) scales linearly with $V_W$. Thus, the resulting value of $A_L^{(\pi)}$ is directly proportional to the values of the weak coupling constants used. In practice, for the parity-violating $NN \rightarrow N\Delta$ transition, the dominant weak vertex in the potential $V_W$ is the one for $NN\pi$. The parity-violating $\Delta N\pi$ vertex is, from angular momentum considerations, $D$-wave rather than $S$-wave and is therefore small. This argument has been borne out explicitly in the calculations of Gorman and Mckellar.\cite{13} They include the weak $\Delta N\pi$ vertex and that changes the original Epstein estimate for $A_L^{(\pi)}$ by only a few percent. Therefore, the scale of the pion-exchange contribution to $A_L$ is set by $f_\pi$, the weak $NN\pi$ coupling constant. Our results for $A_L^{(\pi)}$ are linearly proportional to $f_\pi$. Estimates and discussions of its value are given in Refs. 5, 18, 19.

For $p\bar{p}$ scattering by pion exchange, the terms $V_W$, $V_WG_NT_N$, and $T_NG_NV_W$ in Eq. (3), will not contribute to $A_L^{(\pi)}$ since they involve only the $NN \rightarrow NN$ matrix elements of $V_W$, which vanish. The calculation of $\text{Im } T_{PV}$, and hence $A_L^{(\pi)}$, can then be broken down into two separate contributions. These correspond to the one-loop ($T_{\Delta}G_NV_W$ plus $V_WG_NT_{\Delta}$, two terms which are equal by hermiticity) and two-loop integrations ($T_{\Delta}G_NV_WG_NT_N$ plus $T_NG_NV_WG_{\Delta}T_{\Delta}$, also equal by hermiticity).
The one-loop term, which has the weak $NN \rightarrow N\Delta$ transition potential combined with a strong $N\Delta \rightarrow NN$ transition (or the inverse of this process), is similar in nature to that which Epstein calculated\(^\text{12}\) using an OPE Born approximation for the strong $NN \rightarrow N\Delta$ amplitude. Actually, Epstein assumed production of a physical $\Delta$ (with a sharp, fixed mass), which is presumed to decay into a final-state pion and nucleon. In contrast, we treat the $\Delta$ as a resonating $\pi N$ state which can be virtual as well as physical. Thus, in our calculation, the contribution of the $N\Delta$ intermediate states may be non-zero at all energies. The evaluation of these integrals requires half-off-shell matrix elements of both the strong amplitude and the weak potential.

The two-loop contribution is a double integral over the intermediate state momenta. It involves both initial- and final-state strong distortions of the weak $NN \rightarrow N\Delta$ transition potential. The integrands are composed of half-off-shell strong interaction $T$-matrices together with fully-off-shell weak transition potential matrix elements.

III. WEAK TRANSITION POTENTIAL

In the calculations presented here, we use the strong-interaction $T$-matrices ($T_\Delta$ and $T_N$ above) and propagators ($G_\Delta$ and $G_N$ above) directly from the calculation described in Ref. 16. It is imperative, then, that the calculation of the weak potential $V_W$ be carried out using a set of conventions consistent with Ref. 16.

We obtain $V_W$ for $NN \rightarrow N\Delta$ from a calculation of the OPE Feynman graph corresponding to the center part of Fig. 1. This is evaluated in analogy with the strong-interaction Born term shown in Fig. 9 of Ref. 16. The spin values of the fermion lines in our case are $s_1 = s_2 = s'_2 = \frac{1}{2}$ and $s'_1 = \frac{3}{2}$. The Lagrangians used are

$$\mathcal{L}_{NN\pi}^S = ig_N \bar{\psi} \gamma_5 (\vec{\tau} \cdot \vec{\phi}) \psi,$$

$$\mathcal{L}_{NN\pi}^W = (f_\pi/\sqrt{2}) \bar{\psi} (\vec{\tau} \times \vec{\phi})_3 \psi,$$

$$\mathcal{L}_{\Delta N\pi}^S = (g_\Delta/\mu) \bar{\psi} \mu (\vec{T} \cdot \partial \mu \vec{\phi}) \psi + \text{h.c.},$$

where $\mu$ is the pion mass and superscripts $S$ and $W$ indicate strong (parity-conserving) and weak (parity-violating) couplings, respectively. These Lagrangians in fact define the coupling constants, which are real. Our conventions have $\gamma_5$ with +1's in the off-diagonal elements.

To evaluate the loop integrals discussed in the previous section, we need the $I=1$, $M_I=+1$ matrix element of $V_W$ (since the initial and final states used in applying the optical theorem are $pp$ states with those isospin quantum numbers and the strong-interaction distorting factors preserve isospin). This is

$$\langle N\Delta, I = 1, M_I = +1 | V_W | NN, I = 1, M_I = +1 \rangle = -\frac{\sqrt{3}}{2} \langle n\Delta^{++} | V_W | pp \rangle,$$

where we have used the fact that $\langle p\Delta^+ | V_W | pp \rangle$ vanishes for $\pi$-exchange because of the $(\vec{\tau} \times \vec{\phi})_3$ structure at the parity-violating $NN\pi$ vertex.\(^\text{6}\) Thus we only need to calculate the weak $pp \rightarrow n\Delta^{++}$ transition.

This matrix element, $\langle n\Delta^{++} | V_W | pp \rangle$, is composed of several factors. The pion propagator provides, in the static approximation, a factor of $(q^2 + \mu^2)^{-1}$, where the momentum transfer $q = p + p'$, according to the labeling of momenta shown in Fig. 3.
Following Eqs. (4.20) and (5.7) of Ref. 16, the upper, parity-conserving \( \pi^+ p \to \Delta^{++} \) vertex has a spin and momentum dependence of the form\(^{20}\)

\[
i(g_\Delta/\mu)N(E, m_N)N(E, m_\Delta) \left( \begin{array}{c|c} s_1 & 1 \\ \hline m_1 & \lambda \\ \end{array} \right) \left( \begin{array}{c} s'_2 \\ m'_2 \\ \end{array} \right) V^*_\lambda, \quad (6a)
\]

where the quantity in parentheses with the vertical bar is a Clebsch-Gordan coefficients and \( V \) is a covariant three-vector specifying the relative momentum between the pion and the nucleon. A summation over \( \lambda \) is understood and we have neglected some higher order terms in \( p/m_N \) or \( p/m_\Delta \). In the non-relativistic limit the spinor normalization factors go to 1 and, neglecting terms of order \( \mu/m_N \), the “magic vector” \( V \) reduces to \(-q\).

At the lower parity-violating \( p \to n\pi^+ \) vertex, the spin and isospin dependence reduces to the simple factor\(^{21}\)

\[
\left( f_{\pi}^{p \to n\pi^+} / \sqrt{2} \right) \delta_{s_1 s_2} \delta_{m_1 m_2} = -i f_{\pi} \delta_{s'_1 s'_2} \delta_{m'_1 m'_2}, \quad (6b)
\]

where the relation between the specific coupling, \( f_{\pi}^{p \to n\pi^+} \), and the general weak coupling, \( f_{\pi} \), is given by the \((\bar{\tau} \times \bar{\phi})_3\) factor in the Lagrangian of Eq. (5b). Finally, to simulate the finite size of the nucleon and delta we also include a form factor \( \Lambda^2/(q^2 + \Lambda^2) \). Combining all these factors then gives the plane-wave representation of the weak \( pp \to n\Delta^{++} \) potential.

The partial wave decomposition of \( V_W \) is then done following the methods described in detail in Sec. 5 of Ref. 16. Because of the static propagator and the replacement of \( V \) by \(-q\), the angular integrals lead to an expression in terms of Legendre functions of the second kind. The partial wave amplitude for \( V_W \) shows explicitly that \( L' = L \pm 1 \), i.e., that parity is not conserved. Because the initial and final \( pp \) states have \( I = 1 \) and the Pauli principle requires that \( L + S + I \) be odd, the initial and final \( pp \) states have \( S_i = 0 \) and \( S_f = 1 \), respectively (or vice versa). The only \( pp \) states that have both triplet and singlet spin states are those for which \( J \) is even. As will be seen later, at intermediate energies \( A_L^{(\pi)} \) is in fact dominated by the \( J = 2 \) partial waves. This happens because the most favorable strong-interaction amplitude from the standpoint of an angular momentum barrier argument is the \( N N(1D_2) \to N\Delta(5S_2) \) transition. It is this transition that dominates the inelasticity in \( pp \) collisions at intermediate energies. It is interesting that, because of the \( J = \text{even} \) requirement, the next most inelastic partial wave, \( NN(3F_3) \), doesn’t contribute to \( A_L^{(\pi)} \).

In the results we present below we take the strong coupling constants from Ref. 16: \( g_N = 13.42 \) and \( g_\Delta = 2.43 \). The mass of the \( \Delta \) resonance is \( m_\Delta = 1.236 \text{ GeV} \) and we will generally choose the form factor cutoff as \( \Lambda = 1 \text{ GeV/c} \). For the somewhat controversial\(^{5}\) weak coupling \( f_{\pi} \) we use the DDH “best value”,\(^{18}\) which in our conventions is \(-4.56 \times 10^{-7} \). As emphasized earlier, our results scale linearly with \( f_{\pi} \) and thus can be adjusted accordingly.

**IV. MODEL FOR STRONG DISTORTION**

The remaining factors in the one-loop and two-loop integrals come from the strong interactions, and we will use the unitary three-body theory for strong \( NN \to NN \) and \( NN \to NN\pi \) reactions presented in detail in Ref. 16. The partial wave amplitudes
Fig. 2. Predicted values of $A_L^{(\pi)}$ in units of $10^{-7}$ for our model and for Epstein\textsuperscript{12} (scaled by $-\frac{1}{2}$ to correspond to the same value of $f_\pi$ as used here, the DDH “best value”).

for $T_S$ are found by solving coupled three-body integral equations. The model provides a unified description\textsuperscript{16} of the strongly-interacting $NN$ and $NN\pi$ systems that respects two- and three-body unitarity. The two-body inputs for solving these three-body equations are the $\pi N$ interactions in the $P_{11}$ and $P_{33}$ channels. This model has been extensively compared with experiment. It predicts the inelastic cross section\textsuperscript{22} within 10\% and, despite the simplicity of the input, provides a reasonable description of the $NN$ elastic scattering phase parameters apart from those of the lowest-$L$ partial waves.\textsuperscript{23} It also gives predictions for exclusive and inclusive $NN \rightarrow NNN\pi$ cross sections and spin observables that are generally in qualitative agreement with the available data.\textsuperscript{24}

V. RESULTS AND DISCUSSION

Our results for $A_L^{(\pi)}$ are shown in Fig. 2. The contribution is positive and of a similar size as the $\rho$- and $\omega$-exchange elastic contributions. In Fig. 2 and below we have used Desplanques, Donoghue and Holstein’s “best value”\textsuperscript{18} of the weak $NN\pi$ coupling constant, $f_\pi = -4.6 \times 10^{-7}$. Not shown in Fig. 2 are results from our model at higher energies (some well above the region of its validity). There are no surprises here; $A_L^{(\pi)}$ falls smoothly to about $+0.4 \times 10^{-7}$ by $T_{Lab} = 2$ GeV.

The predicted value at 800 MeV is $A_L^{(\pi)} = +1.5 \times 10^{-7}$. Together with the Henley-
Krejs calculation\(^7\) of the "elastic" contributions from \(\rho\) and \(\omega\) exchange (also positive), it brings the total theoretical \(A_L\) into good agreement with the experimental value measured at LAMPF. On the other hand, combined with Oka's results\(^9\), the total predicted \(A_L\) is somewhat larger than experiment. If \(f_\pi\), as Adelberger and Haxton argue from the \(^{15}\text{F}\) \(\gamma\)-ray data\(^5\), is smaller than the DDH "best value", then in both cases the combined prediction for \(A_L\) would be compatible with experiment.

Figure 2 also shows how the result at medium energies is dominated by the \(J = 2\) partial wave. The reason for this is that the \(N\Delta\) partial wave with the smallest angular momentum barrier occurs in the \(^5S_2\) channel. To indicate the model dependence of our prediction, we also show in Fig. 2 how the \(J = 2\) contribution changes if we use \(J = 2^+\) partial wave amplitudes for the strong \(NN \rightarrow N\Delta\) transition modified by heavy-boson-exchange\(^{25}\) to reproduce the experimental \(^1D_2\) inelasticity.

Note also that \(A_{L}^{(\pi)}\) does not go to zero at the pion-production threshold, because the \(N\Delta\) intermediate state in Fig. 1 will contribute even if it is virtual, i.e., unable to decay into a physical \(NN\pi\) state. Thus \(A_{L}^{(\pi)}\) gets contributions from elastic and as well as inelastic \(pp\) scattering. Below pion production threshold the \(J = 2\) contribution is small and the \(J = 0\) part starts to dominate. The value of \(A_{L}^{(\pi)}\) changes sign between 100 and 200 MeV. This is consistent with earlier calculations for a two-pion-exchange parity-violating potential.\(^{26,27}\) At these energies and lower, the strong interaction model of Ref. 16 may be inappropriate. However, using more realistic \(0^+\) and \(0^-\) phase shifts, we find \(A_{L}^{(\pi)}\) is about \(-0.1 \times 10^{-7}\) at 50 MeV.

Our result for \(A_{L}^{(\pi)}\) at 230 MeV may change the interpretation of the TRIUMF experiment now in preparation.\(^{17}\) The motivation for measuring at this energy comes from the presumption that the weak transitions only involve \(\rho\) and \(\omega\) exchange. In that picture, the usually dominant \(J = 0\) contribution is expected to vanish at 230 MeV. The strongest remaining contribution, that from \(J = 2\), is thought to be mostly due to \(\rho\) exchange, and the experiment is therefore sensitive to the weak coupling constant \(h_{NN\rho}\). However, since the \(\rho\)-exchange contribution actually occurs in combination with \(\pi\) exchange, the extraction of \(h_{NN\rho}\) may not be so clean. The precise contribution of \(A_{L}^{(\pi)}\) is hard to specify in view of the model dependence discussed above.

Figure 3 shows the separate contributions to \(A_{L}^{(\pi)}\) from the one-loop and two-loop pieces. The two-loop piece is itself decomposed into a dispersive part and a "pole" part with the latter coming from the part of the integration that puts the intermediate \(NN\) state on its mass shell. (This pole piece is one of the elastic \(NN\) contributions.) An important point about the curves in Fig. 3 is that the two-loop contribution is, above threshold, of the same sign as the one-loop graph. That is, the strong distortion gives a sizeable enhancement over the one-loop graph. This is at variance with the way Epstein\(^{12}\) estimated these distortion effects; he used a Watson-type factor, \(\sqrt{\eta}\exp^{18}\), which can only reduce the overall magnitude of his effect.

Another model dependence in our predictions is the dependence on the weak form factor cutoff parameter, \(\Lambda\). Depending on what \(\Lambda\) is taken, \(A_{L}^{(\pi)}\) can change by as much as a factor of two. However, since the choice of \(\Lambda = \infty\) is unrealistic, the form factor dependence of our predictions is only moderate. We have used the value \(\Lambda = 1\) GeV/c in Figs. 2 and 3.

To conclude this section, we remark that there may be other pion-exchange contributions to \(A_{L}^{(\pi)}\). Possible crossed two-pion exchange diagrams which might contribute
to $V_W$ would include $N\Delta\pi\pi$ or $NN\pi\pi$ intermediate states and are probably small. One effect, which ought also be included in a pion-exchange model, is a two-loop graph in which both intermediate states are $N\Delta$. In this case the weak interaction could be either at the $NN\pi$ vertex (i.e., also proportional to $f_\pi$) or at the $\Delta\Delta\pi$ vertex (which can now be in a S-wave). One should also consider the heavier $\rho$ and $\omega$ exchanges coupling to the $N\Delta$ states as well. These possibilities are presently under investigation.

VI. CONCLUSIONS

From this consideration of the pion-exchange contribution to the parity-violating asymmetry in $\bar{p}p$ scattering, we have come to the following conclusions:

1) $A_L$ from $\pi$-exchange is positive and comparable in size to the earlier predictions for $A_L$ from $\rho$ and $\omega$ exchange.

2) Assuming the DDH “best value” of $f_\pi$, the theoretical prediction for $A_L$ is in reasonable agreement with the experimental $A_L$ at 800 MeV when we combine our result with the “elastic” $\rho$- and $\omega$-exchange contribution of either Ref. 7 or Ref. 9. One should bear in mind the large error bars on the experimental value, the wide range allowed for $f_\pi$, and the model dependences of both our result and the “elastic” calculations. The possibility remains that the value of $f_\pi$ is indeed smaller than the DDH “best value”, as suggested in Ref. 5 and Ref. 19.

3) We find both elastic and inelastic contributions to $A_L^{(\pi)}$. One consequence of this is that $A_L^{(\pi)}$ is non-zero and sizeable even below pion-
production threshold. This complicates the analysis of the 230 MeV experiment underway at TRIUMF.

4) The two-loop contribution is of the same sign and size as the one-loop contribution. Thus, here, distortion enhances the magnitude of $A_L^{(\pi)}$ (in contrast with a Watson-type treatment, which reduces it).

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B. Desplanques, J. F. Donoghue, and B. R. Holstein, Ann. Phys. (N.Y.) 124, 449 (1980), henceforth referred to as “DDH”. These authors use a \( \gamma_5 \) which is the negative of ours. This can be verified by following the algebra from their Lagrangians corresponding to our Eqs. (5a) and (5b) to the weak PV potential displayed in their Eq. (119). Thus, the value of \( f_\pi \) we use here is the negative of the DDH value. We thank the referee for pointing out this difference between our conventions.


The strong interaction vertices of Ref. 16 [eqs. (4.18), (4.19), (4.20), (5.5), and (5.6)] should include a factor of \( i \) to be consistent with the \( G_N \) and \( G_\Delta \) used in Ref. 16 and the sign used in the optical theorem to obtain \( \sigma_{\text{tot}} \) in our eq. (4) [see eq. (3.5) of Ref. 16]. The computer code that calculates the strong amplitudes has always had this correction properly taken into account.

To be certain that we have the correct relative sign here, we have compared the ratio of the weak to strong \( np \rightarrow pn \) potentials calculated using both our conventions and those of DDH.


RELATIVISTIC AND INELASTIC EFFECTS IN THE PROTON-PROTON PARITY-VIOLATING INTERACTION

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ABSTRACT

Relativistic effects are described in a meson-exchange model of the proton-proton interaction. Relativistic corrections to parity-violating observables depend on the consistent derivation of strong and weak potentials from meson theory. Parity-violating inelastic effects are discussed.

INTRODUCTION

The development of experimental methods to measure parity violation in hadronic systems provides new opportunities for studying the hadronic interaction. While the parity-violating effect itself is small, the ability to measure parity violation in hadronic systems makes it possible to probe hadronic states in manner distinct from other more established means. One immediate application of the new physics is to test conventional meson theory. Can parity-violating interactions be included in the familiar meson-exchange picture to explain parity-violating effects? Or will contradictions with experiment lead to other descriptions of hadronic physics? As more measurements are completed, significant conclusions to these questions are expected to follow.

To test meson theory at energies above 200-300 MeV, calculations must include relativistic and inelastic effects. Strong distortions cannot be neglected, as these modify the energy dependence of parity-violating observables. A recent calculation by Driscoll and Miller includes strong distortions without nonrelativistic approximations. In this calculation, relativistic corrections in the strong interaction model produce increases in the parity-violating observables larger than the increases due to relativistic factors in the weak potential alone. At 800 MeV, the predicted analyzing increases is 44 per cent larger than in the nonrelativistic calculation.

Parity-violating inelastic processes may also have a significant bearing on parity-violating observables above 300 MeV. Calculations of inelastic effects have been reported by Epstein; Gorman and McKeller; and Silbar, Kloet, Kisslinger, and Dubach, but, except for Reference 5, these calculations neglect the strong distortions. Some recent work on inelastic effects is described below.

Meson Theory with Strong and Weak Interactions

To develop a consistent model for the strong and weak hadronic forces, we must derive both forces from a common set of assumptions. This can be accomplished by first adopting an established model for the strong force, such as one of the Bonn meson-exchange models. Weak forces are introduced by including weak, parity-violating meson-nucleon interactions in the original model. In calculations, only parity-violating graphs with one parity-violating vertex need be considered. The remaining strong vertices and propagators are defined in the strong model. It is essential, in our view, to derive both strong and weak forces in such a consistent
manner, using the same strong interactions, with identical strong coupling parameter values and form factors in both strong and weak potentials. Also, any approximations (e.g., the static approximation) assumed for the strong force should be applied in the same manner to the weak. It is the assumptions of meson theory that we wish to test, and these assumptions affect both the strong and weak forces. If the strong and weak potentials are derived using different parameter values, or using other conflicting assumptions, the calculation would have unclear implications.

A number of Bonn model potentials strong potentials, $V_s$, are obtained using various approximations. In each model, there is an assumed set of meson coupling parameters and dipole or monopole form factors. Some of these parameters and form factors are adjusted to fit empirical NN data. The strong potential $V_s$, in any of these models, can be used to calculate off-shell values of the strong $T$ matrix, $T_s$, or the strong reaction matrix, $R_s$. It is these values that determine the strong wave functions, or the “strong distortions”.

Two momentum-space models are developed in alternative calculations of the parity-violating observable. These extend two Bonn models. One, OBEPQ (One Boson Exchange Potential in Q space), is derived with the Blankenbecler-Sugar approximation, which produces an energy-independent, static potential. The other, OBEPPT (One Boson Exchange Potential based on Time-ordered perturbation theory), is nonstatic and energy-dependent. For details, the reader is referred to Reference 2. A third calculation, consistent with the nonrelativistic Bonn OBEPRT (One-Boson-Exchange Potential in R Space) model, is developed in Reference 1. For each of these three models, a parity-violating potential, $V_W$, is derived using the same strong meson-nucleon interactions, the same approximating assumptions, and the same fitted parameter values as assumed in the strong-model potential, $V_s$.

The rho and omega interactions are,

$$
\mathcal{H}_{pp}^{\rho} = g_{\rho} \overline{\psi} \left( \gamma_\mu + \frac{\mu_\nu}{2m_N} \sigma_\mu \kappa_\nu \right) \phi_\rho \psi
$$

$$
\mathcal{H}_{pp}^{\omega} = g_\omega \overline{\psi} \left( \gamma_\mu + \frac{\mu_\nu}{2m_N} \sigma_\mu \kappa_\nu \right) \phi_\omega \psi
$$

The weak parameter values $h_{pp}^{\rho} = -15.47 \times 10^{-7}$ and $h_{pp}^{\omega} = -3.04 \times 10^{-7}$ are the correct isospin combinations, for proton-proton states, of the Desplanques-Donoghue-Holstein (DDH) parity-violating coupling parameters. The theoretical uncertainties of these parameters are rather large: $h_{pp}^{\rho}$ and $h_{pp}^{\omega}$ may vary by ±200 per cent. For definiteness, we use only the given “best values”. In the equations above, we assume Bjorken-Drell conventions. Note that Reference 9 uses the opposite sign convention for $\gamma_5$.

In the static model (OBEPQ), $n_\rho = 2$ and $n_\omega = 2$ (dipole form factors). The strong coupling and cutoff parameter values from OBEPQ are $g_\rho^2/4\pi = 0.90$, $f_\rho/g_\rho = 6.1$, $g_{\omega}^2/4\pi = 24.5$, $f_\omega/g_\omega = 0.0$, $\Lambda_\rho = 1850$ MeV, and $\Lambda_\omega = 1850$ MeV.

The static potentials have “minimal relativity” factors $\sqrt{E_q/m_N}$ and $\sqrt{E_{q'}/m_N}$, where $E_q = \sqrt{q^2 + m_N^2}$, $E_{q'} = \sqrt{q'^2 + m_N^2}$, and $q$ and $q'$ are the initial and final momenta. These bring the scattering equation, Equation 1 or 2, into an effective, nonrelativistic form, $G^{(1)} = (E - H_0 + i\epsilon)^{-1}$ with $H_0(k) = k^2/m_N$.

In the energy-dependent model, $V_W$ is derived in time-ordered perturbation theory consistently with OBEPPT. In this case, $V_S$ and $V_S$ contain energy-dependent meson propagators and the nucleon propagator used with OBEPPT is explicitly relativistic: $G^{(1)} = (E - H_0 + i\epsilon)^{-1}$, with $H_0(k) = 2\sqrt{k^2 + m^2}$. 

$$
\mathcal{H}_{pp}^{\rho} = -h_{pp}^{\rho} \overline{\psi} \gamma_\mu \gamma_5 \psi \phi_\rho
$$

$$
\mathcal{H}_{pp}^{\omega} = -h_{pp}^{\omega} \overline{\psi} \gamma_\mu \gamma_5 \psi \phi_\omega
$$
In this section, we develop the formalism for calculating the parity-violating amplitudes. The $T$ and $R$ operators are the solutions of the Lippmann-Schwinger equations,

$$T = V + V G^{(+)} T,$$
$$G^{(+)} = (E - H_0 + i\epsilon)^{-1}$$

(1)

$$R = V + V G^{(s)} R,$$
$$G^{(s)} = P(E - H_0)^{-1}$$

(2)

with outgoing- and standing-wave boundary conditions for $T$ and $R$ respectively. The Coulomb force, considered in Reference 1, is neglected here. In operator equations, the arguments of energy-dependent operators $T(E)$, $R(E)$, $V(E)$, $G^{(+)}(E)$, and $G^{(s)}(E)$ are generally suppressed. Sums and integrals for intermediate-states quantum numbers are implied. Note that $V$ is actually energy-independent in the static and nonrelativistic models.

We assume $V$ is Hermitian (elastic) and time-reversal invariant. Then, Equation 2 is time-reversal symmetric, and the solution $R$ is Hermitian and time-reversal symmetric. In a representation where $V$ is symmetric and real (below inelastic threshold), $R$ is also. Notice that $T$ and $R$ are related by,

$$T = R - i\pi\delta(E - H_0)T.$$

(3)

In the presence of parity-violating interactions, $V$ acquires a small, weak correction: $V = V_S + V_W$. Neglecting second-order terms, which are of no consequence here, the parity-violating corrections to $T$ and $R$, derived from Equations 1 and 2, are

$$T_W = (1 + T_S G^{(+)} )V_W (1 + G^{(+)} T_S)$$
$$R_W = (1 + R_S G^{(s)} )V_W (1 + G^{(s)} R_S),$$

(4)

(5)

$T_S$ and $R_S$ are the zero-order solutions of Equations 1 and 2, i.e., with $V = V_S$. $T_W$ and $R_W$ are related by,

$$T_W = (1 - i\pi T_S \delta(E - H_0)) R_W (1 - i\pi \delta(E - H_0)T_S)$$

(6)

In numerical calculations, it is simpler to work with $R_W$ than $T_W$ since Equations 2 and 5 are real. Equation 6 is not needed in its present form.

To calculate the scattering observables, we require the values of $T$ or $R$ on the energy shell, where $H_0(k) = E$. In the partial-wave helicity representation\textsuperscript{10}, we write these with lower-case letters $t$ and $r$:

$$t^i_{\lambda'_1\lambda'_2;\lambda_1\lambda_2}(k) = -\pi \rho_E (k \lambda'_1 \lambda'_2 | T^i(E) | \lambda_1 \lambda_2 k)$$
$$r^i_{\lambda'_1\lambda'_2;\lambda_1\lambda_2}(k) = -\pi \rho_E (k \lambda'_1 \lambda'_2 | R^i(E) | \lambda_1 \lambda_2 k)$$

(7)

(8)

Here, $j$ is the total angular momentum, and $\lambda_1, \lambda_2, \lambda'_1$, and $\lambda'_2$ are the initial and final helicities. For the static (OBEPQ) and nonrelativistic (OBEPR) models, $\rho_E = m_N k/2$; for the energy-dependent model (OBEPPT), $\rho_E = E k/2$. In the static model, as mentioned above, the minimal relativity factors in the effective potential bring the scattering equation, Equation 1 or 2, into nonrelativistic form, i.e., with a nonrelativistic $G$.

With the normalization chosen in Equations 7 and 8, the unitary $s$ matrix is

$$s^j_{\lambda'_1\lambda'_2;\lambda_1\lambda_2}(k) = \delta_{\lambda'_1\lambda_2} \delta_{\lambda_1\lambda_2} - 2it^j_{\lambda'_1\lambda'_2;\lambda_1\lambda_2}(k).$$

(9)
and Equation 3 reduces, on shell, to

\[ t'_{\lambda_1'\lambda_2';\lambda_1\lambda_2}(k) = r'_{\lambda_1'\lambda_2';\lambda_1\lambda_2}(k) + \sum_{\lambda_1''\lambda_2''} ir'_{\lambda_1'\lambda_2';\lambda_1''\lambda_2''}(k)t'_{\lambda_1''\lambda_2'';\lambda_1\lambda_2}(k). \] (10)

By suppressing sums and indices, Equations 9 and 10 obtain the simple forms of matrix equations, \( s = 1 + 2it \) and \( t = r + irr \). Using straightforward matrix algebra, \( t \) can be solved from \( r \), or visa-versa.

Now, with parity-violating interactions, \( s, t, \) and \( r \) acquire weak terms, as before. The parity-conserving parts, \( s_S, t_S, \) and \( r_S \), are (to first order) independent of parity-violating corrections and are accurately known from conventional scattering data. For \( t_S \), we use Arndt's empirical values.11 To calculate parity-violating observables, there remain the parity-violating \( s \)-matrix elements, \( s_W = 2it_W \), where \( t_W \) is given by

\[ t_W = (1 + it_S)r_W (1 + it_S). \] (11)

Only \( r_W \), on the right-hand side, is derived from a model. Equation 11 is equivalent to Equation 6 on the energy shell. Alternatively, Equation 11 follows from Equation 10, by treating \( t_W \) and \( r_W \) as differentials of \( t \) and \( r \).

The partial-wave helicity representation matrix elements for \( V_W \) are obtained with the transform,

\[ \langle q' \lambda_1' \lambda_2' | V_{W,\alpha} | q \lambda_1 \lambda_2 \rangle = 2\pi \int_{-1}^{1} d(cos \theta) d^j_{\lambda_1}(\theta) \langle q' \lambda_1' \lambda_2' | V_{W} | q \lambda_1 \lambda_2 \rangle, \] (12)

where \( q'q \cos \theta = q' \cdot q, \lambda' = \lambda_1' - \lambda_2', \lambda = \lambda_1 - \lambda_2, \) and where \( d^j_{\lambda}(\theta) \) are Jacobi polynomials. There are sixteen possible combinations of initial and final state helicities. For identical particles, the exchange symmetry implies (suppressing \( q' \) and \( q \))

\[ \langle \lambda'_1 \lambda'_2 | V^j_S | \lambda_1 \lambda_2 \rangle = +\langle \lambda'_1 \lambda'_2 | V^j_S | \lambda_2 \lambda_1 \rangle \]
\[ \langle \lambda'_1 \lambda'_2 | V^j_W | \lambda_1 \lambda_2 \rangle = +\langle \lambda'_1 \lambda'_2 | V^j_W | \lambda_2 \lambda_1 \rangle \]

Under parity, \( P \rangle j m_j \lambda_1 \lambda_2 \rangle = (-1)^{j-1} \langle j m_j -\lambda_1 -\lambda_2 \rangle \) and

\[ \langle \lambda'_1 \lambda'_2 | V^j_S | \lambda_1 \lambda_2 \rangle = +(-\lambda'_1 -\lambda'_2 | V^j_S | -\lambda_1 -\lambda_2 \rangle \]
\[ \langle \lambda'_1 \lambda'_2 | V^j_W | \lambda_1 \lambda_2 \rangle = -(-\lambda'_1 -\lambda'_2 | V^j_W | -\lambda_1 -\lambda_2 \rangle \]

With these constraints, the number of independent matrix elements is reduced to six for \( V_S \) and four for \( V_W \). We define a set of parity and exchange-symmetry eigenstates, with notation \( \langle j m_j \text{ sign}(\lambda_1 \lambda_2) \rangle, (-1)^{j-1}P \rangle \),

\[ \langle j m_j \text{ + +} | = \frac{1}{\sqrt{2}} \langle j m_j \lambda_1 = +\frac{1}{2} \lambda_2 = +\frac{1}{2} \rangle + \frac{1}{\sqrt{2}} \langle j m_j \lambda_1 = -\frac{1}{2} \lambda_2 = -\frac{1}{2} \rangle \]
\[ \langle j m_j \text{ + -} | = \frac{1}{\sqrt{2}} \langle j m_j \lambda_1 = +\frac{1}{2} \lambda_2 = -\frac{1}{2} \rangle - \frac{1}{\sqrt{2}} \langle j m_j \lambda_1 = -\frac{1}{2} \lambda_2 = -\frac{1}{2} \rangle \] (13)
\[ \langle j m_j \text{ - +} | = \frac{1}{\sqrt{2}} \langle j m_j \lambda_1 = +\frac{1}{2} \lambda_2 = -\frac{1}{2} \rangle + \frac{1}{\sqrt{2}} \langle j m_j \lambda_1 = -\frac{1}{2} \lambda_2 = +\frac{1}{2} \rangle \]
\[ \langle j m_j \text{ - -} | = \frac{1}{\sqrt{2}} \langle j m_j \lambda_1 = +\frac{1}{2} \lambda_2 = -\frac{1}{2} \rangle - \frac{1}{\sqrt{2}} \langle j m_j \lambda_1 = -\frac{1}{2} \lambda_2 = +\frac{1}{2} \rangle \]
where $\langle j m_j \, - , + |$ and $\langle j m_j \, - , - |$ are defined only for $j > 0$. The triplet states $\langle j m_j \, + , + |$ and $\langle j m_j \, - , + |$ are coupled by the strong tensor force. In the representation of symmetrized parity eigenstates, $V^j_S$ and $V^j_W$ have the minimum number of non-zero matrix elements. For proton-proton states, the exclusion principle eliminates even-$j$ triplet states $\langle j m_j \, - , - |$, odd-$j$ singlet states $\langle j m_j \, + , - |$, and odd-$j$ triplet states $\langle j m_j \, + , + |$ and $\langle j m_j \, - , + |$. As a consequence, there are no odd-$j$ weak matrix elements. The transform to the conventional $j(\ell s)$ representation is given by,

\[
\begin{align*}
\langle j m_j \, s=0 \, l=j | & = + \langle j m_j \, + , - | \\
\langle j m_j \, s=1 \, l=j | & = - \langle j m_j \, - , + | \\
\langle j m_j \, s=1 \, l=j-1 | & = + \sqrt{\frac{j}{2j+1}} \langle j m_j \, + , + | + \sqrt{\frac{j+1}{2j+1}} \langle j m_j \, - , + | \\
\langle j m_j \, s=1 \, l=j+1 | & = - \sqrt{\frac{j+1}{2j+1}} \langle j m_j \, + , + | + \sqrt{\frac{j}{2j+1}} \langle j m_j \, - , + |
\end{align*}
\]

NUMERICAL RESULTS

The total parity-violating analyzing power, $A$, is defined by,

\[
A = \frac{\sigma_+ - \sigma_-}{\sigma_+ + \sigma_-}.
\]

For simplicity, we neglect effects of Coulomb forces. (These are discussed in Reference 1.) Here, $\sigma_+$ and $\sigma_-$ are the total scattering cross sections for scattering positive and negative helicity protons from an unpolarized target. Results of calculations with the consistent meson models described above are given in Reference 2.

All parity-violating observables for the two-proton system, including $A$, can be calculated from the s- or t-matrix elements determined, with Equation 11, by $r_W$ and $t_S$. The observables depend linearly on $r_W$. All other factors depend only on matrix elements of $t_S$, which are the strong scattering amplitudes given by Arndt's empirical analysis. Since $t_S$ is energy-dependent, these factors modify the energy-dependence of $r_W$. Theoretical predictions for $r_W$ are given in Reference 2.

These allow a more direct comparison of different models.

There is only one independent parity-violating transition with $j = 0$ and two for each even $j > 2$. The $j = 0$ transition gives the dominant contribution to $A$ at the low energies, and at energies near 800 MeV. Near 230 MeV, the empirical energy-dependent factor for the $j = 0$ contribution to $A$ disappears and the $j = 2$ terms dominate. At 800 MeV, the $j = 2$ scattering amplitudes are small and (with $r_W$ real) the $j = 2$ contributions to $A$ are then small. Elements of $r_W$ with $j > 2$ contribute very little to the parity-violating analyzing power below 1200 MeV. Terms with $j > 2$ account for only a small increase in the magnitude of $A$.

All results scale linearly with the parity-violating parameters $h^{pp}$ and $h^{pp}$. These have large theoretical uncertainties; for definiteness, we use the DDH "best values". These parameter values can be adjusted to fit the 45 MeV measurement, but the ratio of $h^{pp}$ to $h^{pp}$ remains ambiguous. Another precise measurement is needed to constrain both parameters.
The separate contributions of rho and omega parity-violating exchanges are shown for the nonrelativistic model in Reference 1, Figure 8.

Qualitative differences in the energy dependence of the results with and without strong distortions are appreciable, especially for the \( j = 0 \) parity-violating matrix element\(^2\). This matrix element depends on the singlet s-wave, which is grossly distorted by the strong nucleon-nucleon potential.

Differences between the static and nonrelativistic calculations are due to relativistic corrections. These include (1) terms omitted by truncating the momentum expansion of \( V_S \) to obtain a nonrelativistic expression and (2) adjustments in the Bonn model parameters and form factors to refit NN data. The second step is, effectively, a relativistic correction, since the adjustments compensate for terms removed in the nonrelativistic reduction. Changes in the Bonn rho and omega parameters and form factors affect \( V_W \) as well as \( V_S \). The overall increase in \( A \) shown for the fully consistent static model is \(~1\) per cent below 45 MeV, 18 per cent at 300 MeV, 44 per cent at 800 MeV, and 54 per cent at 1200 MeV. These corrections are large considering that \( E/m_N \) increases to only 1.08 at 300 MeV and 1.19 at 800 MeV.

The static and nonrelativistic models use the same Lippmann-Schwinger propagator, so differences in \( r_W \) or \( A \) are induced only by changes in \( V_S \) and \( V_W \). To study the effects of \( V_S \) separately from \( V_W \), consider a calculation, similar to the nonrelativistic calculation, but where \( V_S \) is replaced by the static model (OBEPQ) potential. \( V_W \) remains as in the nonrelativistic model. Then, \( A \) increases by 9 per cent below 100 MeV, 15 per cent at 300 MeV, 11 per cent at 800 MeV, and 7 per cent at 1200 MeV. These numbers are quite different those quoted above (1, 18, 44, and 54 per cent respectively), where both \( V_S \) and \( V_W \) are replaced by the corresponding static model potentials. The differences between the numbers (9, 15, 11, 7) and (1, 18, 44, 54) are the combined effect of including relativistic factors and adjusting the Bonn model parameters and form factors in \( V_W \).

The relativistic factors in \( V_W \) are less important than the adjustments in Bonn model parameters. The relativistic factors, \( E_{q'} \) and \( E_q \) in the static model potential \( V_W \) become \( m_N \) in a nonrelativistic reduction of \( V_W \). In the Born calculation, where \( V_S = 0 \), \( V_W \) is evaluated on shell, that is, with \( E_{q'} = E_q = E \). In that case, the formulae for \( V_W \) reduce to simpler expressions where only the transitions \( \langle q', - | V_W^2 | q, + \rangle \) (and adjoints) include \( E/m_N \) corrections. The other Born amplitudes, which include the dominant \( j = 0 \) matrix element, have no corrections when \( E_{q'} = E_q = E \). In the end, the analyzing power \( A \) shows only small increases of \(~0\) per cent below 100 MeV, 1 per cent at 300 MeV, 2 per cent at 800 MeV, and 5 per cent at 1200 MeV due to \( E/m_N \) corrections. In the static model, where off-shell elements of \( V_W \) do contribute, the effects of \( E_{q'}/m_N \) and \( E_q/m_N \) in \( V_W \) are different: \( A \) increases by 1.4 per cent at 100 MeV, 5 per cent at 300 MeV, 4 per cent at 800 MeV, and 2 per cent at 1200 MeV. These increases are still small.

Results for the energy-dependent model are included in Reference 2. Up to 300 MeV, parity-violating effects are consistently larger than those for the static model. At 300 MeV, \( A \) is 14 per cent larger than the static result and 27 per cent larger than the nonrelativistic result. Differences between the energy-dependent and static models are attributable to the static approximation. Notice that the energy-dependent, static, and nonrelativistic models all give nearly identical results below 15 MeV. This does not occur when \( V_S \) and \( V_W \) use inconsistent strong model parameters.
Above 290 MeV, pion production can introduce absorptive terms in $V_W$ or in the strong distortion factors of Equations 4 or 5. $R_W$ and $r_W$ then become complex. The delta resonance may provide the dominant mechanism for pion emission in parity-violating graphs, as it does for the parity-conserving graphs.

In our treatment, $G$, in Equations 1 or 2 is the nucleon-nucleon propagator. The two-pion exchange box diagram with a delta-isobar intermediate state is irreducible, and the boxed and crossed graphs with a delta isobar are complementary terms of the irreducible kernel $V_W$. Many isospin- and spin-dependent terms in the graphs with parity-violating meson-nucleon vertices cancel at low energy, suggesting that both boxed and crossed graphs must be included to predict the correct energy dependence of parity-violating observables. Above threshold, one of the exchanged pions can become real, which makes an inelastic contribution to $V_W$. To obtain the correct energy-dependence, the boxed and crossed graphs should be calculated together.

Delta-isobar pion-emission transitions also contribute to $T_S$ and $R_S$ in Equations 4 and 5. At 800 MeV, the $j = 2$ strong scattering amplitudes contribute large inelastic components to the total scattering cross section. To estimate the effect of pion emission in $T_S$ on the parity-violating $r_W$-matrix amplitudes, we employ an optical potential $V_S$ to calculate the 1D2-3P2 transition. We do not consider the $j = 0$ parity-violating transition, since inelastic effects are small in the 1S0 and 3P0 channels. (The values of Arndt’s inelastic parameters for these channels are smaller than for 1D2 and 3P2). A reasonable optical potential is obtained by setting (when the potential $V_S < 0$),

$$\text{Im } V_{S;1D2-1D2}(r) = \alpha_{1D2} \text{ Re } V_{S;1D2-1D2}(r)$$
$$\text{Im } V_{S;3P2-3P2}(r) = \alpha_{3P2} \text{ Re } V_{S;3P2-3P2}(r)$$

The imaginary parts remain zero when the real parts are positive (repulsive). For the real parts, we use the nonrelativistic Bonn (OBEPR) 1D2 and 3P2 channel potentials. For simplicity, $V_{S;3P2-3F2}(r)$ and $V_{S;3F2-3F2}(r)$ remain real.

Solutions of Schrödinger's Equation are obtained with the optical potential, which depends on $\alpha_{1D2}$ and $\alpha_{3P2}$. By adjusting these to fit Arndt’s inelastic parameters, we obtain $\alpha_{1D2} = .9$ and $\alpha_{3P2} = .24$. With these values, the DWBA calculation gives

$$\text{Re } r_{W;1D2-3P2} = 3.01 \times 10^{-8} \quad \text{Im } r_{W;1D2-3P2} = -.15 \times 10^{-8}$$

For comparison, the result with no imaginary potential ($\alpha_{1D2}, \alpha_{3P2} = 0.0$) is

$$\text{Re } r_{W;1D2-3P2} = 2.95 \times 10^{-8} \quad \text{Im } r_{W;1D2-3P2} = .00 \times 10^{-8}$$

The analyzing power $A$ is roughly four times more sensitive to the imaginary part than to the real part of $r_W$ for the 1D2-3P2 transition at 800 MeV. (We assume Arndt’s values for the 1D2 and 3P2 phases). Consequently, the imaginary part of $r_W$ has a twenty per cent effect on the 1D2-3P2 contribution to $A$. Compared with the dominant 1S0-3P0 contribution, however, this effect is small.
REFERENCES

6. R. Machleidt, K. Holinde, and Ch. Elster, Phys Rep 149, No. 1 (1987). The parameter values for OBEPT are given in Table 8, p. 53 of this reference. The r-space potential, OBEPR, is defined in Appendix F.
LOW-ENERGY WEAK \( pp \) ASYMMETRY: ENERGY DEPENDENCE

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ABSTRACT

The weak interactions giving elastic \( pp \) weak asymmetry \( (A^W) \) are short-ranged, so theoretical calculations can be done in quark models, where the weak quark interactions from particle physics can be used directly. We use two quark models: the HQHM, which has been used extensively for studies of electromagnetic form factors and other nuclear properties, and a generator coordinate type quark model, which is currently being developed for the two-nucleon system. It is found that the energy-dependence of the weak asymmetry is model dependent, in contrast to the conventional meson-exchange model of weak nuclear interactions—in which the energy dependence at low energy is given as \( \sqrt{E} \). In the version of the HQHM that has been most extensively used the magnitude of the weak asymmetry is found to increase in going from 45 MeV (where an accurate measurement has been made) to 15 MeV, rather than decreasing as \( \sqrt{E} \). A major new element of the quark models is that the weak \( P-D \) transition is important at lower energies than expected in the local meson-exchange model. Nonlocalities are the source of these differences. Accurate measurements of \( A^W \) at about 15 MeV could give valuable information about the nature of weak nuclear interactions and possibly nuclear structure.

INTRODUCTION

Today I will discuss the elastic \( pp \) weak asymmetry, \( A^W(E, \theta) \), from the point of view of quark models of nucleons and nuclei. Since the weak one-pion exchange potential cannot give a weak elastic \( pp \) asymmetry, the weak interactions must take place at short distance, or possibly at midrange. The physical picture for short-range interactions is quite different in quark models than in conventional meson-exchange models, as illustrated in Fig. 1.

![Fig. 1](image-url)

Although models of \( \rho \) and \( \omega \) exchange between point nucleons can provide all of the operators expected to be effective for low-energy parity violating effects, it is not obvious that the nonlocal distribution of the weak quark interactions [Fig. 1(a)] can be accurately represented by averaging over the positions of the quarks to give the conventional local meson-exchange model [Fig. 1(b)].
A recent study\textsuperscript{1} showed that the energy dependence of the total weak asymmetry in a quark model can be quite different from the meson-exchange models currently used at low energy, where the energy dependence of the meson-exchange model is known to be $\sqrt{E}$, as in any static local model. In the present work we explore the model-dependence of $A_{\text{tot}}^W(E)$. We also examine the angular dependence, particularly with regard to the importance of the $P\cdot D$ weak parity-violating amplitude, which is not important in meson-exchange models at low energy, but is seen to be very important for quark models, because of the large nonlocality.

A measurement of the total weak asymmetry at 45 MeV has been carried out recently\textsuperscript{2} to an accuracy of $2 \times 10^{-8}$. One main conclusion of the present work is that an experimental determination of $A_{\text{tot}}^W$ at a lower energy would be very worthwhile. We were pleased to learn at this workshop that such an experiment at 15 MeV is in progress at Bonn.\textsuperscript{3}

THE WEAK QUARK EFFECTIVE HAMILTONIAN

The weak parity-violating asymmetry is defined as

$$A^W(E, \theta) = \frac{(\sigma(E, \theta) - \sigma(E, \theta))}{(\sigma(E, \theta) + \sigma(E, \theta))},$$

(1)

where $\sigma$ is the cross section for the scattering of a longitudinally polarized proton by a target proton. The theoretical expression for $A^W$ is

$$A^W(E, \theta) = \text{Re} \left[ \sum_{s=\text{spins}} F_{s}^{st}(E, \theta) f_{s}^{W}(E, \theta) / \sum_{s} |F_{s}^{st}(E, \theta)|^2 \right],$$

(2)

where the strong + electromagnetic interaction amplitude, $F^{st}$, can be obtained from experimental $NN$ phase shifts. The theoretical problem is to calculate the weak amplitudes,

$$f_{s}^{W}(E, \theta) \propto \langle pp | H^W | pp \rangle_s.$$  

(3)

To do this one needs a model for the weak Hamiltonian, $H^W$, and for the $pp$ states.

An effective weak quark Hamiltonian is needed for the quark model calculations, since there are important QCD corrections to the lowest-order weak quark interactions (given by $W$ and $Z$ exchange between quarks, as illustrated in Fig. 2). The method used to obtain the effective weak quark interaction is to calculate the one-gluon loop corrections to one-$W$ or -$Z$ exchange processes and the “penguin” diagrams at high $Q^2$.\textsuperscript{4} The form of the resulting parity-violating Hamiltonian in the $pp$ system is\textsuperscript{1}

$$H_{pp}^{p.v.(\text{eff})} = K(\mu^2) \frac{G}{\sqrt{2}} \bar{q} \gamma^\mu q \bar{q} \gamma^5 q + \text{c.c.}$$

(4)

which is just the standard Cabbibo model for $\Delta S = 0$ multiplied by the numerical function $K(\mu^2)$, which is given in Ref. 1. The quantity $\mu^2$ is the “renormalization point,” corresponding to typical momenta in the system. Since in our calculations the wave functions are determined...
by the treatment of the strong interaction, the value of $\mu^2$ is the main parameter in our formulation. In most cases $\mu^2$ is chosen to fit the 45-MeV experimental value for $A^W$. Since the corresponding value of $\mu^2$ is less than $1.0 \text{ (GeV)}^2$, the one-loop corrections are not expected to be reliable. Therefore, one must consider our procedure to be a prescription for including nonperturbative QCD corrections rather than a calculation using perturbative QCD.

QUARK MODEL CALCULATIONS

Two quark models of protons and the $NN$ system are used for calculating $A^W$. The most extensive and complete studies have been done in the Hybrid Quark Hadron Model (HQHM). In this model conventional hadronic models are used to describe systems of baryons when the centers of the baryons are separated by a distance greater than an assumed transition length, $r_0$. When the centers are separated by a distance less than $r_0$, we assume that a six-quark system has formed with the quarks all in color contact. The size of the six-quark clusters is a parameter, $R_6$. The main assumption of the model is that at the transition distance $r_0$ the confinement of quarks into three-quark baryon clusters disappears, and that partially deconfined six-quark clusters form.

The HQHM has been used for extensive studies of exclusive form factors of few-nucleon systems. The main parameters of the HQHM for the elastic scattering of electrons from $^2\text{H}$, $^3\text{He}$, and $^3\text{H}$ targets, and for threshold electrodisintegration of the deuteron, are the two parameters, $r_0$ and $R_6$. The results of the fits, which are discussed below are:

$$r_0 = \text{transition length} \approx 1.0 \text{ fm}$$
$$R_6 = 6q \text{ cluster size} \approx 3/2 R_3,$$

where $R_3$ is the size of a three-quark cluster (a nucleon).

In the study of elastic electron scattering with a deuteron target$^6$ in the HQHM, the magnetic form factor, $B(q^2)$, is dominated by the six-quark cluster term at rather low values of $q^2$, as shown in Fig. 3(a). The model works well up to about $q = 1.0 \text{ GeV}$, as shown in Fig. 3(b). Since the $B$ form factor is very sensitive to the six-quark clusters, the parameters
are well-determined in this calculation, with the results shown in Eq. (5). With these same values for the parameters, the threshold deuteron electrodisintegration is also fit well up to 1 GeV, as shown in Fig. 4.

Studies of the elastic form factors of $^3$He and $^3$H have also been carried out in this model. It was found that the charge form factors could be fit well with the same parameters [Eq. (5)]. Recently calculations with the model have been carried out for the magnetic form factors of $^3$He and $^3$H with the same parameters. The results are shown in Fig. 5. Once more, one obtains very good fits to the experimental data up to about 1.0 GeV.

Overall, the HQHM provides excellent fits to exclusive form factors of few-body systems for momentum transfers up to $q^2$ about 1.0 (GeV)$^2$. However, this should be interpreted as showing how little we know about the short-range two-body currents rather than a detailed test of a particular quark cluster model. The fits in the HQHM up to 1.0 GeV depend mainly on the parameter $r_0$, which determines the probability of six-quark clusters, and $R_0$, which introduces a new length scale into nuclear physics. The effects are illustrated in Fig. 6. The impulse approximation form factor for a $A$-mass nucleus [$F_A(q^2)$ in Fig. 6] falls much faster with $q^2$ than does the smaller
nucleon’s form factor $[F_N(q^2)]$. The form factor associated with the six-quark cluster, $[F_{6q}(q^2)]$, has a $q^2$ dependence in between the two, reflecting the intermediate length scale given by $R_6$. Meson exchange currents also contain this intermediate length scale. This illustrates how little we learn about meson-exchange currents or quark-cluster effects up to momentum transfers of 1.0 GeV.

It should be emphasized that the six-quark cluster probabilities, $P_{6q}(E)$, are not parameters in the HQHM. They are determined by probability current conservation and continuity. The method used to find these probabilities for scattering states is to start with the experimentally determined $NN$ phase shifts in each SLJI channel, use a potential that can approximately produce the phase shifts at that energy, and integrate the wave function determined from the potential to the matching point $r_0$.

An essential point for the understanding of the results on the weak asymmetry is that the changes in the interior quark configuration which are discussed below with regard to $A^w$ have little effect on the exclusive form factors just discussed, since the latter depend on the $P_{6q}$ and not on the detailed configurations within a particular channel, to a good approximation.

Let us now return to the study of $A^w$. In the conventional static local meson exchange model the weak parity-violating amplitude at low energy is independent of energy. It follows from Eq. (2) that the low-energy energy dependence of $A^w$ is prescribed:

$$A_{tot}^{W(me)}(E) \propto f^{W(me)} \times F^{(st)}(E) \propto \sqrt{E}.$$  
(6)

This result is valid for any local model. Therefore, in meson exchange models, the value of $A^w$ at 15 MeV is determined by the value at 45 MeV.

In the HQHM the weak parity-violating amplitude has an energy dependence. Although the interaction, given in Eq. (4), is energy independent, the weak matrix element of Eq. (3) has an energy dependence that arises from the energy dependence of $P_{6q}(E)$, the probability of the interior six-quark cluster in each channel. Therefore, the energy dependence of the parity-violating asymmetry in the HQHM is given (schematically) by

$$A_{tot}^{W(HQHM)}(E) \propto \sqrt{E} \times P_{6q}(E).$$  
(7)

The results for $A^w(E)$ at energies up to 100 MeV are shown in Fig. 7. The experimental point shown at 45 MeV is the recent accurate measurement at SIN, $^2 A^w = -1.5 \pm 0.2 \times 10^{-7}$. At about 15 MeV, there is a LANL measurement, $^{11} A^w = -1.7 \pm 0.8 \pm 10^{-7}$; and at this workshop, a Bonn measurement was reported $^3$ as $A^w = -3.0 \pm 1.6 \times 10^{-7}$. Any meson exchange model would predict that $A^w$ is less in magnitude than $1.0 \times 10^{-7}$ at 15 MeV, while the HQHM predicts that $A^w$ is about $-1.7 \times 10^{-7}$. The errors in the LANL and Bonn measurements are too large to test this important difference at the present time.

![Fig. 7](image-url)
The difference between the meson exchange and the HQHM can only be nonlocality. At such low energies the spin-angle operators given by $\rho$ or $\omega$ exchange must be similar to the quark operators resulting from the Hamiltonian of Eq. (4). To test the sensitivity of the results to details of HQHM and to try to understand the physical difference between the nonlocal quark picture and the local meson-exchange picture, a series of calculations was carried out changing the $pp$ deformation while keeping all of the $P_{6q}(E)$ constant. Let us explain what is meant by the deformation. The gluonic interactions have tensor forces, as all vector exchange interactions must. All quark model calculations of the deuteron or the $S$-states of the $NN$ system based on QCD-motivated $qq$ interactions show that these states are deformed. I.e., in addition to the $6s_{1/2}$ configurations, there are configurations containing $\ell \neq 0$ quark orbits. In all of the previous calculations in the HQHM, the deformation was fixed, and we believe that this is correct.

The results obtained by changing the deformation but keeping the six-quark probabilities constant are shown in Fig. 8. An important observation is that the spherical case resembles the meson-exchange model. The most significant difference between the deformed and spherical versions of the HQHM for $A^W$ is that in the spherical version no weak $P-D$ amplitude can be present, and the weak $S-P$ transition completely dominates. In the deformed version, however, there is an important weak $P-D$ component. This arises solely from the nonlocality, and is discussed further below.

To study these ideas further, we have used a continuous quark model to calculate $f^W$. This is a generator coordinate method of the general form that has been used by a number of theorists (see, e.g., Oka and Yazaki for a review and references). It is a quark model for nucleons and $NN$ systems with a Hamiltonian of the form

$$H = H_{\text{conf}} + H_{\text{gluon}} + H^{(\pi,\sigma)}.$$  (8)

The parameters of the confinement potential, $H_{\text{conf}}$, which is of the Dirac harmonic oscillator form, are determined by fitting the proton mass and radius and the nucleon magnetic dipole moments. The long- and medium-range $NN$ interactions are given by $H^{(\pi,\sigma)}$, with coupling to quarks. The gluon-quark exchange interactions give short-range effects which are important as nucleons overlap. The size of the six-quark cluster as observed in the HQHM is represented by polarization as nucleons overlap. Typical fits to phase shifts are shown in Fig. 9. In general, we have not been able to obtain good fits to $P$-wave phase shifts, so our results for the calculation of $A^W$ must be considered preliminary.

The calculation of $A^W(E,\theta)$ in this model is tedious but straightforward. Having determined the wave functions via fits to nucleon and $NN$ observables, one calculates $f^W(E,\theta)$ as in Eq. (3) using $H^{P}\mathcal{P}$ of Eq. (4). Preliminary results for $A^W_{\text{tot}}(E)$ are shown (labeled RGM).
in Fig. 10. The most important observation is that the $P-D$ transition is very important at low energy. This is seen to be quite dramatic in the present calculation, with qualitatively different results when the $P-D$ contribution is omitted. This is also illustrated in Fig. 11, where the angular dependence of $A^w(\theta)$ is shown. One can see the important effect of the $P-D$ transition as one goes from the spherical to deformed HQHM. The angular dependence of the RGM can be understood by adding some additional $P-D$ transition. Note that the normalization of the curve "omit P-D" is arbitrary.
CONCLUSION

In conclusion, in quark models the energy dependence of $A^W(E)$ is model-dependent rather than being given as $\sqrt{E}$ at low energy. A main new feature in comparison with meson exchange models is that the $D$-state enters at rather low energy. This can be understood by realizing that although the centers of nucleons are well-separated due to the centripetal potential in the $D$-state, the quarks can still interact. This is illustrated in Fig. 12. This peripheral interaction is magnified by the spin-flip aspects of $A^W$. I believe that it should be possible to represent this in static models by using a weak parity-violating potential of the form

$$V^{p.v.} = [a(E) + b(E) \ell \cdot \ell] V^{p.v.}(r), \quad (9)$$

where $V^{p.v.}(r)$ is a conventional meson exchange weak local potential and $\ell$ is the relative $NN$ orbital angular momentum. It would be interesting to investigate other phenomena to see if such $\ell \cdot \ell$ potentials appear. Finally, an accurate measurement of $A^W$ at 15 MeV would give valuable information about the nature of weak quark interactions.

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AN INTRODUCTION TO CP VIOLATION

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ABSTRACT

A review of the status of CP violation in kaons is given. An up to date knowledge of quark mixing angles in the standard six quark model is presented. The role $B_d^0 - B_d^0$ transition plays in this study is examined. A comparison of the estimates of CP violation effects from models beyond the standard one is given. Other experiments that have the capability of testing different CP violation models are also discussed.

I. INTRODUCTION

This year is the twenty-fifth anniversary of the discovery of CP violation\(^1\) effects in kaon decays. The field has grown from an esoteric area of physics into one that bridges particle and nuclear physics as well as cosmology. There were major discoveries made during these years. Just last year we have established $B_d^0 - B_d^0$ mixing.\(^2\) Although it is not a CP violating phenomenon per se, it has important implications in our study of CP violation. Also, direct CP violation was first reported.\(^3\)\(^\text{a}\) However, this is not being confirmed by the FNAL experiment.\(^3\)\(^\text{b}\) On the theoretical side a framework of describing CP violation in the context of modern spontaneously broken gauge theories of electroweak interactions is now firmly in place. This is the standard model Kobayashi-Maskawa (KM) description of quark mixing and CP violation.\(^4\) The next step is clearly to determine how well the KM paradigm describes quark mixings and CP violation. At the same time we should also be on the lookout for CP violating effects that are not within KM. It is important to emphasize now that we do not have a deep understanding of the KM matrix. The KM description is a convenient and successful parametrization of data. Hence, the origin of CP violation still alludes us even if we have established that the KM matrix is all there is to CP violation. It is certainly our hope that this is not the case.

The importance of CP violation in nuclear and hadronic physics lies in its connection to violation of time reversal invariance (T). Due to the CPT theorem which is central to local quantum field theories, CP violation implies T violation. Currently the best candidate of strong interaction is quantum chromodynamics (QCD). This theory has a non-trivial vacuum structure. The effective Lagrangian of QCD consists of a term $\theta_{QCD} G^{\mu\nu} \tilde{G}_{\mu\nu}$ where $G^{\mu\nu} (\tilde{G}_{\mu\nu})$ is the field tensor (dual tensor) of the gluon gauge fields of QCD. This term violates both $P$ and $T$ and hence can induce a electric dipole moment for the neutron. The non-observation of this quantity sets $\theta_{QCD} \leq 10^{-9}$. It is one of the unresolved problems of the standard model, why is this parameter so small. CP violation also has profound connection with cosmology. The most important one being the observed asymmetry between baryons and antibaryons in our part of the universe. It is observed that there are hardly any antibaryons around. If we denote the average number densities of baryons, antibaryons and photons in the present universe by $n_B$, $n_B$ and $n_\gamma$ respectively, the cosmological baryon asymmetry (CBA) is

$$\frac{n_B - n_B}{n_\gamma} \sim \frac{n_B}{n_\gamma} \sim (4 - 7) \times 10^{-10}. \quad (1)$$

This means that during the evolution of the universe, the baryons and antibaryons did not totally annihilate leaving a small amount of baryon matter left. The explanation of this
asymmetry almost certainly requires CP violation. This is presently a very active research area in cosmology and particle physics.

In this lecture I will discuss the present status of quark mixing and CP violation. Sources of CP violation beyond the standard model will also be described. Since these models are mostly speculations this section carries the subtitle "Lampposts in fairyland". Next I will describe some current discussions of CP violation experiments beyond the usual $K \to \pi\pi$ mode. I view this as an important dialogue between experimentalists and theorists out of which new ideas and results will emerge. These are important first steps. Regrettably, I have to omit the interesting question of CP violation in the $b$-quark system. This is simply due to a lack of time. The presentation is aimed at the non-experts such as graduate students, post-doctoral fellows and active physicists in other areas of sub-atomic physics, who are interested in the profound problem of CP violation. The literature on CP violation is extensive. It is impossible to do justice to all the authors who made important contributions in this area. More detailed discussions of the topics covered here can be found in Ref. 6.

II. QUARK MIXINGS AND THE K-M MATRIX

In the standard model the quarks come in three families listed below

$$
\begin{pmatrix}
  u \\
  d \\
  c \\
  s \\
  t \\
  b
\end{pmatrix}
= \begin{pmatrix}
  u_R, d_R \\
  c_R, s_R \\
  t_R, b_R
\end{pmatrix}.
$$

(2.1)

The left-handed quarks come in a doublet representation of the SU(2) part of the SU(2)$_L \times$ U(1) gauge interactions whereas the right-handed quarks are SU(2) singlets. These are eigenstates of the weak interactions. In general they are not the mass eigenstates. However, they are related to each other by a unitary transformation. In the minimum case of 3 quark families, the unitary transformation is represented by a $3 \times 3$ unitary matrix. Hence, we have

$$
\begin{pmatrix}
  d' \\
  s' \\
  b'
\end{pmatrix}
= \begin{pmatrix}
  V_{ud} & V_{us} & V_{ub} \\
  V_{cd} & V_{cs} & V_{cb} \\
  V_{td} & V_{ts} & V_{tb}
\end{pmatrix}
\begin{pmatrix}
  d \\
  s \\
  b
\end{pmatrix},
$$

(2.2)

where the prime denotes weak eigenstates and the unprimed ones are mass eigenstates. The Lagrangian of the charged weak interactions is then given by

$$
L = -\frac{g}{\sqrt{2}} \sum_{i=u,c,t, j=d,s,b} \bar{q}_i \gamma^\mu \left(1 - \gamma^5\right) \frac{1}{2} V_{ij} q_j + h.c.
$$

(2.3)

where $g$ is SU(2) gauge coupling. A $3 \times 3$ unitary matrix can be written in a form resembling the Euler rotations of rigid bodies, i.e. in terms of 3 mixing angles plus a physical phase $\delta$. There are many forms of writing this and a commonly used one is given below

$$
V_{ij} = \begin{pmatrix}
  c_{12} c_{13} & s_{12} c_{13} & s_{13} e^{-i\delta} \\
  -s_{12} c_{23} - s_{23} s_{13} c_{12} e^{i\delta} & c_{12} c_{23} + \ldots & c_{13} s_{23} \\
  s_{12} s_{23} - s_{23} s_{13} c_{12} e^{i\delta} & -s_{23} c_{12} + \ldots & c_{13} c_{23}
\end{pmatrix},
$$

(2.4)
where \( c_{ij} (s_{ij}) \) denotes the \( \cos (\sin) \) of the mixing angle between \( i \) and \( j^{th} \) families of quarks. The angles are such that \( c_{13} \approx c_{23} \approx 1 \) and the dots in Eq. (2.4) are terms much smaller than the ones we retained.

We have a fair bit of knowledge of the elements of \( V_{ij} \). The first element is very well measured by the \( O^+ \to O^+ \beta \)-transitions. Currently, we have\(^7\)

\[
V_{ud} = 0.9740 \pm 0.00010 ,
\]

which is the culmination of many years of careful experimentation and calculations. Ambiguities in dealing with the necessary radiative corrections of nuclear matrix elements prevent us from using this route for an improvement on the measurement of \( V_{ud} \). Pion \( \beta \)-decay and free neutron \( \beta \) decay would be ways out of this difficulty. However, both of these experiments present formidable challenges if one wishes to achieve the desired accuracy.

The element \( |V_{us}| \) is also very well determined from \( K_{e3} \) and hyperon decays. The world average is\(^8\)

\[
|V_{us}| = 0.2196 \pm 0.0014 \pm 0.0018 .
\]

A breakdown of the errors in the analysis \( K_{e3} \) reveals:

\[
|V_{us}| = 0.2196 \pm 0.0014 \pm 0.0018 .
\]

The last uncertainty is theoretical and one should aim at reducing it. The first uncertainty comes from radiative effects and experimental uncertainties. It appears to be possible to reduce these uncertainties to \( \pm 0.0005 \) which will be a job for the kaon factories.

We come to the last element in the first row of the KM matrix; namely \( V_{ub} \). Unitarity dictates that

\[
|V_{ud}|^2 + |V_{us}|^2 + |V_{ub}|^2 = 1 .
\]

whereas experimentally we have \( 0.9970 \pm 0.021 \) from \( V_{ud} \) and \( V_{us} \) alone. We can use this to set an upper limit on \( |V_{ub}| \). An important point to note is that without taking into account radiative corrections, the left-hand side of Eq. (2.8) would be 1.035. This is clearly in violation of unitarity. We will have more to say about \( V_{ub} \) later.

We quickly move on to the second row. The elements \( V_{cd} \) and \( V_{cs} \) are fairly well determined from neutrino production of charmed particles. From these measurements we obtain

\[
|V_{cd}| = 0.22 \pm 0.03
\]

and

\[
|V_{cs}| = 0.95 \pm 0.14 .
\]

This brings us to \( |V_{bc}| \). This is measured by the semileptonic branching ratio of \( b \)-mesons into charmed ones plus \( \nu_e \) pair at the quark level depicted in Fig. 1. In the limit that one can neglect the masses of the final state quarks, the partial width of \( b \to c\ell\nu \) is given by the usual formula for muon lifetime with the appropriate substitution of \( m_\mu \to m_b \) and including the mixing element \( V_{bc} \). Explicitly we have

\[
\Gamma(b \to c\ell\bar{\nu}) = \frac{G_F^2 m_b^5}{192 \pi^3} |V_{bc}|^2 .
\]

Fancier calculations exist that include the correct phase-space factors to account for \( c \)-quark mass and QCD corrections. From this we can extract

\[
|V_{bc}| = 0.046 \pm 0.010 .
\]
Now we can return to $|V_{ub}|$. The ratio $|V_{ub}/V_{bc}|$ can be obtained from the semileptonic decay of $B$-mesons by fitting the charged lepton spectrum as a sum of two components containing $b \to c\ell\nu$ and $b \to u\ell\nu$. Since the $u$-quark is almost massless whereas the $c$-quark is about 1.5 GeV one expects more high energy charge leptons if $|V_{bu}|$ is large. Preliminary reports on these high energy leptons gives

$$\frac{|V_{bu}|}{|V_{bc}|} \simeq 0.09 \pm 0.02.$$ \hfill (2.12)

If this is true, it will imply that charmless $B$ meson decays such as $B^0 \to p\bar{p}\pi^+\pi^-$ modes will have to be observed at the level $10^{-4} - 10^{-5}$.

With the knowledge we gained thus far we can constrain the mixing elements of the third row of $V_{ij}$. This is shown in (2.13)

$$|V_{ij}| = \begin{pmatrix} 0.974 & 0.22 & 0.003 - 0.01 \\ 0.22 & 0.974 & 0.05 \\ 0.001 - 0.023 & 0.05 & 0.999 \end{pmatrix}. \hfill (2.13)$$

The elements $|V_{ij}|$ are known within two and three decimal places without having found the $t$-quark! This certainly is a measure of how successful the standard model is. A closer look at Eq. (2.13) reveals the KM matrix elements exhibit a hierarchical structure. The diagonal elements are close to unity. Going from family 1 to family 2 a factor of $\lambda = 0.22$ is involved. Going from family 2 to family 3 a factor of $\lambda^2$ is involved. This observation leads to the Wolfenstein-Miani parametrization\(^\text{10}\) of KM matrix in powers of $\lambda$ as follows.

$$V \simeq \begin{pmatrix} 1 - \frac{1}{2} \lambda^2 & \lambda & A\rho \lambda^2 e^{i\delta} \\ -\lambda(1 + A^2\lambda^4\rho e^{-i\delta}) & 1 - \frac{1}{2} \lambda^2 & A\lambda^2 \\ A\lambda^3(1 - \rho e^{-i\delta}) & -A\lambda^2(1 + \lambda^2\rho e^{-i\delta}) & 1 \end{pmatrix}. \hfill (2.14)$$

This is an approximate representation in power series of $\lambda$ and we have only kept the leading terms. $A$ is given by $|V_{bc}|$ and is found to be

$$A = 1.05 \pm 0.17.$$ \hfill (2.15)

This is a very convenient form for the study of CP violation. The physical phase $\delta$ always appears with at least a factor of $\lambda^3$. Hence, the observed weakness of CP violation effect is understood to be due to the smallness of the quark mixing angles. Secondly, CP violation is associated with either the $c$, $b$ or $t$ quark. We shall see that CP violation in the $s$-quark system involves both $c$ and $t$-quarks in an essential way. This is due to the fact that there is no direct flavor changing neutral current in the standard model.

![Fig. 1. Decay modes of $b$-quark in the standard model.](image-url)
III. $B_d^0 - \bar{B}_d^0$ MIXING

In the last two years we have seen the discovery and confirmation of the $B_d^0 - \bar{B}_d^0$ mixing of phenomena by ARGUS and CLEO. It is fair to mention that indications of this mixing appeared previously in the UA1 detector although no firm conclusion can be drawn from them. In the ARGUS and CLEO experiment, the $B_d^0$ and $\bar{B}_d^0$ mesons are pair produced at the $\Upsilon(4S)$ state:

$$e^+e^- \rightarrow \Upsilon(4S) \rightarrow B_d^0 \bar{B}_d^0 .$$

Ordinary $B_d^0$ decays into $l^-\bar{\nu}^+ +$ charmed particles semileptonically whereas $\bar{B}_d^0$ into $\bar{l}\nu^+$ charmed particles, see Fig. 2. Hence, we have opposite sign dileptons, in the final states. If mixing occurs, i.e. $B_d^0 \rightarrow \bar{B}_d^0$ or $\bar{B}_d^0 \rightarrow B_d^0$, then we will have a number of $l^+l^+$ or $(l^-l^-)$ events in the final state. The actual number of these same sign dilepton events depends on the strength of this mixing. This is indeed what was found and the ratio of same sign dileptons to opposite dileptons is given by

$$r_d \equiv \frac{N(e^+e^+) + N(e^-e^-)}{N(e^+e^-)} = 0.21 \pm 0.08 \quad \text{ARGUS}$$

$$= 0.15 \pm 0.06 \quad \text{CLEO} .$$

In SM, $B_d^0$ being a $bd$ meson and $\bar{B}_d^0$ a $\bar{b}d$ meson, can make transition into each other by exchanging $W^+W^-$, i.e. a higher order effect. This is due to the fact that there are no direct or tree-level flavor changing neutral currents in SM. This is depicted in Fig. 3. The calculation of the diagram gives the mixing parameter $x_d$ which is related to $r_d$ by

$$r_d \simeq \frac{x_d^2}{2 + x_d} .$$

Explicitly, $x_d$ is

$$x_d = \frac{\tau_B}{\tau_B} \frac{G_F^2}{6\pi^2} \frac{\eta_{QCD}}{\eta_{QCD}} \frac{m_{t}^2 M_{B}}{m_{t}^2 M_{B}} \left[ B_{BF} \right] \left[ f \left( \frac{m_{t}^2}{M_{w}^2} \right) \right] |V_{td}|^2 .$$

The various factors are understood as follows:

i) $\eta_{QCD} \approx 0.85 - 0.63$ is a factor due to gluonic exchanges between quarks. Without including them we would have $\eta_{QCD} = 1$ which is the free quark result.

ii) $\tau_B$, and $M_B$ are simply the lifetime and the mass of the $B$-meson which is experimentally known.
iii) $B_1 f_B^2$ represents the hadronic matrix element and at present is not reliably calculable. $f_B$ is the $B$-meson decay constant similar to the familiar $f_\pi$ and $f_K$ for pions and kaons. The uncertainty in this quantity is about a factor of 2, i.e. $[B_1 f_B^2] \sim 100 - 200 \text{ MeV}$.

iv) The function $f \left( \frac{M_B^2}{M_\pi^2} \right)$ is the result from calculating the box-diagram of Fig. 3 and is given by

$$f(y) = 1 - \frac{3y(1+y)}{4(1-y)^2} \left( 1 + \frac{2y^2}{(1-y^2) \ln y} \right).$$  \hspace{1cm} (3.5)

v) $|V_{td}|^2$ is the KM matrix which can be read from Eq. (2.4) or (2.14). We have also used $|V_{tb}|^2 \simeq 1$.

The interesting point is that the measured value of $x_d$ implies a heavy $t$-quark. Using

$$|V_{td}|^2 = \lambda^2(1 + \rho^2 - 2\rho \cos \delta),$$  \hspace{1cm} (3.6)

and Eq. (3.3) $x_d$ gives a constraint equation on $\rho$, $\delta$ and $m_t$. Equations (3.3) and (3.5) illustrate the point I made in the beginning; namely $B_1^0 - \bar{B}_1^0$ mixing is not a CP violating effect per se, i.e. even if $\delta = 0$, $x_d$ is not vanishing. However, in SM it is closely linked to CP-violation because it relates the parameters $\rho$, $\delta$ and $m_t$ non-trivially.

IV. CP-VIOLATION IN KAONS

We return now to the discovery that started the whole game; namely CP violation in $K \to \pi\pi$ decays. The two most important parameters are $\epsilon$ and $\epsilon'$. The first one measures the CP violating admixture in the wave functions of the neutral kaons $K_L^0$ and $K_S^0$. We need here a few standard notations: the definite CP states are defined by CP $|K^0 > = -|\bar{K}^0 >$, CP $|\bar{K}^0 > = \bar{K}^0 >$, $|K_1 > = K^0 > - 1\bar{K}^0 > \sqrt{2}$ and $|K_2 > = |K_2 > + \sqrt{2}$ with CP $|K_1 >$ to $= |K_1 >$ and CP $|K_2 > = -|\bar{K}^2 >$. The states $K_L^0$ and $K_S$ is then given by

$$K_L = \frac{K_2 + \epsilon K_1}{1 + |\epsilon|^2},$$  \hspace{1cm} (4.1)

![Fig. 3. The box diagram for $B_1^0 - \bar{B}_1^0$ mixing.](image-url)
and

\[ K_S^0 = \frac{K_1 + \epsilon K_2}{1 + |\epsilon|^2}. \]  

(4.2)

Thus for a small \( \epsilon \) \( K_L^0 \) is mostly CP odd and \( K_S^0 \) is mostly CP even. Experimentally one measures the amplitudes \( K_L^0 \to \pi^+\pi^- \) and \( K_S^0 \to \pi^+\pi^- \). The neutral models of \( \pi^0\pi^0 \) final states are also measured. The ratio of the amplitudes gives

\[ \eta_{+-} = \frac{A(K_L \to \pi^+\pi^-)}{A(K_S \to \pi^+\pi^-)} = \epsilon + \epsilon', \]

(4.3)

\[ \eta_{00} = \frac{A(K_L \to \pi^0\pi^0)}{A(K_S \to \pi^0\pi^0)} = \epsilon - 2\epsilon'. \]

(4.4)

Equations (4.3) and (4.4) show that CP violation can arise from the wave function as given by \( \epsilon \) as well as directly in the decay as given by \( \epsilon' \). It is customary to discuss the relative strengths of these in terms of the ratio \( \frac{\epsilon'}{\epsilon} \). The superweak theory\(^{11}\) gives \( \frac{\epsilon'}{\epsilon} = 0 \).

In SM, \( \epsilon \) is given by the imaginary part of the box diagram as depicted in Fig. 4. Standard analysis gives

\[ |\epsilon| = \frac{G_F^2 M_K M_{\pi^0}}{6\sqrt{2}\pi^2 \Delta M_K} B_K f_K^3 A^2 \rho \delta \sin \delta \]

\[ \left[ y_c \{ \eta_3 f_3(y_c, y_t) - \eta_1 \} + \eta_2 y_t f_2(y_t) A^2 \lambda^4 (1 - \rho \cos \delta) \right] \]

+ small term ,

(4.5)

where \( \Delta M_K \) is the \( K_L^0, K_S^0 \) mass difference which is measured. Again \( \eta_1 \simeq 0.4, \eta_2 \simeq 0.6 \) and \( \eta_3 \simeq 0.7 \) are QCD correction factors. The function \( f_3 \) is

\[ f_3(y_c, y_t) = \ln \frac{y_t}{y_c} - \frac{3y_t}{4(1 - y_t)} \left( 1 + \frac{y_t}{1 - y_t} \ln y_t \right), \]

(4.6)

where \( y_{c,t} = \frac{m_{c,t}^2}{M_w^2} \) and \( f_2 \) is given in (3.4). \( \epsilon \) is a genuine CP violating quantity as seen explicitly by the \( \sin \delta \) factor. It is important to note that both the \( c \) and \( t \) quarks are of equal importance for \( m_t \approx M_W \). This is seen by comparing the first and second terms of Eq. (4.5). Although the \( t \) quark contribution is down by \( \lambda^4 \), \( m_t \) being much larger than \( m_c \) compensates for this. The hadronic uncertainty is similar to \( B_s^0 B_d^0 \) mixing. Combining these two calculations a constrain on \( \rho \) and \( \delta \) is obtained. For numerical details see Ref. 12.
Let me reiterate the physics behind the KM description of CP-violation in the kaon system. The observed weakness of the strength of CP-violation is due to the smallness of the quark mixing angles, i.e. involves \( \lambda^6 \), which in turn is due to the essential role played by the heavy \( c \) and \( t \) quarks.

What about \( \frac{\epsilon'}{\epsilon} \)? The experimental situation is given below

\[
\text{Re} \left( \frac{\epsilon'}{\epsilon} \right) = (3.3 \pm 1.1) \times 10^{-3} \text{ NA31} \\
( -0.5 \pm 1.4 \pm 0.6) \times 10^{-3} \text{ FNAL} .
\]  

(4.7)

This is clearly an unsettled situation. In SM, \( \frac{\epsilon'}{\epsilon} \) is given by the interplay between the electroweak and the strong or color interaction via the penguin diagrams. This is depicted in Fig. 5. There are a lot more theoretical uncertainties involved here than in the calculations of \( B_d^0 - \bar{B}_d^0 \) mixing and \( \epsilon \). It is simply impossible to get into details here. Calculation with SM incorporating the range of \( t \)-quark mass allowed by \( B_d^0 - \bar{B}_d^0 \) mixing. Buras et al.\(^{14} \) obtains

\[
\frac{\epsilon'}{\epsilon} = (2.1 \pm 0.5) \times 10^{-3} \left[ \frac{150 \text{ MeV}}{m_t} \right]^2 \left( \frac{2/3}{B_k} \right) \\
\left[ \frac{f_{\bar{B}_d}^2 B_d}{(145 \text{ MeV})^2} \right]^{3/4}.
\]  

(4.8)

Agreement with experiment is good and the sign is correct. This by no means implies quantitative understanding. One potentially troublesome point should be made. It is found that the calculation of \( \epsilon' \) is insensitive to \( m_t \); whereas \( \epsilon \) is increasing rapidly with \( m_t \). However, \( \epsilon \) is fixed by experiment. This in turn implies that if the \( t \)-quark is heavy; e.g. \( m_t \geq 100 \text{ GeV} \) as preferred by \( x_d \), then the parameters in Eq. (4.5) will have to take the lower values that are allowed by experiments. This then feeds back into a smaller value of \( \epsilon' \). In short, a heavy \( t \)-quark favors a lower value of \( \frac{\epsilon'}{\epsilon} \). A second point concerns the importance of the penguin diagrams. They were first invoked to explain the \( \Delta I = \frac{1}{2} \) rule. This it fails to do.\(^{15} \)

Optimistically, one can take the view that \( \frac{\epsilon'}{\epsilon} \) is a better probe of penguins and \( \Delta I = 1/2 \) rule is the wrong place for sighting them. Certainly much more work is required.

I can now summarize the status of CP violation in SM as follows: it works very well and gives quantitative predictions. The KM paradigm is consistent with all data. It is urgent to
settle the issue of $\frac{\epsilon'}{\epsilon}$, especially its sign. The range of $m_t$ is $60 < m_t < 200$ GeV. Only four parameters need to be determined, $\rho$, $\delta$, $m_t$ and $m_H$ and the picture will be complete.

V. LAMPPPOSTS IN FAIRYLAND

Despite its remarkable success, the standard model does not look like the final theory. It contains too many free parameters. Depending on the counting one gets 21 free parameters. The Higgs mechanism is \textit{ad hoc} and suffers from naturalness and fine tuning problem. There are many models that go beyond SM. For the discussion of CP violation, I simplify the situation by classifying them into two categories. The first class is to keep the gauge symmetry as $SU(2) \times U(1)$ with the possible addition of a global $U(1)$ Peccei-Quinn symmetry.\(^{17}\) This $U(1)_{PQ}$ symmetry is involved to solve the strong CP problem.\(^{18}\) In brief the strong CP problem arises in SM due to tunneling effect between different QCD vacua. This leads to a $P$ and $T$ violating term in the SM Lagrangian which can be written as $\theta_{QCD} G_{\mu\nu} \tilde{G}^{\mu\nu}$ where $G^{\mu\nu}$ ($\tilde{G}^{\mu\nu}$) is the gluonic field tensor (dual). This induces neutron electric dipole moment (EDM) given by\(^{19}\)

$$d_n \simeq 10^{-16} \theta_{QCD} . \quad (5.1)$$

<table>
<thead>
<tr>
<th>Model</th>
<th>$kM$</th>
<th>$\text{HM}$</th>
<th>$\text{sign} \left( \frac{\epsilon'}{\epsilon} \right)$</th>
<th>$\text{solve}$</th>
<th>$d_n$</th>
<th>Ref.</th>
</tr>
</thead>
<tbody>
<tr>
<td>SM</td>
<td>yes</td>
<td>+</td>
<td>no</td>
<td>$\theta_{QCD}$</td>
<td>$\sim 10^{-31}$</td>
<td></td>
</tr>
<tr>
<td>2 Higgs explicit CPV</td>
<td>yes</td>
<td>+</td>
<td>yes</td>
<td></td>
<td>$\sim 10^{-31}$</td>
<td>Peccei-Quinn(^{16})</td>
</tr>
<tr>
<td>3 Higgs spont. CPV</td>
<td>yes</td>
<td>+</td>
<td>yes</td>
<td></td>
<td>$\sim 10^{-26}$</td>
<td>to $10^{-25}$</td>
</tr>
<tr>
<td>2 Higgs doublet &amp; 2 singlet</td>
<td>yes</td>
<td>not fixed</td>
<td>yes</td>
<td></td>
<td>$\sim 10^{-26}$</td>
<td>to $10^{-25}$</td>
</tr>
<tr>
<td>4(^{th}) generation</td>
<td>yes</td>
<td>fixed</td>
<td>no</td>
<td></td>
<td>$\sim 10^{-31}$</td>
<td></td>
</tr>
<tr>
<td>left-right symmetric</td>
<td>no</td>
<td>fixed</td>
<td>user STV probably small</td>
<td></td>
<td>$\sim 10^{-26}$</td>
<td>to $10^{-27}$</td>
</tr>
<tr>
<td>SUSY</td>
<td>yes</td>
<td>+</td>
<td>?</td>
<td></td>
<td>$\sim 10^{-27}$</td>
<td></td>
</tr>
</tbody>
</table>
Experimentally $d_n < 10^{-25}$ e.c.m. Thus $\theta_{QCD}$ has to be less than $10^{-9}$ which is unnatural since all other couplings in SM are of order $\sqrt{\alpha}$.

The second class of models involves an extension of the gauge group or the symmetry of the SM. Prominent examples are left-right symmetric models which has the gauge group $SU(2)_L \times SU(2)_R \times U(1)$ and super symmetric standard models which give a bosonic partner to every fermion in SM and vice versa. In all such models considered thus far an extended Higgs structure is required.

In the extended Higgs model where the gauge symmetry is $SU(2) \times U(1)$ one can either keep the KM phase, i.e. complex Yukawa couplings, or hypothesize that Yukawa couplings are real. In the latter case all CP violation effects arise from the Higgs potential. This is known as spontaneous CP or T violation, (STV). The Weinberg three Higgs doublets model is archetypical. One can also combine both KM phase and STV. A minimal model is constructed by Geng and Ng. The new CP violation mechanism in these type of models is always Higgs-like boson exchanges. They involve either charged Higgs boson or scalar-pseudoscalar mixings.

In $SU(2)_L \times SU(R)_R \times U(1)$ models it is customary not to include STV to give a physical relative phase between left and right $W$-bosons. The Higgs sector is then ignored for the study of CP violation. Attention is then focused on $W_L - W_R$ mixing and the phase involved.

Since models beyond SM have too many free parameters to give quantitative prediction for CP-violation phenomenology I will summarize how well they do qualitatively in Table I.

$$d_n = \begin{cases} 
-(1.4 \pm 0.6) \times 10^{-25} \text{ e-cm} & \text{Lobashov et al.} \\
-(1.2 \pm 0.6) \times 10^{-25} \text{ e-cm} & \text{Grenoble} 
\end{cases}$$

VI. SOME CP VIOLATION EXPERIMENTS BEYOND $K \to \pi \pi$

While some theorists are wandering in fairyland, experimenters are busy planning the next generation of experiments. I will give a list of what has been discussed. This is certainly an incomplete list and more details can be found in Ref. 25.

i) $K^0 \to \pi^0 e^- e^+$

This has become a favorite decay for the next generation of CP experiment. The branching ratio is expected to be small; $10^{-11} \sim 10^{-12}$. It proceeds through a CP violating one-photon exchange diagram and a CP conserving two-photon exchange diagram (see Fig. 6). It is a test of the electromagnetic penguin mechanism.

ii) $K^0_L \to \mu \bar{\mu}$

The longitudinal polarization asymmetry, $P_L$, of the final state muon defined by

$$P_L \equiv \frac{N(R) - N(L)}{N(R) + N(L)}$$

where $N(R/L)$ denotes the number of right-handed/left-handed muon, in a a CP violating observable. In SM $P_L$ is expected to be $10^{-3}$. However, if the Higgs-boson is light, this can be enhanced to $10^{-2}$ to $10^{-1}$. The lighter the Higgs boson is the larger $P_L$ will get. It is remarkable that experimenters are now measuring different quantities in rare decays.

iii) Transverse muon polarization in $K^+ \to \pi^0 l^+ \nu$

Here, the CP violating observable one hopes to measure is the transverse polarization of the lepton $S_e \cdot (\vec{p} \times \vec{K})$ where $S$ is the spin of the lepton, $\vec{p}$ and $\vec{K}$ are the momentum of...
iv) Electric dipole moment of charge leptons
These measurements are challenging and theoretically interesting. Again the SM predicts very small value coming from 3-loop effects. Large values can arise in extended Higgs models, left-right symmetric models, SUSY, and horizontal symmetry models. An important test here is to measure $\frac{d}\delta dr$ and see how this ratio will scale as a function of the lepton mass, i.e. $\frac{d\mu}{d_e} \sim \left( \frac{m_\mu}{m_e} \right)^3$ or $\left( \frac{m_\mu}{m_e} \right)^2$. Predictions of models differ in this ratio. Obviously, $d_\mu$ is free of any QCD ambiguities and thus observations of these EDM is a stringent test of electroweak models of CP violation.

In conclusion it is clear that new experiments in CP violations are necessary for a deeper understanding of CP violation in SM as well as for discovering new physics. These will require new facilities and new detectors. The next step looks difficult and challenging. As a compensation the rewards for this work will be great indeed.

ACKNOWLEDGEMENT

I would like to thank the organizers for a very stimulating conference that brings together nuclear and particle physicists to discuss many interesting and deep problems of physics.
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15. See e.g. A. Buras in Ref. 6.
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27. See G. Belanger et al., in Ref. 5.
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33. W. Morse in Ref. 25.
34. For a review see S. Barr and W. Marciano in Ref. 5.
TOWARD A MEASUREMENT OF CP VIOLATION IN HYPERON-ANTIHYPHERON SYSTEMS

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ABSTRACT

Violations of CP symmetry would be exhibited if differences exist in the parameters which describe the decays of hyperons as compared to antihyperons. Such differences are predicted in many of the models which accommodate CP violation, including the standard model. However none have ever been observed. This paper explores those experimental obstacles which must be overcome in order to realize a sensitive test of CP symmetry in hyperon-antihyperon systems.

INTRODUCTION

Violations of CP symmetry were originally discovered in the decays of neutral kaons but, even though the integrity of this symmetry has been tested in many other systems, no other violations have been found. The possibility of detecting CP violation (at predicted levels) in hyperon-antihyperon systems is explored in this paper. Such a measurement would be an outstanding, complimentary and new contribution to our present limited understanding of this odd violation of nature's symmetries. The outlined program is difficult, perhaps not even possible, but the spin / symmetry community represented at this workshop is quite accustomed to those barriers.

By observing the decays of equally well-prepared samples of hyperons and antihyperons, one can test various fundamental symmetries. For example, a comparison of the lifetimes of lambdas and antilambdas is a test of CPT, and a comparison of the decay asymmetry parameters (as outlined below) is a test of CP. The hyperon-antihyperon event samples can be prepared in a number of ways. One such method is to produce a large number of \( J/\psi \) particles by \( e^+e^- \) collisions and then to observe the hyperon-antihyperon decays (total <1\%). Another method is to produce them in antiproton-proton (\( \overline{p}p \)) collisions in either a "colliding-beam" or a "fixed-target" mode. In any case, in order to study the relevant decay properties for CP tests, polarized samples of the hyperons are required. This type of production has already been demonstrated at the low-energy antiproton ring (LEAR) at CERN. An extension of the \( \overline{p}p \) production technique, based on a hydrogen cluster-jet target intersecting an antiproton storage ring, is possible. Antiproton storage rings are presently available and possibly upgradeable to a level necessary to make the measurements. However, the detection technique which is described is equally appropriate for high-intensity extracted beams and small cells of liquid hydrogen. Secondary beam lines with sufficient antiproton flux (for example at KAON) could possibly be tailored from the outset to fit the needs of this experiment.

The PS185 collaboration\(^1\) has been using the LEAR facility with an extracted antiproton beam to study reactions of the type, \( \overline{p}p \rightarrow \overline{Y} \) where \( Y \) is a hyperon. The effort has been mainly concentrated on the \( \overline{\Lambda}\Lambda \) channel where a few times \( 10^4 \) events have been accumulated. While the motivation for these studies was mainly to learn more about the dynamics of strange-
antistrange quark production, the measurements and the experimental techniques serve nicely as a platform on which one can construct "second-generation" experiments devoted to precision tests of symmetries. A brief résumé of the PS185 results will therefore be included as reference material.

A suggestion has been made\textsuperscript{2} to study the more experimentally challenging system of $p\bar{p} \rightarrow \Xi^- \Xi^+ \rightarrow \Lambda \pi^- \bar{\Lambda} \pi^+ \rightarrow p \pi^- \pi^- \bar{p} \pi^+ \pi^+$, as a full analysis of this reaction has the potential to test a CP nonconserving quantity which is predicted to be relatively large. An outline of such an experiment, including preliminary thoughts on a necessary detector, are given. A comparison of sensitivities to CP violation in this and in the better-known $\Lambda\Lambda$ system is made. The conclusion drawn is that, at least to this point, the feasibility study should continue.

\section*{STATUS OF CP VIOLATION}

On the 25th anniversary of the discovery of CP violation in the decays of neutral kaons, one must contend with the fact that no other system has exhibited such a violation and that the only "well-known" information is contained entirely in the parameter $\epsilon$. This complex parameter measures the difference between (neutral kaon) mass and CP eigenstates with $|\epsilon| \neq 0$ indicating CP non-invariance. It can be interpreted as a measure of the CP impurity in the $K_L$ system and has the experimental value\textsuperscript{3} of $\epsilon = (2.27 \pm 0.02) \times 10^{-3}$. In the picture in which a $K_L$ is a linear combination of CP-odd ($K_2$) and CP-even ($K_1$) eigenstates, then the following diagram represents a path through which a (physical) $K_L$ particle may decay to a CP-even, two-pion final state. In the process represented by the black dot in diagram (1), a strangeness change of two units ($\Delta S=2$) takes place.

\begin{center}
\includegraphics{diagram1.png}
\end{center}

In the exact limit of the $\Delta I=1/2$ rule, this decay should proceed equally to both $\pi^+\pi^-$ and $\pi^0\pi^0$ final-state particles. Experimentally, the actual CP-odd observables which are determined are the complex quantities $\eta_{+-} = |\eta_{+-}|e^{i\phi_{+-}}$ and $\eta_{00} = |\eta_{00}|e^{i\phi_{00}}$. These are ratios of the amplitudes for $K_L$ and $K_S$ decays to $\pi\pi$ pairs; the subscripts indicate the charges of the final-state pions and the phases, $\phi_{ij}$, are equal if CPT is a good symmetry. Both $\eta_{+-}$ and $\eta_{00}$ are approximately equal to $\epsilon$ as given above.

In the "superweak" model of Wolfenstein,\textsuperscript{4} an interaction is postulated which provides a $\Delta S=2$ mechanism. It also requires that $\eta_{+-}$ and $\eta_{00}$ are exactly equal. However, within the context of the standard six-quark model, expectations exist for $\Delta S=1$ CP violation. A consequence of this is that $\eta_{+-}$ and $\eta_{00}$ are not expected to be equal. The difference from unity for the ratio of ratios, $|\eta_{00}| / |\eta_{+-}| \neq 1$, is related to a second important CP quantity termed $\epsilon'$. This quantity is proportional to the "direct" decay process (below) involving $\Delta S=1$ transitions.
Recently, ambitious experimental efforts\textsuperscript{5,6} have been mounted both at CERN and at FNAL in order to measure all four relevant decay rates in the same experiment in order to determine the quantity $\varepsilon'/\varepsilon$ which is defined as

$$\frac{\varepsilon'}{\varepsilon} = \frac{(1 - R)}{6}; \quad \text{where} \quad R = \frac{\Gamma(K_L \rightarrow 2\pi^0)/\Gamma(K_L \rightarrow \pi^+\pi^-)}{\Gamma(K_S \rightarrow 2\pi^0)/\Gamma(K_S \rightarrow \pi^+\pi^-)}.$$  

The NA31 collaboration\textsuperscript{5} at CERN has reported a value of $\varepsilon'/\varepsilon = (3.3 \pm 1.1) \times 10^{-3}$, while the FNAL E761 group\textsuperscript{6} has reported $\varepsilon'/\varepsilon = (3.2 \pm 3.0) \times 10^{-3}$. Both of these experimental results are consistent with standard model predictions and have greatly rekindled excitement in CP tests. For example, a non-zero value of $\varepsilon'$ eliminates the superweak model. Within the context of the standard model, one can use the value of $\varepsilon'$ to make predictions of the magnitude of CP violation in related systems such as in hyperon and antihyperon decays. Here only a $\Delta S=1$ process is allowed since baryon number conservation prohibits the mixing of $\Lambda$ and $\bar{\Lambda}$ particles.

The apparent agreement of these two collaborations on the value of $\varepsilon'/\varepsilon$ may be short-lived. The FNAL group has very recently announced a new value which is based on a significantly larger event sample. Their preliminary statement is that $\varepsilon'/\varepsilon$ is more closely consistent with zero.\textsuperscript{7} The resolution awaits the final analysis and a new value from the CERN team. Additionally, a third collaboration\textsuperscript{8} at CERN is also preparing to make an $\varepsilon'/\varepsilon$ measurement. They will use a decidedly different technique using $\bar{p}p$ annihilations at rest to produce "tagged" neutral kaon "beams." They will begin data collection in earnest in 1990.

In a few years, the value of $\varepsilon'/\varepsilon$ should be known reliably. To reemphasize, this number is a critical input to calculations of CP violation in hyperon-antihyperon systems. For calculations within the standard model, the value of the top quark mass is also important as it enters into the Kobayashi-Maskawa mixing matrix. This quantity, too, might be known by the time an experimental program as discussed below is launched.

**CP VIOLATION IN THE HYPERON - ANTIHYPERON SYSTEM**

The decay properties of hyperons and antihyperons should be the same in the limit of CP symmetry. These decays are weak and parity violating, manifest by an anisotropic angular distribution of the decay baryon with respect to the hyperon spin axis. If CP is violated, one would expect that the anisotropy of the hyperon decay would be different from that of the antihyperon decay. Why CP might be violated in this system may be seen by noting that there is a connection between a hyperon's decay and the kaon system described above. Diagram (3) shows how an "ordinary" CP-violating process such as those represented in diagrams (1) and (2) can be embedded into hyperon decay. In this particular representation of a lambda decay
channel, one might expect that the \( K \to \pi\pi \) piece of the diagram might have a different amplitude compared to the similar process for an antilambda decay where an antikaon is involved.

\[ (3) \]

What makes this a particularly interesting picture, as discussed by G. Miller\(^9\), includes the fact that the \( \Lambda-n-K \) and the \( n-\pi-n \) vertices are both strong; this is in contrast to a similar diagram for the electric dipole moment (EDM) of the neutron in which the lambda is replaced by a neutron in the above figure and the corresponding vertex is weak. Based on the strong vs weak argument, the hyperon process is relatively stronger by many orders of magnitude as compared to the EDM of the neutron. As pointed out at this workshop, predicted values for the EDM of the neutron are extremely small and heroic experimental efforts are still many orders of magnitude away from the "sensitive" level.

The \( K_L \to \pi\pi \) process in (3) could occur by either of the more elementary diagrams shown previously. However, the process which governs diagram (1) is proportional to \( M_{12}/(q^2 - M^2) \) where \( M_{12} \) is the matrix element for mixing \( K^0 - \bar{K}^0 \) particles. For the decays of free kaons, \( q^2 = M^2_L \) so that the denominator nearly vanishes. This results in \( \varepsilon \) being large. But the mass of the off-shell kaon in diagram (3) is significantly different from \( M_L \) such that the CP-violating diagram (1) does not contribute. Instead, the "direct" diagram (2) is the most important one and thus the link to the knowledge of \( \varepsilon' \).

John Donoghue and collaborators have presented\(^2\) a "model-independent" analysis of CP violation in hyperon-antihyperon decays. Their studies lead to a set of experimental "tests" that can be made and to corresponding sets of (model dependent) predictions of these effects. Their work is summarized below.

A hyperon decays to a "final" baryon and a pion, \( B_1 \to B_f \pi \), via a combination of S and P waves. The decay can be described by these amplitudes and their relative phase. If the amplitudes are denoted by "S" and "P", then the decay parameters, in the formalism of Lee and Yang\(^10\), are given as

\[
\Gamma = |S|^2 + |P|^2, \quad \alpha = \frac{2\text{Re}(S^*P)}{|S|^2 + |P|^2}, \quad \beta = \frac{2\text{Im}(S^*P)}{|S|^2 + |P|^2}.
\]

Here \( \Gamma \) is the decay rate for the process of interest. The quantity \( \alpha \) describes the angular distribution of the decay baryon, \( B_f \), with respect to the hyperon spin axis. If \( \theta \) represents the angle between the polarization direction for produced hyperons and the outgoing baryon direction, then the angular distribution in the hyperon rest frame is

\[ I(\theta) = I_0 \left[ 1 + \alpha \gamma P \gamma \cos \theta \right], \]
where $P_Y$ is the magnitude of hyperon polarization and the measurable quantity is the up-down asymmetry of the decay baryon with respect to this vector. The parameter $\beta$ is related to the polarization of $B_f$ itself along the normal to the plane containing the $B_f$ momentum and the hyperon polarization:

$$P_{B_f} = \frac{\beta_Y P_Y \sin \theta}{1 + \alpha_Y P_Y \cos \theta}.$$

In the limit of CP conservation, $\alpha = -\bar{\alpha}$ and $\beta = -\bar{\beta}$.

The amplitudes which describe the decay of a hyperon and antihyperon, have the following structure:

$$S_Y = + S_1 \exp(i(\delta_1^s + \phi_1^s)) + S_3 \exp(i(\delta_3^s + \phi_3^s))$$

$$S_Y = - S_1 \exp(i(\delta_1^s - \phi_1^s)) - S_3 \exp(i(\delta_3^s - \phi_3^s))$$

$$P_Y = + P_1 \exp(i(\delta_1^p + \phi_1^p)) + P_3 \exp(i(\delta_3^p + \phi_3^p))$$

$$P_Y = + P_1 \exp(i(\delta_1^p - \phi_1^p)) + P_3 \exp(i(\delta_3^p - \phi_3^p)),$$

where $\delta$ and $\phi$ are the strong and weak phases, respectively, and the subscript "1" or "3" refers to either a $\Delta I=1/2$ or $\Delta I=3/2$ transition. The sign of the weak phases reverses in the antiparticle description which causes the values of $\alpha$ and $\beta$ to differ from $\bar{\alpha}$ and $\bar{\beta}$.

Donoghue et al.\textsuperscript{2} construct difference tests of the decay parameters for which a non-zero value indicates CP violation. They construct the tests in terms of the amplitudes in order to predict the magnitude of violation. The expressions below are evaluated to leading order.

$$\Delta = \frac{(\Gamma - \bar{\Gamma})}{\Gamma + \bar{\Gamma}} = \frac{S_3}{S_1} \sin(\delta_3^s - \delta_1^s) \sin(\phi_3^s - \phi_1^s) + S \leftrightarrow P$$

$$A = \frac{(\alpha + \bar{\alpha})}{(\alpha - \bar{\alpha})} = \sin(\delta_1^p - \delta_1^s) \sin(\phi_1^s - \phi_1^p)$$

$$B' = \frac{(\beta + \bar{\beta})}{(\alpha - \bar{\alpha})} = \sin(\phi_1^s - \phi_1^p).$$

The last factor in each expression takes into account the weak phases on which the expected CP violation is dependent. Noting that the $\Delta I=1/2$ amplitudes are about 20 times stronger than the $\Delta I=3/2$ amplitudes due to the $\Delta I=1/2$ rule, and that $\sin(\delta_1) \approx 1/10$, Donoghue finds that $\Delta = A/10 = B'/100$. Thus, CP violation as measurable in the quantities $\Delta$ and "A" is reduced in magnitude by factors which are "uninteresting" while the ratio $B'$ has the feature that it isolates the weak phases and is correspondingly the test with the largest expected non-zero value (model independent).

To evaluate these tests in a theoretical model, one needs the parameters $\epsilon$ and $\epsilon'$, values for the relevant hadronic matrix elements, and the top quark mass. From the present knowledge of these parameters and uncertainties with the models, one can expect\textsuperscript{2} that
\( A_\Lambda = (0.3 - 3.0) \times 10^{-4} \) (with \( A_\Xi \approx 20\% \) larger) and \( B' = 10A \). The preliminary calculations of Miller and Iqbal (for these same quantities but using a different approach) give similar magnitudes for the violations.\(^9\) It should be noted that a result of \( e'/e = 0 \) is technically allowed in the standard model description and thus such a result does not by itself guarantee CP conservation in hyperon-antihyperon systems. In fact, the predicted magnitude of the violation would be quite uncertain in such a case and could possibly be large.\(^9\) Allowing for the present non-zero value of \( e'/e \), it would be eminently worthwhile if a measurement of "\( A' \)" to an accuracy of \( 10^{-4} \) or \( B' \) to \( 10^{-3} \) could be made. These are the benchmarks which will be assumed below.

**ONGOING MEASUREMENTS IN THE \( \Lambda\Lambda \) SYSTEM**

The reaction \( \bar{p}p \to \Lambda\Lambda \) is being thoroughly explored at LEAR by the PS185 collaboration.\(^1\) The technique is to produce \( \Lambda\Lambda \) pairs using an extracted beam of cooled antiprotons and then to observe their characteristic two-"\( Vee \)" charged decays with a set of tracking wire chambers. The ideal antiproton momentum for a dedicated CP experiment on the \( \Lambda\Lambda \) system is about 1.650 GeV/c, which is enough above the threshold that the cross section is high, yet still below the next hyperon-antihyperon threshold (\( \Lambda\Sigma + \Lambda\Xi \)). The proper lifetime, \( \epsilon \tau \), of a \( \Lambda \) is 7.9 cm and hence \( \Lambda \)'s recoil downstream approximately 3 to 10 cm before they decay. At this momentum the outgoing decay proton is constrained to a lab angle of less than 30\(^\circ\). The branching ratio to the all-charged final state, \( \Lambda \to p\pi^- \), is 64\%.

![Fig. 1. A plan view of the PS185 detector. The \( \bar{p} \) beam enters the target cells (blowup) from the left and the apparatus is triggered when any of the veto counters which surround the cells do not fire and when the hodoscope, which follows the wire chambers, does fire.](image-url)
Although the total $\bar{p}p$ cross section is large (about 100 mb), and hyperon production is only a tiny fraction of the produced events, very high quality data can still be accumulated at high beam intensities ($10^6$ incident $\bar{p}$/sec) by use of a "charged-neutral-charged" trigger. The process detected is $\bar{p}p \rightarrow \Lambda\Lambda \rightarrow \bar{p}\pi^+\pi^-$. An incoming $\bar{p}$ enters a target system consisting of four individual 2.5 mm long cylinders of CH$_2$, each surrounded by thin veto scintillators which do not trigger on outgoing neutrals. The hyperon decay volume is sandwiched between the target and a scintillator hodoscope, which detects the outgoing charged particles from the hyperon decay. These elements form the online trigger.

The decay volume is filled with both MWPC planes and drift chamber planes in order to record the charged-particle tracks. Finally, a low-field solenoidal magnet, mounted behind the hodoscope, contains three additional drift chamber planes which measures the curvature of the penetrating baryons and thus tags the respective lambda and antilambda. The detector is shown in Fig. 1.

The differential cross section at 1.695 GeV/c incident antiproton momentum is shown in Fig. 2. An anisotropy with a strong forward peaking of the antilambda in the center-of-mass system is seen. The forward-peaked differential cross section means that the $\Lambda$'s travel, on average, further downstream before decaying than the $\Lambda$'s. This means that there is an inherent asymmetry built into the reaction process which cannot be avoided. Therefore, in constructing a dedicated CP-testing apparatus, one must be able to insure that decays measured throughout the detector are equally well understood.

Fig. 2. PS185 results for the reaction $\bar{p}p \rightarrow \Lambda\Lambda$ at 1.695 GeV/c incident antiproton momentum. The abscissa is the cosine of the antilambda in the center-of-mass system.
The polarization of the outgoing lambdas is based on the intensity distribution of the decay protons with respect to the hyperon spin axis. This can be determined if the lambda sample is polarized with respect to some measurable plane. In the strong (parity-conserving) process, $\bar{\text{p}} p \rightarrow \Lambda \bar{\Lambda}$, the hyperons can only be polarized along the direction which is normal to the production plane; the production plane contains the interaction point and the decay vertices of the lambda and antilambda). The product $\alpha_{\Lambda} P_\Lambda$ is measured by determining the asymmetry in the number of decay protons which are emitted above vs below this plane and $P_\Lambda$ is derived by dividing the asymmetry by the known value for $\alpha_\Lambda = 0.642 \pm 0.013$.

In Fig. 3a, the "raw" measured asymmetries $\alpha P_\Lambda$ and $\bar{\alpha} P_\bar{\Lambda}$, are plotted vs the cosine of the antilambda production angle in the center-of-mass system. The fitted lines are included as guides and one can see that the polarization of the lambda is relatively large and negative but crosses zero for forward angles. The CP measurement sensitivity depends directly on the degree of polarization. It can be seen that maximum sensitivity is realized where the differential cross section is isotropic (Fig. 2) and minimal sensitivity is realized where most of the events fall -- in the forward region.

The PS185 collaboration has begun to evaluate the quantity "$A_{\Lambda}$" for the higher-statistics event samples. The ratio $(\alpha P_\Lambda + \bar{\alpha} P_\bar{\Lambda}) / (\alpha P_\Lambda - \bar{\alpha} P_\bar{\Lambda})$ can be constructed directly from the data of Fig. 3a. This ratio reduces to "$A_{\Lambda}$" when charge invariance in the strong production mechanism is invoked ($P_\Lambda = P_\bar{\Lambda}$) as shown in Fig. 3b. Thus the result is obtained directly and without modifications to the analysis. A result of this work,\textsuperscript{12} based on a total sample of 16,000 events, is $A_{\Lambda} = 0.023 \pm 0.057$.

While the PS185 method is quite straightforward and is developed from a very clean trigger concept, it is limited in maximum event rate accumulation. When commissioned, the rate was limited to the intensity of the extracted beam at LEAR. However, since the ACOL at CERN has been completed, beams of up to $5 \times 10^6 \bar{\text{p}} / \text{sec}$ are available. The antiproton beam passes directly through the center of the PS185 detector, including the wire chamber. All wires which "see" the primary beam are made inactive over a length comparable to the beam spot size. This method works well up to about $10^6 \bar{\text{p}} / \text{sec}$. Above that rate, prohibitive noise in the chambers and excessive trigger rates for the electronics influence the quality of the data and the actual throughput of the data acquisition environment. The situation can be rectified with appropriate, but major, upgrades.

Another problem in extracted beam experiments is the maximum density of protons which can be provided in a target "cell" whose length is short compared to the decay distance of the lambdas. The PS185 CH$_2$ targets total 1 cm in length. This ultimately limits the luminosity. Any future experiment would require a (less dense) liquid hydrogen target. This does not easily permit the introduction of very thin veto counters to surround the cell(s).

The extension of the PS185 technique to a sensitivity of, at best, $\sim 5 \times 10^{-3}$ might be possible. A sample of more than $10^6$ events would be required along with a sizeable amount of computer time for the analysis. Beyond this, an alternate technique should be considered. Accordingly, a feasibility study of an experiment which would be designed from the beginning to explore CP violation in a hyperon-antihyperon system at a level in line with predictions has begun among a few members of the collaboration.\textsuperscript{13} One such proposal is to build a symmetric version of the PS185 apparatus using a cluster-jet target and high-speed electronic detectors throughout in order to obtain the very large sample of $\Lambda \bar{\Lambda}$ events which are required in order to measure "$A_{\Lambda}$" to a level of interest. Looking into a different reaction is also a possibility. This is discussed next.
Fig. 3. (a) Asymmetry measurements for the reaction $\bar{p}p \rightarrow \bar{\Lambda}\Lambda$ at 1.695 GeV/c. The fitted lines are third-order polynomials which are included to guide the eye to the trend of the polarization. In (b), the quantity "$A_{\Lambda}$" is calculated from the data.

DISCUSSION OF THE $\Xi\Xi$ SYSTEM

The CP test, $B'$, is predicted to be about 10 times larger than "$A\Lambda" but to test it one needs to determine the outgoing baryon polarization. In the case of lambda and antilambda decays, this means measuring the outgoing polarization of the proton and antiproton which would be impossible to do symmetrically for the particle and antiparticle. However in the more exotic $\Xi^{-}\Xi^+$ system, the $\Lambda$ is the "decay baryon" and its polarization is revealed in the usual manner through the parity violation asymmetry in its weak decay. By observing the complete decay chain, one can measure the quantities $A_\Xi$, $A_\Lambda$, and $B'_\Xi$. The three experiments could be done simultaneously. With the device which is described below, other hyperon-antihyperon reactions could be included in the online trigger as well.
To produce $\Xi\Xi$ pairs in this manner requires an antiproton beam with a momentum above 2.6 GeV/c. This exceeds the capabilities of LEAR, whose highest momentum is 2.0 GeV/c, but is within the working range of the FNAL antiproton accumulator. The peak of the total cross section is expected to be only about 2 $\mu$b (to be compared with about 80 $\mu$b for the $\bar{\Lambda}\Lambda$ system) so a high-luminosity environment is also required. An appropriate setup could be achieved using the existing cluster-jet target which presently intersects the FNAL antiproton storage ring. Such configurations are presently being exploited by collaborations mounting other experiments at both LEAR\textsuperscript{14} and at FNAL\textsuperscript{15}.

The event topology is shown in Fig. 4 with the relevant planes of interest indicated. One seeks to measure the angular distribution of decay protons having the general form

$$I(\theta) = I_0(1 + b \cos \theta),$$

where $\theta$ is the angle of the proton relative to the normal to the $\vec{S}_\Xi - \vec{p}_\Lambda$ plane; $\vec{S}_\Xi$ is the polarization direction of the $\Xi$, and $\vec{p}_\Lambda$ is the direction of the $\Lambda$. The coefficient $b$ is given by

$$b_\Xi = P_\Xi \alpha_\Lambda \beta_\Xi$$

and similarly for $b_\Lambda$. The difference between $b_\Xi$ and $b_\Lambda$ is proportional to $B'$, the CP non-conserving quantity. For a sample of $N$ events, the statistical uncertainty on $B'$ is

$$\sigma_{B_\Xi} = \frac{1}{P_\Xi \alpha_\Lambda \alpha_\Xi} \sqrt{\frac{3}{2N}}.$$ 

By comparison for the $\bar{\Lambda}\Lambda$ experiment, $b_\Lambda$ is given by $P_\Lambda \alpha_\Lambda$ and for $N$ events the statistical uncertainty on $A'$ is given by

$$\sigma_{A_\Lambda} = \frac{1}{P_\Lambda \alpha_\Lambda} \sqrt{\frac{3}{2N}}.$$ 

The number of events required to achieve the desired sensitivities in the $\bar{\Lambda}\Lambda$ and $\Xi\Xi$ experiments depends, in part, on the polarization of the produced hyperons and on other "dilution" factors which enter into the experimental observables as described above. The absolute error ($\pm 0.057$) on the value of $A_\Lambda$ as determined in the 16,000 event sample of PS185 is consistent with the above expressions and an average polarization of $|P_\Lambda| = 0.26$. Extrapolating to an absolute error of $10^{-4}$ implies that an event sample exceeding $5 \times 10^8$ is required. The production polarization of the particles (if any) in the $\Xi\Xi$ system are not known. The optimistic assumption that $P_\Xi = 0.26$, coupled with the expression for the error on $B'$ implies that $2.5 \times 10^8$ events are needed for this experiment to obtain an absolute error on $B'$ of $10^{-3}$. Since the cross section for $\Xi\Xi$ production is lower by a factor of 40 compared to $\bar{\Lambda}\Lambda$ production one can see that the 20-times fewer events required is offset by the rate of accumulation. The two experiments should be explored in parallel -- detector and accelerator considerations have not yet been folded in.

With an upgraded FNAL main injector,\textsuperscript{16} a peak luminosity of $5 \times 10^{31}$ sec$^{-1}$ cm$^{-2}$ could be achieved at 3.5 GeV/c using the present cluster-jet target and antiproton accumulator combination. This leads to about 100 events per second. Assuming one could develop a detector with an overall collection efficiency of 10%, then the needed $2.5 \times 10^8$ events (or...
more) would require about 10 months of running time. These numbers are given in order to illustrate that such a program is reasonable.

\[ \bar{p}p \rightarrow \Xi^- \Xi^+ \rightarrow \Lambda \pi^- \bar{\Lambda} \pi^+ \rightarrow p \pi^- \bar{p} \pi^+ \pi^+ \] in the laboratory frame.

Note that the polarization of the outgoing $\Xi$ is allowed only in the "y" direction. The relevant measurement is based on whether the final decay proton is "aligned" or "antialigned" with respect to the normal vector to the plane formed by "y" and the direction of the lambda, indicated as $S_x \times \vec{p}_\Lambda$ in the figure.

**A STRAW-MAN DETECTOR FOR CP STUDIES**

Some considerations for an appropriate detector for the $\Xi \Xi$ measurement are outlined next. This detector is also *entirely appropriate* for a dedicated $\bar{\Lambda} \Lambda$ study but the discussion is focused on $\bar{p}p \rightarrow \Xi^- \Xi^+ \rightarrow \Lambda \pi^- \bar{\Lambda} \pi^+ \rightarrow p \pi^- \bar{p} \pi^+ \pi^+$ since it is the harder experiment. At 3.5 GeV/c antiproton momentum, the $\bar{p}p$ total cross section is about 70,000 $\mu$b. This means an interaction rate of 3.5 MHz at a luminosity of $5 \times 10^{31}$ sec$^{-1}$ cm$^{-2}$. The immediate consequence is that the first-level online trigger must be both very fast and rather selective. Here, "first level" means decisions in less than 100 nsec. A solution to the trigger is proposed which is entirely based on conventional photomultiplier tube signals and fast electronic logic.

The final state of the $\Xi \Xi$ reaction has six charged particles. However, six-pronged events are quite common in such $\bar{p}p$ collisions, including the non-annihilation, "direct"
The production of $\bar{p}p \rightarrow \bar{p}p\pi^+\pi^-\pi^+\pi^-$. Monte-Carlo studies indicate that less than 5% of the events of interest have charged particles backward of $80^\circ$ in the laboratory frame. A fast veto of any charged particles in the backward direction is included in the design.

The $pp$ total cross section produces a significant number of events with gammas from the copious production of neutral pions. The desired event features no neutrals. A gamma veto shield around the entire detector can potentially eliminate a large fraction of the online triggers.

Finally, and most importantly, none of the final state particles has a trajectory which intersects the jet-beam crossing. The final state pion and proton tracks appear to emerge from an infinite number of "virtual sources" due to the delayed decays of the hyperons. A detector which is "hard-wired" to be sensitive to all of these allowed trajectories is difficult to conceive. But essentially all of the background reactions involve particles which do point to the interaction region, and the number of paths (or "roads") which point there is relatively finite. These paths can be programmed into the detector design and put into veto.

Thus one would like to develop a trigger based on only six charged particles in the forward direction, none of which points to the interaction region. The realization of the first-level trigger concepts is included in Fig. 5. The overall geometry of the detector is cylindrical, surrounding the beam pipe with the tracking elements downstream of the gas jet. A segmented "all-veto" shield surrounds the beam line backwards of $80^\circ$ and is made from a combination of lead and scintillator having a depth of about 4 radiation lengths and read out on the upstream end by phototubes. Both charged particles and gammas will trigger this device equally well with any hit rejecting the event.

Two concentric shells of projective scintillator strips surround the tracking chambers. The respective strips in each layer are placed in coincidence with each other, forming fast coincident signals which indicate that a charged particle was emitted from the interaction region. This geometry is often used for calorimeter modules (called "pointing geometry"), but here is simply made from relatively thin layers of scintillator which are separated in distance, and segmented such that ample angular resolution is achieved. In the limit of perfect resolution, all six charged particles pass this device without triggering a coincidence, however, many of the decays will occur close to the target and undoubtedly some delayed-decay tracks will satisfy the coincident condition. The actual trigger used can allow $\leq 2$ coincident paths for example, but this needs to be carefully simulated along with the actual granularity of the device.

So far, only backgrounds have been addressed. Elementary Monte-Carlo calculations of the reaction of interest have been made with 3.5 GeV/c incident antiprotons. The results are shown in Figs. 6. The $\Xi$ is constrained to a laboratory angle of less than $24^\circ$ (Fig. 6a) and it decays close to the beam axis, due in part to the production angle and to the fact that $c\tau_\Xi = 4.9$ cm. The radius of the decay (Fig. 6b) indicates that the majority of decays would occur inside the standard 2-cm radius beam pipe which exists at the FNAL accumulator.

The decay distributions of the lambdas with respect to the radius from the beam axis and the distance downstream of the interaction region are shown in Fig. 6c and 6d. These plots indicate the maximum extent of the tracking detector which is required, approximately 50 cm in length and 30 cm in diameter. The particles which pass this detector travel mainly downstream with the angle of the decay proton constrained to within $40^\circ$ of the beam axis, Fig. 6e. Therefore the wires of the tracker should point along the radius of the cylinder as in a radial drift chamber. Some of the pions do travel at steeper angles (Fig. 6f). These would be better measured if a second tracker, made from straw drift tubes with wires pointing along the
beam direction, surrounded the radial drift chamber. The two chambers would team up to
detect particles which pass roughly orthogonal to the orientation of their respective wires.

Finally, as in the PS185 setup, a low-field solenoidal magnet would be placed at the
downstream end of the apparatus to serve as a baryon number identifier. The beam pipe would
pass through the center of this magnet, but would be shielded by soft iron from any disturbing
field. The magnet volume would be instrumented with at least five planes of drift chambers.
The back of the magnet could also incorporate segmented liquid Cherenkov counters which
would help to identify protons from pions. But otherwise, every attempt is made to minimize
the mass seen by the charged particles.

It should be noted that, with the exception of the downstream solenoid, the entire
detector is non-magnetic. The events are reconstructed from the kinematics imposed by the
opening angles of the decay kinks and Vee's. This avoids the precession of the spins of the
Ξ's and Λ's before decay.

Fig. 5. A "cutaway" view of the outlined detector. In this view, a gas jet crosses the
storage ring vacuum tube and the antiproton beam enters from the left. The "all veto"
covers the angular region from $80^\circ$ to $135^\circ$. The package of detectors immediately
downstream of the intersection includes, from the inside, a radial drift chamber, an annulus
of straw-tube chambers, and the projective "veto" scintillator hodoscope. A large solenoid
magnet is positioned furthest downstream and contains additional drift chambers which are
not shown.
SOME WORRIES AND SOME CURES

One experimentally-induced asymmetry which cannot be avoided is based on the fact that the detector is entirely made of matter, yet is expected to detect both matter and antimatter without bias. Antimatter annihilations occur preferentially in the more massive parts of the detector. This must be minimized by the choice of very thin detectors. The question also arises whether there is selective depolarization of the positive $\Xi$'s vs the negative $\Xi$'s in crossing the beam pipe?\(^\text{17}\)

Must the $\Xi$ decay vertex occur inside a tracking chamber? The answer to this question is probably "yes" since knowledge of the orientation of the production plane is important. The limitations on the minimum size of storage-ring beam tubes are related to the physical dimensions of the trajectories of the accumulated (or injected) beam before phase-space cooling. When the beam is ready to be used (cluster jet turned on) it typically has small enough dimensions to satisfy the experimental conditions. Perhaps these storage rings, or future versions of them, might incorporate a "spur track" where the detector would be located. After the antiproton beam is cooled inside the main ring, the beam could be rerouted via fast-kicking
magnets along the "alternate" route, presumably through the very small beam pipe at the position of the detector. The device would be costly, but it may be paramount to the realization of the experiment as a whole -- an item that is probably necessary in order to exploit the luminosity gained from this setup as compared to a conventional extracted-beam approach. At KAON, a dedicated extracted beam line could be prepared that would not have such problems. The question to the designers of this machine is, "what is the maximum intensity of antiprotons that can be delivered, and with what momentum resolution?". An alternate question to be considered, for LEAR, for FNAL, or for KAON, is, "what would it cost to build a dedicated storage ring devoted only to this problem?".

The list of potential troubles and false asymmetries in such an experiment is long and continues to grow (see Hamann14). But there are a few positive features too. The desired signal is an asymmetry involving the knowledge of several production and decay planes for each event. Even the simple measurement of the (yet unknown) polarization of the outgoing $\Xi$ depends on measuring an asymmetry. In order to determine whether the asymmetries which are observed are real signs of physics or just false signals, the analysis of the events can be "replayed" with the introduction of arbitrary orientations of the production (or decay) planes. For example, the production plane can be "assigned" the horizontal direction in the analysis program. The observed polarization of the outgoing $\Xi\Xi$ pairs should then vanish when averaged over the full event sample. The key is that both particle and antiparticle samples are prepared in equal numbers and that analyzed events are only those events which contain both particle and antiparticle topologies completely. A decay-antiproton annihilation in some part of the detector eliminates the whole event. Such asymmetries should be able to be sorted out using clever analysis procedures.

SUMMARY

The preceding discussion represents an interim report of a study aimed at measuring CP-odd signals in hyperon-antihyperon systems. Having reviewed some of the requirements to measure the "more sensitive parameter" $B'$ in the $\Xi\Xi$ system compared to "$A$" in the $\Lambda\Lambda$ system, one is led to believe that a dedicated $\Lambda\Lambda$ experiment is more promising. The event topology is simpler, the cross section is much higher, and the experiment can be run below the thresholds of "look-alike" background reactions. To achieve an interesting sensitivity requires less running time -- and that is based on unproven and optimistic factors folded in for the $\Xi\Xi$ system. In reality, the $\Lambda\Lambda$ system probably wins in other ways too.

The detector which was outlined is appropriate for either measurement. A dedicated (but yet non-existent) antiproton storage ring may also be required. LEAR is not presently useful due to technical reasons. With an aggressive approach to the design of beam lines for KAON, one might be able to design a high-intensity antiproton beam which could be used. If so, it would surely be one of the most outstanding experiments to be mounted there.

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NUCLEAR ENHANCEMENT OF TIME-REVERSAL NON-INVARIANCE

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ABSTRACT

I consider the five-fold correlation experiment planned at LANSCE, and present, for the corresponding asymmetry observable, a detailed prediction of the root-mean square sensitivity realised in practice to a nuclear time-reversal noninvariant interaction.

It is now an article of nuclear folklore that low-energy nuclear reactions [i.e. compound-nucleus (CN) reactions] can display enhanced sensitivity to the violation of fundamental space-time symmetries.\(^1\) The optimal conditions for such measurements have been clarified by several authors, notably Bunakov and co-workers.\(^2\) I have previously discussed\(^3\) the advantages and disadvantages of this approach, as well as the potential of a contrasting method, the application of statistical reaction theory\(^4\) in the interpretation of measurements. It, in particular, allows one to make simple and detailed predictions of (and not just speculate on) the magnitude of effects. In this contribution, I illustrate this for an example appropriate to this session, namely the five-fold correlation (FC) measurements planned by the TRIPLE group.\(^5\)

These experiments involve the study of the transmission of transversely-polarised neutrons through spin-aligned targets, specifically $^{165}$Ho and $^{235}$U (since, of suitable targets, these are the easiest to construct). The natural observable is the asymmetry

$$\epsilon = \frac{1}{P_n} \frac{I_+ - I_-}{I_+ + I_-}.$$  \hspace{1cm} (1)

Here, $P_n$ is the neutron polarisation and $I_+$ ($I_-$) is the intensity of transmitted neutrons for the spin "up" ("down") configuration — i.e. incident neutron polarisation parallel (antiparallel) to $\mathbf{k} \times \mathbf{c}$, $\mathbf{k}$ being the beam wavevector and $\mathbf{c}$ a unit vector along the direction of the "crystal" axis of the target. Any deviation of $\epsilon$ from zero would be a manifestation of a Parity (P) conserving violation of Time-reversal (T) invariance. Even if one supposes that such processes do occur, the characteristics of CN reactions imply that $\epsilon$ would vary randomly with changing energy. In fact, its energy average $\langle \epsilon \rangle$ would vanish identically. The basic nontrivial measure of the strength of fluctuations in $\epsilon$, and hence of its potential sensitivity to a P-even T-odd nuclear interaction, is the variance $\langle \epsilon^2 \rangle^{1/2}$. It is on the behaviour of this variance which I will report.

A crucial ingredient in determining the behaviour of $\langle \epsilon^2 \rangle^{1/2}$ realised in practice, is the effect of restricted experimental energy resolution. I treat the case of "thin" targets — i.e. those for which $n \left( \sigma - \langle \sigma \rangle \right) \ll 1$, where $n$ is the number of target nuclei per unit area and $\sigma = (\sigma_+ + \sigma_-)/2$, $\sigma_+$ ($\sigma_-$) being the total spin up (down) cross-section.\(^6\) Then, following for example Ref. 7, the asymmetry at an energy $E_0$ is

$$\epsilon(E_0) \sim -\frac{n}{2} \int_{-\infty}^{\infty} dE' R(E' - E_0) \left[ \sigma_+(E') - \sigma_-(E') \right],$$  \hspace{1cm} (2)

where the function $R$ is assumed to describe the energy resolution characteristics of the experimental apparatus.\(^8\)
Via the optical theorem, \( \delta \sigma \equiv (\sigma_+ - \sigma_-)/2 \) is related to the anti-symmetric part \( S^{(a)} \) of the (elastic) \( S \)-matrix by an expression of the form

\[
\delta \sigma = 2 k^2 \sum_{a < b} A_{ab} \text{Im}(S_{ab}^{(a)}) ,
\]

where \( a, b \) denote elastic neutron channels, \( S_{ab}^{(a)} = (S_{ab} - S_{ba})/2 \) and \( A_{ab} \) is real and energy-independent with \( A_{ab} = -A_{ba} \). It accommodates the polarisation states of beam and target, a statistical spin factor and various angular momentum recoupling coefficients. Exploiting the randomness of \( S^{(a)} \), the corresponding expression for \( \langle \epsilon^2 \rangle \) is

\[
\langle \epsilon^2 \rangle = 2 \frac{\pi^2}{k^4} \sum_{a < b} (A_{ab})^2 \int_{-\infty}^{\infty} dE \, W(E) \text{Re}[\langle S_{ab}^{(a)}(-E) (S_{ab}^{(a)}(E))^* \rangle_e] ,
\]

where \( W(E) \) is the weight

\[
W(E) \equiv 2 \int_{-\infty}^{\infty} dx \, R(x - E) R(x + E) ,
\]

which is normalised in the case that \( R \) is normalised.

The non-trivial component of Eq. (4) is the auto-correlation function

\[
v_{ab}(E) \equiv \langle S_{ab}^{(a)}(-E) (S_{ab}^{(a)}(E))^* \rangle_e .
\]

In line with statistical reaction theory, it involves an average \( \langle \ldots \rangle_e \) at fixed energy over an ensemble of Hamiltonian matrices with randomly distributed entries.\(^4\)\(^9\) It is in this way that the fluctuation properties of CN states with changing energy are simulated, the equivalence with explicit energy-averaging resting on the well-founded assumption of Ergodicity. The strength of the P-even T-odd interaction \( H^{(T)} \) in the CN system is conveniently parametrised in terms of the associated spreading width \( \Gamma^{(T)} \) for states of a given spin and parity: if their average spacing is \( D \),

\[
\Gamma^{(T)} \equiv \frac{2\pi}{D} \langle (H^{(T)}_{\mu_\mu'})^2 \rangle ,
\]

where \( H^{(T)}_{\mu_\mu'} \) denotes any matrix element of \( H^{(T)} \) in an arbitrary orthonormal basis. All the other physically relevant parameters of the model are tied down by specifying, in addition to the average level-spacings \( D \), the energy-averaged \( S \)-matrix \( \langle S \rangle \), which at these energies is responsible for shape-elastic scattering. Actually, it is more convenient to work with the related transmission coefficients \( T_a \equiv 1 - |\langle S_{aa} \rangle|^2 \). It is the fact that this input is among some of the best known nuclear data that enables one to make detailed predictions.

Using the exact super-symmetric generating-functional formalism adapted to the CN problem in Ref. 10, it is possible to analytically reduce the ensemble average \( v_{ab}(E) \) to a compact expression in which the dependence on the above average parameters is explicit:\(^1\)\(^1\)

\[
v_{ab}(E) = T_a T_b f_{ab}(E, \{T_c\}) \frac{2\pi \Gamma^{(T)}}{D} ,
\]

where \( f_{ab}(E, \{T_c\}) \) is the three-dimensional integral

\[
f_{ab}(E, \{T_c\}) = \frac{1}{8} \int_0^\infty d\lambda_1 \int_0^\infty d\lambda_2 \int_0^1 d\lambda \, \mu(\lambda) \pi(\lambda) p_{ab}(\lambda) \exp \left[ -2i \frac{\pi E}{D} (\lambda_1 + \lambda_2 + 2\lambda) \right] ,
\]
in which
\[
\mu(\lambda) \equiv \frac{(1 - \lambda)\lambda|\lambda_1 - \lambda_2|}{(\lambda_1(1 + \lambda_1)\lambda_2(1 + \lambda_2))^{1/2}(\lambda + \lambda_1)^2(\lambda + \lambda_2)^2},
\]
\[
\pi(\lambda) \equiv \prod_{c \text{ open}} \frac{1 - T_c \lambda}{(1 + T_c \lambda_1)^{1/2}(1 + T_c \lambda_2)^{1/2}},
\]
and
\[
p_{ab}(\lambda) = \{y_a(1) y_b(2) + 2y_a [y_b(1) + y_b(2)] + 3y_a y_b\} + \{a \leftrightarrow b\}
\]
with
\[
y_c(k) = \lambda_k(1 + \lambda_k)/(1 + T_c \lambda_k)
\]
and
\[
y_c = \lambda(1 - \lambda)/(1 - T_c \lambda).
\]
Observe that, consistent with its role as auto-correlation function, \(v_{ab}(E)\) decreases with increasing \(E\) [because of the oscillatory factor in Eq. (6)].

In the applications I have in mind (i.e. to the resolved resonance regimes of \(^{165}\)Ho and \(^{235}\)U), the open channels comprise a large number of gamma channels, and a few weakly-absorbing neutron and, in the case of \(^{235}\)U, fission channels, meaning that all transmission coefficients \(T_c \ll 1\). In my computations, I use the exact result for \(f_{ab}(E, \{T_c\})\) [Eq. (6)], but it is instructive to note that, under these conditions, \(f_{ab}(E, \{T_c\})\) is in effect a function of just one variable, namely the combination \(\bar{t} \equiv t(1 + 2iE/\Gamma)\), where \(t \equiv \sum_c T_c\), the sum [like the product in \(\pi(\lambda)\)] being restricted to the open channels of the same spin and parity as \(a\) and \(b\), and \(\Gamma\) is the average width of resonances of this spin and parity. The observation hinges on two facts: first, to obtain the contribution to \(v_{ab}(E)\) of leading order in each \(T_a\), it suffices to set
\[
\pi(\lambda) \rightarrow \exp \left[-\frac{t}{2}(\lambda_1 + \lambda_2 + 2\lambda)\right]
\]
and
\[
p_{ab}(\lambda) \rightarrow 2[\lambda_1(1 + \lambda_1) + 2\lambda(1 - \lambda)] \cdot [\lambda_2(1 + \lambda_2) + 2\lambda(1 - \lambda)] - 2\lambda^2(1 - \lambda)^2;
\]
second, in the resolved resonance regime, \(t\) can be identified with the ratio \(2\pi \Gamma/D\). Recalling that \(v_{ab}(E)\) is the auto-correlation function of two matrix elements of \(S^{(a)}\) separated by an energy of \(2E\), this property of \(f_{ab}(E, \{T_c\})\) implies that the corresponding correlation length is, to a good approximation, \(\Gamma\), which is again intuitively plausible.

One can now state qualitatively the conditions under which the effects of restricted energy resolution should become important. If one makes the reasonable assumption that \(R\) is a peaked function of full-width half-maximum (FWHM) \(\Delta E\), then the weight \(W\) in Eq. (4) is also peaked and of FWHM \(\sim \Delta E\). As noted above, the function \(\text{Re}[v_{ab}(E)]\) in Eq. (4) decreases with increasing \(E\), and, given the identification of \(\Gamma\) as the correlation length, the decrease becomes significant for \(E \gg \Gamma\). Consequently, it is when \(\Delta E \gg \Gamma\) that fluctuations in \(\epsilon\) should be noticeably washed out. For the LANSCE facility where the TRIPLE group's FC measurements will be performed, this corresponds to neutron energies as little as 50 to 100 eV: with the projected neutron flight path of 60 m, \(\Delta E\) equals \(\Gamma^{(165}\)Ho) (\(\approx 0.06\) eV) and \(\Gamma^{(235}\)U) (\(\approx 0.3\) eV) at neutron energies of 55 eV and 75 eV, respectively.

In dealing with \(\Gamma(T)\), there are two options. The one would be to parametrise \(\left<\epsilon^2\right>\) in terms of \(\Gamma(T)\). This would have the advantage of making clear the natural division between reducing a null measurement of \(\epsilon\) to a bound on \(\Gamma(T)\), and extracting information on more fundamental parameters from this bound. However, the information content of \(\Gamma(T)\) is crucial to the motivation for the FC measurements. A first attempt at a systematic reduction of \(\Gamma(T)\)
has been undertaken by French et al. For the purpose of order of magnitude estimates, the result can be summarised as

$$\frac{1}{2\pi} \Gamma^{(T)} \sim (10^5 \text{eV}) \alpha_T^2,$$

where $\alpha_T$ is the strength of the P-even T-odd part of the effective nucleon-nucleon interaction (relative to the standard P-even T-even part). A similar reduction to a bare strength parameter has, to my knowledge, not been attempted, but the example of the parity-violations observed in neutron transmission experiments suggests that it should be $\sim \alpha_T$. I shall present parametrisations of $\langle e^2 \rangle$ in terms of $\alpha_T$, using Eq. (7).

Figure 1 contains a plot of the root-mean-square (rms) sensitivity parameter

$$S \equiv \frac{1}{np} \langle e^2 \rangle^{1/2}$$

as a function of neutron energy for both $^{165}$Ho and $^{235}$U with Lorentzian and Gaussian resolution functions $R$. In addition to $\alpha_T$, I have divided out the dependence on the two parameters under the exclusive control of the experimentalist: the number of target nuclei per unit area $n$ and the polarisation parameter

$$p \equiv \frac{3}{2} \left( \frac{5}{2} \right)^{1/2} i_{10} i_{20},$$

where $i_{10}$ and $i_{20}$ are the first and second order irreducible statistical tensors describing the degree of polarisation of the beam and alignment of the target, respectively, in their symmetry frames. At each energy, $\Delta E$ is set equal to the value appropriate to the LANSCE facility (Flight path = 60 m). Note that, in Fig. 1, $S$ has units of barns (amounting to the choice of inverse barns for $n$).

The most striking feature of the prediction for $S$ is the non-trivial dependence on the neutron energy $E$. Below about 50 eV, $S$ rises as $\sqrt{E}$. Above about 100 eV ($\Delta E \gg \Gamma$), $S$ decreases as $1/E$ approximately. The former behaviour is governed by the penetrability factors which kinematically suppress the FC as $E \to 0$, and the latter is a consequence of the restricted resolution, and is consistent with the analytic estimate that it be given by $\sqrt{E}/\Delta E$. (For LANSCE, $\Delta E \propto E^{3/2}$ for $E \gg 1$.) Were the effect of restricted energy resolution omitted, $S$ would continue to increase as $\sqrt{E}$ for the energies depicted. So, for the LANSCE facility, inclusion of restricted energy resolution has a very important effect on the energy dependence.

The choice of $R$, on the other hand, would appear not to be crucial. The examples in Fig. 1 illustrate that the global shape of $S$ is insensitive to the choice of $R$. The primary difference is in the magnitude of $S$ at energies for which the effects of restricted energy resolution become important. Note that, for both $^{165}$Ho and $^{235}$U, the ratio of the Gaussian $R$ ($R_G$) to Lorentzian $R$ ($R_L$) predictions for $S$ is approximately 1.5. That the $R_G$ prediction should be greater is not surprising: given that $R_G$ dies off exponentially for large arguments and $R_L$ only as a power law, the former “smears” out less. In fact, it would seem that the difference between the $R_L$ prediction (which has the advantage that excellent simple closed-form approximations to it are available) and the prediction for other choices of $R$, can be related to the ratio of the respective resolution functions evaluated at the origin. For example,

$$R_G(0) / R_L(0) = \sqrt{\pi} \ln 2 \approx 1.5.$$

In discussing the implications of Fig. 1 for the TRIPLE group experiment, I adopt the $R_G$ prediction (as this is the more realistic choice of resolution function). If one takes the
Figure 1: Sensitivity $S$ versus neutron energy for the FC measurement at LANSCE (Flight path = 60 m). The four curves are: dot-dashed (dotted) — $^{165}$Ho ($^{235}$U) target [with circles (diamonds) — Gaussian $R$; without — Lorentzian $R$).
Table 1: Target characteristics for FC measurements

<table>
<thead>
<tr>
<th></th>
<th>$^{165}$Ho</th>
<th>$^{235}$U</th>
</tr>
</thead>
<tbody>
<tr>
<td>$np \ [b^{-1}]$</td>
<td>$\sim 0.1$</td>
<td>$\sim 0.01$</td>
</tr>
<tr>
<td>$\langle \epsilon^2 \rangle^{1/2}/ \alpha_T$</td>
<td>$\sim 1$</td>
<td>$\sim 0.1$</td>
</tr>
</tbody>
</table>

results at face value, then one would conclude that there is an optimal window in which the measurements should be performed. A reasonably conservative definition of this window is that it spans the interval form $\sim 75 \text{ eV}$ to $\sim 400 \text{ eV}$. It must be stressed that the potential sensitivity of $\epsilon$ to $\alpha_T$ is greatest in this window, precisely because the fluctuations in $\epsilon$ are greatest. Note, in particular, that, below this window, the sensitivity of $\epsilon$ is severely suppressed. As regards the choice of target, the sensitivity parameters for $^{165}$Ho and for $^{235}$U differ by only a factor of two. From a theoretical viewpoint, $^{165}$Ho is marginally superior. The deciding factor is then the value of the geometrical parameter $np$. The approximate projected values of $np$ for the TRIPLE measurement are given in Table 1. It is obvious that, in this respect, $^{165}$Ho is markedly superior. My advice to TRIPLE is to forget about using $^{235}$U, and focus all its resources on making a really good measurement with $^{165}$Ho. Note that the ultimate maximal rms sensitivity of $\epsilon$ to $\alpha_T$ with $^{165}$Ho is "only" of order unity. Another of life’s hard facts.

These results apply strictly to the somewhat idealised limit of “thin” targets. However, their net implication is that measurements are best performed at neutron energies where $\Delta E \approx \Gamma$ — i.e. at energies just below the point at which the consequences of restricted energy resolution become important. This conclusion should thus survive the inclusion of the effects omitted in this analysis, as should the prediction of the order of magnitude of $S$ at these optimal energies.

Let me close by broaching very briefly two general issues.

First, what ultimately is the advantage of FC neutron transmission measurements over the established P-even T-odd test furnished by the comparison of the cross-sections for an inelastic reaction with unpolarised participants and its inverse (Traditional Detailed Balance)? An obvious advantage of traditional detailed balance is the greater flexibility with respect to choice of target and reaction. Actually, traditional detailed balance measurements are also potentially more sensitive to $\alpha_T$ by as much as an order of magnitude.\(^3\) However, past experience in traditional detailed balance experiments (which have established a bound on $\alpha_T$ of $\sim 10^{-3}$)\(^{13}\) indicates that it will be difficult to achieve a precision which allows one to set a bound of less than $10^{-4}$. The advantage then of FC measurements rests on the level of precision which can be achieved.

The second issue concerns the likelihood of the FC analysing power being $\sim 10^3 \alpha_T$, the conservative estimate for FC measurements performed under optimal conditions.\(^{14}\) For a target about one mean-free path thick, this would correspond to finding $\epsilon \sim 10^3 \alpha_T$. Given
Table 2: Sensitivity estimates for Parity measurements

<table>
<thead>
<tr>
<th></th>
<th>$^{139}$La</th>
<th>$^{93}$Nb</th>
<th>$^{127}$I</th>
<th>$^{81}$Br</th>
<th>$^{238}$U</th>
</tr>
</thead>
<tbody>
<tr>
<td>P-violation</td>
<td>Yes</td>
<td>No</td>
<td>No</td>
<td>Yes</td>
<td>No</td>
</tr>
<tr>
<td>$S_P/10^6$</td>
<td>$\sim 10^{-1}$</td>
<td>$\sim 10^{-2}$</td>
<td>$\sim 10^{-2}$</td>
<td>$\sim 10^{-2}$</td>
<td>$\sim 10^{-1}$</td>
</tr>
</tbody>
</table>

the above results on the rms of $\epsilon/\alpha T$ in $^{165}$Ho and $^{235}$U, one can see that such a fluctuation in $\epsilon$ is, at least, $\sim 10^3$ ($\sim 10^4$) $\sigma$ event in $^{165}$Ho ($^{235}$U): these numbers are a measure of how improbable it is that the optimal conditions for a FC measurement identified in Ref. 14 will be realised. A more quantitative statement is not possible at present, but an interesting cross-check with the P-violation measurements by the Alfimenkov group is. For P-violation, the sensitivity parameter which replaces $S$ is

$$S_P \equiv \frac{1}{n p'} \frac{\langle \epsilon^2_p \rangle^{1/2}}{\alpha_P},$$

where $\epsilon_p$ is the asymmetry appropriate to P-violation measurements, $p'$ the corresponding polarisation parameter and $\alpha_P$ the relative strength of the P-odd part of the effective nucleon-nucleon interaction. One can ask whether, by considering $S_P$, one can understand the success rate of the Alfimenkov group. To see a P-violation at the 1% level, one would expect $\epsilon_P \sim 10^6 n p' \alpha_P$ or $S_P \sim 10^6$. (I assume $n p' \sim 0.1$ and take $\alpha_P \sim 10^{-7}$.) Table 2 lists some of the nuclei studied by the Alfimenkov group, whether a P-violation was seen or not and gives an order of magnitude estimate of the corresponding value of $S_P/10^6$. Use of $S_P$ is only partially successful. If one supposes that $S_P$ explains the $^{139}$La, $^{93}$Nb and $^{127}$I results, then one cannot understand why an effect was seen in $^{81}$Br and was not seen in $^{238}$U. Resolution of this dichotomy is a priority.

Interaction with C. Gould, F. Fröhner and G. E. Mitchell is gratefully acknowledged. I would like to thank TUNL for its hospitality during preparation of this manuscript.

REFERENCES


5. C.f. C. Gould's contribution.

6. More generally, the restriction on the validity of my analysis is that \( n (\sigma - \lambda \langle \sigma \rangle) \ll 1 \), where \( \lambda \) is some arbitrarily chosen constant.


8. I assume that effects not catered for in Eq. (2) [such as Doppler-broadening of \( \sigma_\perp \)], although they would further quench the fluctuations in \( \epsilon \), are of less importance. They can be included by resorting to Monte Carlo simulations.

9. Details of the distribution are discussed and motivated in the second reference of 4 above.


13. The most satisfactory discussion of the best bound to emerge to date is given by J. B. French et al., Tests of Time Reversal Invariance in Neutron Physics, edited by N. R. Roberson et al. (World Scientific, Singapore, 1987), p. 80. They quote a confidence level of 85% for the bound, but this presupposes that the statistical reduction of the raw data was correct. This would seem not to have been the case (H. L. Harney, private communication). An improved analysis would worsen the confidence level.

14. C.f. first reference in 2 above.
SEARCH FOR PARITY AND TIME REVERSAL SYMMETRY VIOLATION USING POLARIZED NEUTRON-NUCLEUS SCATTERING*

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ABSTRACT

In this paper the TRIPLE collaboration (Time Reversal Invariance and Parity at Low Energies) reviews parity (P) and time reversal (T) violation experiments that are possible with epithermal polarized neutrons. The scientific motivation for the experiments is discussed and sensitivities that can be achieved at the Los Alamos Neutron Scattering Center (LANSCE) are given. Strategies for dealing with systematic errors in transmission measurements are presented.

INTRODUCTION

Spallation neutron sources provide unique opportunities for studying Standard Model predictions of P- and T-violation in epithermal neutron-nucleus scattering. Due to the small level spacings in the compound nucleus, the mixing between energy levels is much larger than for levels near the ground state. As a result, the dependence of a scattering process on a symmetry-violating observable exhibits a large enhancement. The measurements involve transmission of polarized neutrons through targets which are unpolarized, polarized or aligned. The total cross section will depend on Cartesian invariants made up from the dynamical variables s (neutron spin), k (neutron momentum), and I (target spin). A complete list is given in Table I. The terms of particular interest are the lowest order symmetry violating terms in s; these are the helicity term s·k (P-odd, T-even), the three fold correlation (TC) s·(I×k) (P-odd, T-odd), and the five fold correlation (FC) s·(I×k)·(I·k) (P-even, T-odd).

* presented by C. R. Gould
The existence of P-violation in the nucleon-nucleon (N-N) interaction is of course well established. The characteristic size is $\sim 10^{-7}$; considerable effort has been devoted to extracting the strengths of the P-violating meson exchange couplings from p-p scattering data and from $\gamma$ decays in light nuclei. 1. The interest in compound nucleus (CN) resonance studies lies not so much in showing that P violation occurs, but rather in clarifying how the underlying symmetry violation in the effective N-N interaction manifests itself in the complicated many body nuclear system. This general issue is of considerable current interest in quantum chaos2-4 and has been studied in recent years by the techniques of Random Matrix Theory. P-violation is characterized by a spreading width $\Gamma_{PV}$ which is related to the variance of the CN matrix elements of the many body P-violating Hamiltonian $H^P$. The spreading width is in turn related to $\alpha_P$, the ratio of the P-odd to P-even effective N-N interaction. We expect $\alpha_P$ to be of order $10^{-7}$. The goals of the P-violation studies are to determine $\Gamma_{PV}$, either via isolated resonance or energy averaged measurements, and to validate the procedure relating $\Gamma_{PV}$ to $\alpha_P$. This procedure is of great interest in its own right and will be crucial to the interpretation of T-violation experiments where the underlying T-violating N-N potential is not known. Successfully testing the conversion from experimental P-violation in neutron-nucleus scattering to P-violation in the N-N interaction would provide confidence that a similar procedure would yield meaningful limits for T-violation when transforming from the nucleon-nucleus regime to the nucleon-nucleon regime.

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Table I. Terms in forward scattering amplitude$^a$. Four additional terms enter each time I increases by half; these correspond to the interaction of $I$ with $k$, and the interaction of $s$ with the three directions $I, k$ and $I \times k$.

<table>
<thead>
<tr>
<th>$I$</th>
<th>Term</th>
<th>Odd in</th>
</tr>
</thead>
<tbody>
<tr>
<td>0</td>
<td>$1$</td>
<td>P</td>
</tr>
<tr>
<td></td>
<td>$s \cdot k$</td>
<td></td>
</tr>
<tr>
<td>$\geq 1/2$</td>
<td>$(s \cdot I)(I \cdot k)$</td>
<td>P</td>
</tr>
<tr>
<td></td>
<td>$s \cdot I$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$s \cdot k$</td>
<td>PT</td>
</tr>
<tr>
<td></td>
<td>$s \cdot (I \times k)$</td>
<td></td>
</tr>
<tr>
<td>$\geq 1$</td>
<td>$(s \cdot I)(I \cdot k)$</td>
<td>P</td>
</tr>
<tr>
<td></td>
<td>$s \cdot k$</td>
<td>P</td>
</tr>
<tr>
<td></td>
<td>$(I \cdot k)^2$</td>
<td>P</td>
</tr>
<tr>
<td></td>
<td>$s \cdot (I \times k)(I \cdot k)$</td>
<td>T</td>
</tr>
<tr>
<td>$\geq (1+n/2)$</td>
<td>${(I \cdot k)^n}$</td>
<td></td>
</tr>
</tbody>
</table>

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$^a$ from reference 14
The helicity parity-violation test is carried out with a polarized beam and an unpolarized target. Figure 1 shows schematically how such measurements are performed. A polarizing filter prepares a neutron beam in either of two helicity states. The neutron beam is transmitted through a thick unpolarized sample and the asymmetry is determined from the count rates in the two helicity states. Remarkably large parity-violating asymmetries have been observed for a number of nuclei, the most spectacular case being the 0.734 eV p-wave resonance in $^{139}$La for which the effect is between 7 and 10% (Refs. 5 and 6).

While $P$-odd $T$-violation is inferred from the decays of the neutral kaon system, there is no experimental evidence for $P$-even $T$-violation. For this latter case, we define a $T$-violating spreading width $\Gamma^T\nu$ by analogy with $\Gamma^{PV}$ and relate it to $\alpha_T$, the ratio of the $P$-even, $T$-odd to $P$-even, $T$-even effective $N\cdot N$ interaction. The quantity $\alpha_T$ is known to be $\leq (1-2) \times 10^{-3}$ from the analysis of the fluctuations of level spacings in heavy nuclei, and $\leq (2-3) \times 10^{-3}$ from an analysis of the $^{24}$Mg($\alpha,p$)$^{27}$Al experiment of Blanke et al. The planned experiments appear to be able to reach bounds of $\alpha_T \leq 10^{-6}$, corresponding to two to three orders of magnitude improvement in these limits.

A $P$-odd $T$-violation test is carried out with a polarized target and a $P$-even $T$-violation test with an aligned target. In both cases the neutron beam is polarized perpendicular to the ($\mathbf{l},\mathbf{k}$) plane. The tests are more complicated than the parity measurements because of problems relating to precession of the neutron spin in the target. An experiment is illustrated in Figure 2. When $s$ is reversed the $T$-violating correlations change sign and the spin dependence of the transmission is measured. However, a transmission difference does not necessarily imply $T$-violation, because changing the state of the polarizer does not correspond to time reversing the apparatus.

Fig. 1. Experiment to search for helicity $s\cdot\mathbf{k}$ interaction with an unpolarized target.

Fig. 2. Experiment to search for $T\cdot C$ or $F\cdot C$ interaction with a polarized or aligned target.
Stodolsky\textsuperscript{9} showed how to deal with this by adding an analyzer between the polarized target and the detector. The dependence of the transmission on the neutron spin is measured by simultaneously reversing the states of the polarizer and the analyzer. The two states of the apparatus are now time reversal partners of one another and detailed balance is tested. Later Kabir\textsuperscript{10} showed that a comparison of induced polarization and asymmetry could be made with a single polarizer/analyzer, and that this test was fully equivalent to the double polarizer/analyzer detailed balance test.

Both T-violation experiments are conceptually feasible, but the P-odd test suffers from a particularly severe problem, precession of the neutron spin in the magnetic and pseudomagnetic fields of the polarized target. In a finite size target the neutron will precess many times and the T-violating signal will be averaged out. This problem is essentially absent for an aligned target (no external magnetic field and no pseudomagnetic fields). Precession is due only to the P-violating terms of even order in $I (\langle s I \rangle (I \cdot k)$ for example). As a result, it is possible to identify P-even T-violation not only from the polarizer/analyzer comparison, but also from measurements of the target thickness dependence of the FC signal. Simulation of T-odd effects by the sequential action of two T-even interactions depend quadratically on the target thickness, while true T-violation depends only linearly on target thickness.

**THEORETICAL BACKGROUND**

The quantities to be determined are the spreading widths of the symmetry-violating Hamiltonians. For the P-even, T-odd Hamiltonian $H^T$ this is

$$\Gamma_{TV} = 2\pi \langle (H_{\mu\mu'}^T)^2 \rangle / D^{J\pi}$$

where $\langle (H_{\mu\mu'}^T)^2 \rangle$ represents a local energy average of matrix elements squared and $D^{J\pi}$ is the average spacing of the CN levels $\mu, \mu'$ of the given $J^\pi$ mixed by $H^T$. The spreading width is related to $\alpha_T$, the ratio of the P-even, T-odd to P-even, T-even effective nucleon-nucleon interaction\textsuperscript{2,3}

$$\Gamma_{TV} \sim 2\pi (10^5 \text{ eV}) (\alpha_T)^2$$

A similar spreading width can be defined for parity mixing between states due to a P-odd Hamiltonian $H^P$:

$$\Gamma_{PV} \sim 2\pi \langle (H_{\mu\mu'}^P)^2 \rangle / \sqrt{D^{J\pi} D^{(-J\pi)}}$$

and assuming the generality of the expression used by French et al. $\Gamma_{PV}$ is related to $\alpha_P$, the ratio of the P-odd, T-even to P-even, T-even effective N-N interaction via

$$\Gamma_{PV} \sim 2\pi (10^5 \text{ eV}) (\alpha_P)^2$$

The P-violation experiments consist of measurements of the helicity dependence of the neutron total cross section at epithermal neutron energies in heavy nuclei. At resonance, the analyzing power $Q$ is defined by $\sigma^\pm = \sigma_0 (1 \pm Q)$ where $\sigma^\pm$ is the cross section for +(-) helicity neutrons and $\sigma_0$ is the unpolarized neutron resonance cross section. A T-violating analyzing power can be defined similarly. In this case $\sigma^\mp$ denote total cross sections for neutrons polarized parallel (antiparallel) to $k \times I$.

The expression $\delta \sigma = \sigma^+ - \sigma^-$ for P-violation has been given by many authors for the case of s- and p-wave interfering resonances.\textsuperscript{11} For target spin $I = \frac{1}{2}$ and compound nucleus spin $J = 1$, we have, for example,
where $\Delta E_S = E_S - E$, $x_j$ is the entrance channel mixing parameter, $V = H_{\mu\mu}^P$, and $\sigma_p$ is a Breit-Wigner like cross section for an isolated $p$-wave resonance:

$$\sigma_p = 4\pi\chi^2 \frac{2J+1}{(2J+1)(2I+1)} \frac{\Gamma_p \Gamma_n}{(E_p - E)^2 + (\Gamma_p / 2)^2}$$

The analyzing power at the peak of the $p$ wave resonance is thus of the form

$$Q \sim \frac{V}{|E_S - E_p|} \sqrt{\frac{\Gamma_n}{\Gamma_p}}$$

The factor $V / |E_S - E_p|$ is called the dynamical enhancement. Clearly energy spacings in the compound nucleus are much smaller than typical single particle energies. On the other hand, $V$ in the compound system is also smaller than that for the single particle system. If the compound nuclear wavefunction has $N$ components, and a random phase approximation is assumed, then $V$ is expected to be reduced by $\sqrt{N}$ while energy spacings are reduced by $N$. This leads to a dynamical enhancement of order $\sqrt{N}_{12}$. The specific value depends on the level density -- typical estimates lead to an enhancement of $10^3$.

In the same spirit as that employed for the parity violation analysis, one can introduce a $T$-violating Hamiltonian $H^T$ with matrix elements $W = H^T_{\mu\mu}$, and carry through a similar procedure. The essential property of $W$ is that it is purely imaginary. For the $P$-even, $T$-odd case, one can calculate the difference in total cross section for the FC for a single $p$-wave resonance or for two interfering $p$-wave resonances. Detailed results are summarized by Gould et al.\textsuperscript{14} The cross section difference for two interfering $J = 3$ $p$-wave resonances for an $I = 7/2$ holmium target is (at the center of resonance 2)

$$\delta\sigma = \hat{t}_{20}(7/2) (x_{j}^{(1)} y_{j}^{(2)} - x_{j}^{(2)} y_{j}^{(1)}) 4\pi\chi^2 \frac{W}{E_1 - E_2} \sqrt{\frac{\Gamma_{p1} \Gamma_{n2}}{\Gamma_{p2}}}$$

where $\hat{t}_{20}(7/2) \sim 0.7$ is the statistical tensor describing the alignment of the holmium target, and $x_j$ and $y_j$ are entrance channel mixing parameters. The $T$-violating analyzing power is

$$Q \sim \frac{W}{|E_{p1} - E_{p2}|}$$

There is no structural enhancement factor $\sqrt{\Gamma_n / \Gamma_p}$. On the other hand, the argument employed with parity violation indicates that this test of $T$-invariance still benefits from dynamical enhancement. Significantly lower limits on $T$ violation in the parity-even nucleon-nucleon interaction are thus attainable.

Energy averaged measurements determine the variance of the cross section asymmetry averaged over the spectrum, on and off resonance.

$$\langle (\delta\sigma)^2 \rangle = \langle (\sigma^+ - \sigma^-)^2 \rangle$$

where for parity, $\sigma^+(-)$ denote total cross sections for $+(-)$ helicity neutrons, and for time reversal, $\sigma^+(-)$ denote total cross sections for neutrons polarized parallel (anti parallel) to $k \times I$. These measurements are discussed by Davis.\textsuperscript{15} The sensitivity is not as great as for
optimal pairs of resonances, but the measurements have the advantage of probing the systematic response of the compound system to an underlying symmetry violation. The interpretation of a null result is therefore simpler than for the case of a single resonance where the relevant spectroscopic information must be known and where, even then, the measurement is sensitive to only a single value of a stochastically distributed CN matrix element.

SENSITIVITY OF LANSCE TESTS OF P AND T

The LANSCE facility for testing P and T violation is described by Bowman et al. It consists of an LMN polarized proton target, a magnetic spin flipper system and neutron detector stations at flight paths of 11, 25 and 60 m. A schematic diagram is shown in Figure 3.

![Schematic of LANSCE Beam Line](image)

To estimate the sensitivity we need the neutron count rate as a function of energy. We consider an 18" diameter by 4 cm thick $^{10}$B detector with efficiency ~ 0.65, situated at 60 m. The effective solid angle is $2.7 \times 10^{-6}$ sr, less than the geometric solid angle because we view the $13 \times 13$ cm$^2$ neutron moderator through the 10 cm$^2$ aperture of the LMN polarizer. The polarizer transmission is 0.14, the neutron polarization $P_n \sim 0.5$, and we consider a sample of thickness $n = 0.06$ atoms/barn and transmission 0.55. (Thicker samples are possible for parity violation experiments. This transmission corresponds to the aligned holmium target presently under construction at TUNL.) For a 50 $\mu$A beam, the neutron count rate is

$$N \sim \frac{2.0 \times 10^{12}}{E} \times (0.65)(0.14)(0.55)(2.7 \times 10^{-6})$$

$$\sim \frac{2.7 \times 10^5}{E} \text{ neutrons/eV-sec}$$

A new butanol polarizer is envisioned which will have a 30 cm$^2$ aperture and a higher transmission (0.30 compared to 0.19). The count rate compared to LMN will be a factor of 6 larger, and the polarization of the neutrons will also be higher, typically 0.85.

P- and T-violation measurements for isolated resonances are performed on p-wave resonances. We restrict our analysis to neutron energies for which $\Gamma_{np} \ll \Gamma_p \sim \Gamma_\gamma$, and for which the experimental resolution $\Delta_{exp} \leq \Gamma_p$. Typically this means $E \leq 100$ eV. The peak cross section difference (for neutron polarization $P_n$) determines an experimental analyzing power $Q$ given by

$$\delta \sigma = 2 P_n \sigma_p (\text{max}) \times Q$$

which is related to the symmetry violating matrix elements via
\[ Q \sim \frac{H^T_{\mu \mu'}}{|E_{p1} - E_{p2}|} \quad \text{for T-violation} \]

and

\[ Q \sim \frac{H^p_{\mu \mu'}}{|E_{s} - E_{p}|} \sqrt{\frac{\Gamma^n_s}{\Gamma^n_p}} \quad \text{for P-violation} \]

Neglecting Doppler broadening and taking the spin statistical factor \( \sim 0.5 \),

\[ \sigma_p (\text{max}) \sim \frac{2\pi}{k^2} \frac{\Gamma^n_p}{\Gamma_p} \]

The bound on the analyzing power \( Q \) in \( t \) seconds is

\[ \delta Q \sim \frac{1}{(P_n \sigma_p (\text{max}) \sqrt{N \Delta E t})} \]

We take the experimental measuring interval \( \Delta E = \pi \Gamma_p \), where \( \Gamma_p \) is the width of the \( p \)-wave resonance.

From the spreading width relation given earlier, we have

\[ H^T_{\mu \mu'} \sim x_{\mu \mu'} 10^3 \left( \frac{D^{J\pi}}{10} \right)^{1/2} \alpha_T \]

where \( D^{J\pi} \) is in eV and \( x_{\mu \mu'} \) is a Gaussian distributed variable of unit variance. At a 95% confidence level, a bound \( \delta Q \) implies

\[ \alpha_T \leq 3 \times 10^{-3} \delta Q (|E_{p1} - E_{p2}|)^{1/2} \quad \text{for T-violation} \]

Similarly

\[ \alpha_p \leq 3 \times 10^{-3} \delta Q (|E_{s} - E_{p}|)^{1/2} \sqrt{\frac{\Gamma^n_p}{\Gamma^n_s}} \quad \text{for P-violation} \]

In both cases a smaller than average spacing will enhance the sensitivity (by the square root of the spacing). In general \( \alpha_p \sim 1 / \sqrt{\frac{\Gamma^n_p}{\Gamma^n_s}} \), and strong interfering s and p wave resonances will give the best sensitivity.

These spectroscopic parameters \( (|E_{s} - E_{p}|, \Gamma^n_s, \Gamma^n_p, \text{etc.}) \) will vary randomly with choice of resonance (or pair of resonances), and much of the potential of our experiments rests on the fact that there may be fortuitous combinations of these parameters giving rise to unusually large enhancements. Nevertheless, for the purpose of discussion, we shall estimate "reasonable" values for bounds on \( \alpha_T, \alpha_p \) by assuming that these parameters have their average values. Thus the energy differences between resonances are set equal to the average s-wave spacing: \( |E_{s} - E_{p}| \sim |E_{p1} - E_{p2}| \sim D_0 \). Neutron widths are evaluated from the strength functions \( S_\perp \) and average spacings \( D_\perp \)

\[ \Gamma^n_s \sim S_0 D_0 E^{1/2} \quad \text{and} \quad \Gamma^n_p \sim S_1 D_1 E^{3/2} \quad (0.088 A^{2/3} \times 10^{-6}) \]

For convenience we tabulate the values used\textsuperscript{17}:
This procedure gives the following bounds on $\alpha_T$ and $\alpha_P$ ($n = 0.06$ atoms/b except for $^{235}\text{U}$ T-violation where $n = 0.006$ atoms/b, $E$ in eV and $t$ in seconds):

<table>
<thead>
<tr>
<th></th>
<th>$\alpha_P$</th>
<th>$\alpha_T$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{139}\text{La}$</td>
<td>$6 \times 10^{-6} \sqrt{\frac{E}{t}}$</td>
<td>$5 \times 10^{-3} \sqrt{\frac{1}{t}}$</td>
</tr>
<tr>
<td>$^{165}\text{Ho}$</td>
<td>$2 \times 10^{-5} \sqrt{\frac{E}{t}}$</td>
<td>$2 \times 10^{-2} \sqrt{\frac{1}{t}}$</td>
</tr>
<tr>
<td>$^{235}\text{U}$</td>
<td>$1.0 \times 10^{-4} \sqrt{\frac{E}{t}}$</td>
<td>$4 \times 10^{-1} \sqrt{\frac{1}{t}}$</td>
</tr>
<tr>
<td>$^{238}\text{U}$</td>
<td>$4 \times 10^{-6} \sqrt{\frac{E}{t}}$</td>
<td>--</td>
</tr>
</tbody>
</table>

Parity-violation measurements are most sensitive at lowest energies, while T-violation measurements have the same sensitivity at all energies ($\delta Q$ independent of energy). P-violation studies must reach $\alpha_P < 10^{-7}$ to see effects. This is readily attainable in $10^6$ seconds up to a few hundred eV in $^{238}\text{U}$. Holmium is less satisfactory, and $^{235}\text{U}$ appears unsuitable for isolated resonance tests because the p-wave strength is fragmented over so many states and because the width of each state is large. (In fact for equal strength functions $\alpha_P$ scales as $\sqrt{\frac{\Gamma_P}{D}}$). T-violation experiments in aligned $^{165}\text{Ho}$ will set bounds on $\alpha_T$ which approach a few $\times 10^{-5}$ in $10^6$ seconds. $^{235}\text{U}$ is not competitive since a thick target cannot be prepared. $^{139}\text{La}$ shows good sensitivity, but there is no evidence that two suitable p-wave resonances exist.

With the new butanol polarizer, the sensitivities are about a factor of four smaller. A long run of $10^7$ seconds will set bounds approaching $10^{-6}$ on $\alpha_T$ for holmium, a factor of $10^3$ smaller than the best current limit of $10^{-3}$.

The sensitivity of an energy average test of T-violation in holmium is also excellent. At around 100 eV Davis\textsuperscript{15} finds the bound on $\alpha_T$ to be of the order of the error in the experimental asymmetry, $\alpha_T \leq \delta \varepsilon$. With the new polarizer we will count $6 \times 2.7 \times 10^4$ neutrons per second in a 10 eV range. Hence a bound of a few $\times 10^{-6}$ is achievable in $10^7$ seconds, comparable to the ‘typical’ resonance sensitivity discussed above.
SYSTEMATIC ERRORS

In this section we discuss systematic errors associated with transmission measurements of P- and T-violation. We consider instrumental effects, as well as physical effects arising from interactions of the neutron spin. For T-violation, we discuss only the five fold correlation (FC). The systematic errors associated with the three fold correlation (TC), although tractable, are more formidable.\textsuperscript{18}

The experiments will be performed using a double modulation technique. The neutron spin will be reversed every few seconds by the use of a magnetic "spin flipper" and at intervals of tens of hours by changing the pumping frequency of the dynamically polarized proton filter. The spin flipper reversals are not as precise as the polarized proton filter reversals, but have the advantage that they can be carried out in milliseconds as compared to hours.

Instrumental systematic errors arise from the fact that the spin flipper fringe field, although small, might effect the photomultipliers of the detector, and from the fact that the spin flipper reversals are not exactly 180°. Both these effects are cancelled by an eight step reversal sequence and the double modulation technique.

The spin flipper's operation depends on using adiabatic and diabatic magnetic field transitions. If a neutron spin is in an adiabatically (slowly) varying magnetic field the projection of the spin on the field direction is a constant:
\[
\frac{dB}{dt} \ll \omega_L B
\]
If on the other hand the field changes diabatically (suddenly) the spin direction does not change and the projection of the spin on the field direction is reversed:
\[
\frac{dB}{dt} \gg \omega_L B
\]

The neutron spin is created in the homogeneous field of the polarized proton filter and is parallel or anti-parallel to the field depending on the radio frequency pumping field. The longitudinal (z) fringing field adiabatically guides the longitudinal neutron spin to the spin flipper. The spin flipper has a time independent longitudinal magnetic field which reverses sign at the midpoint. A transverse (y) field is turned on or off to reverse or not reverse the spin. When the transverse field is off the neutron spin diabatically passes through the field reversal region, and emerges without reversing its direction. When the transverse field is on (up or down) the spin adiabatically follows the magnetic field direction and rotates by 180 degrees in a distance of one meter. The spin reversal is not exactly 180 degrees because the magnetic field configurations in the reversal and non-reversal configurations are different.

The neutron spin is reversed equally well by an up or down transverse field, and the eight step sequence will consist of the transverse field off, up, down, off, up, off, off, down. To the extent that the detector response is a linear function of the magnetic field, this reversal sequence eliminates the influence of the field on the detector. Furthermore, the use of this reversal sequence eliminates the effects of slow drifts up to terms of order \((dt / \tau)\) \(^2\) where \(dt\) is the measurement time in one spin state and \(\tau\) is a time characteristic of the drifts.

The transmission through the target depends linearly on each component of the neutron average spin vector, and data will be accumulated in eight step sequences for both polarizer states. Let \(S(P,F)\) be the average spin vector at a particular neutron time of flight. The quantity \(S(P,F)\) is a function of the polarizer state \(P\) (pumping frequency), either A or B, and of the flipper state \(F\) (transverse field), either off, up or down. The four spin vectors for which the transmissions are measured are:
The spin flipper approximately reverses the spin direction and therefore: \( S_1 \sim -S_2 \) and \( S_3 \sim -S_4 \). Since changing the polarizer state exactly reverses the spin direction, \( S_1 = -S_3 \) and \( S_2 = -S_4 \).

For small changes in the average spin vector the transmission will consist of a spin independent transmission \( T_0 \) plus a linear function of the average spin vector. There will also be possible contributions \( D(F) \) which represent the influence of the flipper stray field in state \( F \) on the detector. The four measured transmissions are \( T(1 \ldots 4) \)

<table>
<thead>
<tr>
<th>Pumping Frequency</th>
<th>Flipper Field</th>
<th>Spin Components</th>
</tr>
</thead>
<tbody>
<tr>
<td>A</td>
<td>off</td>
<td>( S_1 = S(A,\text{off}) )</td>
</tr>
<tr>
<td>A</td>
<td>(up+down)/2</td>
<td>( S_2 = 1/2(S(A,\text{up}) + S(A,\text{down})) )</td>
</tr>
<tr>
<td>B</td>
<td>off</td>
<td>( S_3 = S(B,\text{off}) )</td>
</tr>
<tr>
<td>B</td>
<td>(up+down)/2</td>
<td>( S_4 = 1/2(S(B,\text{up}) + S(B,\text{down})) )</td>
</tr>
</tbody>
</table>

A desired combination which eliminates the direct influence terms as well as the spin independent term is:

\[
1/4(T_1 - T_2 - T_3 + T_4) = E \cdot [1/2S(A,\text{off}) - 1/4S(A,\text{up}) - 1/4S(A,\text{down})] = E \cdot <S>
\]

The double modulation procedure therefore eliminates the effects of the spin flipper field acting directly on the detector and compensates for non-exact neutron spin reversals.

In order to control systematic errors which depend on the neutron spin direction, we also must develop a procedure which establishes that \( E \cdot <S> \) is non zero only as a result of \( T \)-violation. We consider two sources of systematic error: sequential \( T \)-even effects and misalignments. The first can be detected by their quadratic dependence on target thickness, and that the second can be identified by subsidiary measurements with the neutron spin aligned along different axes.

Polarized neutrons propagate through a medium according to a spin dependent index of refraction. If \( \chi_i \) is the initial spin state and \( \chi_f \) is the final spin state then

\[
\chi_f = \exp \left( i \left( n-1 \right) k z \right) \chi_i
\]
where
\[(n-1) = \frac{2\pi}{k^2} \rho f(0^\circ) = bs_x + cs_z + ds_y\]

Here \(n\) is the index of refraction, \(k\) is the neutron momentum, \(\rho\) is the number of scattering centers per unit volume, \(s_x, s_y\) and \(s_z\) are Pauli operators and \(b, c\) and \(d\) are complex coefficients which depend on the polarization or alignment of the target. The neutron momentum \(k\) is taken to be in the \(z\) direction. For the five fold correlation the alignment axis is in the \((z, x)\) plane at an angle of 45 degrees to the \(z\) axis and the polarizer prepares the spin to be either up or down (\(\pm y\)). For a purely aligned target the average value of \(I\) is zero, and the coefficients \(b\) and \(c\) are small, arising only from the weak interaction. It is the absence of a large spin-spin interaction which makes the FC less subject to systematic error than the TC.

The coefficients \(b, c\) and \(d\) which arise from the spin dependence of the forward elastic scattering amplitude are obtained from Table I. (Note that an invariant which involves \(I\) to an odd power averages to zero in a purely aligned target and induces no spin dependence in the neutron index of refraction. Also terms in \(k^3\) and higher will not contribute if only \(s\) and \(p\) wave neutrons are considered.)

\[
\begin{align*}
    b &= \sqrt{2} \sum_{N}^{\text{odd}} (sI) (I\cdot k)^N \equiv \sqrt{2} \, b_1 \\
    c &= \sum_{N}^{\text{even}} s \cdot k (I\cdot k)^N + \sqrt{2} \sum_{N}^{\text{odd}} (s \cdot I) (I\cdot k)^N = c_1 + \sqrt{2} \, b_1 \\
    d &= \sum_{N}^{\text{odd}} (s \cdot I \times k) (I\cdot k)^N
\end{align*}
\]

The FC is the leading term in \(d\) and the problem is to measure \(d\) in the presence of \(b\) and \(c\).

First consider the transmission through the polarizer and target. The operator which represents the combined action is

\[
U = \frac{1}{2} [(1 + g s_y) \exp \{iz(bs_x + cs_z + ds_y)\}],
\]

where \(g = +1\) for neutron polarization up and \(g = -1\) for polarization down. Let \(m^2 = b^2 + c^2 + d^2, P = \cos(mz),\) and \(Q = \sin(mz),\) and define \(T(g)\) as the transmission through the apparatus when the polarizer is in state \(g.\) Then

\[
T(+) - T(-) = \text{Im}(Pd^*Q^*/m^*) + (QQ^*/mm^*)\text{Im}(bc^*) = \text{Im}(Pd^*Q^*/m^*) + \sqrt{2}(QQ^*/mm^*)\text{Im}(b_1 c_1^*)
\]

The three interactions represented by \(b, c\) and \(d\) act coherently. The real parts precess the neutron spin and the imaginary parts attenuate the neutron flux. The first term depends linearly on the imaginary part of \(d.\) The flux is attenuated differently when the imaginary part of the five fold correlation is positive or negative. The second term represents a precession of the \(y\) component of the spin about the \(x\) axis by the real \((s \cdot I)(I \cdot k)\) interaction followed by an attenuation due to the imaginary \(s \cdot k\) interaction acting on the resulting \(z\) spin component. This term represents the sequential action of two \(T\)-even interactions.
resulting in a false T-odd interaction. In neutron transmission these terms can be removed by the addition of an analyzer after the target, or by comparing induced polarization with transmission asymmetry. As we now discuss, these effects can also be removed by less expensive techniques. For \( m_z \) small

\[
\text{Im}(Pd^*Q^*/m^*) \sim -z\text{Im}(d)
\]

and

\[
\sqrt{2} (QQ^*/mm^*)\text{Im}(b_1c_1^*) \sim \sqrt{2} z^2 \text{Im}(b_1c_1^*).
\]

The effect of the five fold correlation increases linearly with target thickness, while the sequential interaction increases quadratically with target thickness. This difference provides a direct means of testing whether or not a transmission asymmetry is the result of T-violation or of this sequential process. Since the coefficient \( c_1 \) may depend on whether or not the target is aligned, establishing that the longitudinal asymmetry is zero when the target is warm does not ensure the absence of this false simulation of T-odd effects.

In treating misalignments of \( s, k \) and \( I \), we may assume without loss of generality that \( k \) is in the \( z \) direction and \( I \) is in the \((z, x)\) plane. The effects of misalignments of the polarizer and the target are calculated by inserting a rotation operator between the polarizer operator and the target operator. Since we are interested in asymmetries which falsify time reversal, we set \( d = 0 \).

\[
U = (1+gs_y)(1+(1/2)(\theta s_x + \phi s_y + \eta s_z)) \exp[i\zeta(b s_x + c s_z)]
\]

The change in the transmission asymmetry is given by

\[
[T'(+) - T'(-)] - [T(+) - T(-)] = \theta \text{Im}(P^*Qc/m) + \eta \text{Im}(P^*Qb/m)
\]

This expression has a straightforward interpretation. The polarizer produces neutrons polarized in the \( y \) direction. A misalignment of the polarizer relative to the target corresponds to a rotation about the \( x \) axis (\( \theta \)) and gives a small polarization in the \( z \) direction. The imaginary part of \( c \) then produces a transmission asymmetry. The same argument applies to misalignment about the \( z \) axis which rotates \( y \) into \( x \) and induces a term proportional to \( \text{Im}(b) \). Misalignment with respect to the polarization axis \( y \) has no effect on the transmission asymmetry. The terms \( \text{Im}(P^*Qc/m) \) and \( \text{Im}(P^*Qb/m) \) can be measured directly by polarizing the beam in the \( x \) and \( z \) directions, respectively (\( \theta = \pi/2 \) and \( \eta = \pi/2 \)). Any small asymmetries which arise will be due to P-odd effects. These can now be calibrated for the situation where the neutron spin is along the \( y \) axis and will be suppressed by one to two orders of magnitude by the angle term.

**SUMMARY**

We have reviewed the scientific motivation and estimated sensitivities for parity and time reversal measurements with epithermal polarized neutron beams at LANSCE. The parity studies will provide a measurement of the weak interaction spreading width and test its relation to the effective weak N-N interaction in nuclei. P-even T violation will be tested at the level of \( 10^{-5} - 10^{-6} \) of the strong interaction, a factor of \( 10^2 - 10^3 \) better than the best present limits. Systematic errors associated with the time reversal study are shown to be tractable.

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Helicity dependence by the neutron radiative capture reaction of $^{139}\text{La}$ was measured by means of the capture $\gamma$-ray detection and the neutron transmission method. The asymmetry in the $\gamma$-ray detection is $9.5\pm0.3\%$ in the $p$-wave resonance of $^{139}\text{La}$ at 0.734 eV and the transmission asymmetry is $9.7\pm0.5\%$ in the same resonance.

Measurement of CP-violating direct amplitude in $K_L^0 \rightarrow \pi^0 e^+ e^-$ with the sensitivity of $1\times10^{-10}$ is being prepared.

An ultra cold neutron source by a method to cool the neutron by liquid helium II has been proposed to use for spallation neutrons from Japan Hadron Facility.

§1 Introduction

We are going to give a report on the experimental tests of T and P invariances, which have been carried out or are in preparation at KEK.

In the first half of the report, the production of polarized slow neutrons and the tests of P and T violations using these neutrons at KEK will be explained.

A project of production of high intensity ultra cold neutrons for the measurement of the electric dipole moment using superfluid helium will be mentioned.

Then, we will show the present status of the preparation of the measurement of CP-violating direct amplitude in $K_L^0 \rightarrow \pi^0 e^+ e^-$ decay.

§2 Slow Neutron Polarization by Longitudinally Polarized Proton Filter

The polarized slow neutron beam is quite important for studying the structure of condensed matter as well as the particles and nuclei. Especially the high intensity polarized neutron beam in the energy
region higher than 100 meV can open the new field of the neutron physics.

Slow neutrons transmitted through a polarized proton filter become polarized, because the cross section for the neutron-proton scattering in the singlet state is substantially larger than that in the triplet state. Recently, it was found that the neutron can be polarized in the direction of the beam as well as perpendicular to the beam direction.\(^1\) The results showed the possibilities to use a polarized proton filter as an effective polarizer or an analyzer of neutrons within the energy range between \(10^{-6}\) eV and \(10^7\) eV.

The method to polarize protons in the spin filter is similar to the one for the polarized proton target used in high energy physics experiments. Protons in the organic material containing \(\text{Cr}^V\) complexes can be polarized by means of the dynamic method in which material is irradiated by microwave of 70 GHz in a 2.5T magnetic field at a temperature lower than 0.5K. At KEK, ethylene glycol was used as filter material, which must be small blocks to get good thermal contact with liquid He, since the proton polarization depends strongly on the temperature. In order to satisfy the above requirement five layers of ethylene glycol plates were immersed in liquid \(^4\text{He},\) which was cooled by liquid \(^3\text{He}\) from the outside through a heat exchanger of copper fins.\(^2\)

A small NMR coil was embedded in the ethylene glycol in order to measure the proton polarization. The proton polarization was also determined from neutron transmission through the filter. The values of the proton polarization measured by the NMR and the neutron transmission are in good agreement within the errors. The polarization of protons was \(70\pm80\%\).

The neutron polarization \(P_n\) is determined by the formula

\[
P_n = \sqrt{1 - (T_o/T)^2},
\]

where \(T\) and \(T_o\) are the neutron transmissions when protons in the filter are polarized and unpolarized, respectively. It is a good approximation at energies above 200 meV. Typical values of the neutron polarization were \((72\pm2)\%\) and \((70\pm2)\%\) at \(0.7\) eV and \(3\) eV, respectively.
§3 Longitudinal Asymmetry in Neutron Radiative Capture Reaction of $^{139}$La

In the previous experiments at KEK$^3$ and at Dubna,$^4$ longitudinal asymmetries in neutron radiative capture reactions using polarized neutron beams have been found to be very large. The results show the large parity-nonconserving effect in nucleon-nucleus interactions, which may be explained by a parity mixing between p-wave and s-wave states of compound nuclei. Recently, several possibilities including CP-conserving and non-conserving nuclear interactions have been discussed to explain the large helicity asymmetry.$^5,6$ Sizable asymmetry on the p-wave resonance of $^{139}$La was explained by a parity mixing in the eigenstate of the S-matrix by Y. Yamaguchi. He also predicted the asymmetry in s-wave elastic scattering and s-wave radiative capture processes in the case of mixing of s-wave and p-wave. On the other hand, large asymmetry can be observed only in the p-wave resonance, if p-wave and d-wave resonances are mixed. In order to study the mechanism of the large enhancement of the parity violating effect at the p-wave resonance of $^{139}$La, simultaneous measurements of the capture $\gamma$-ray and the transmitted neutrons are carried out in a wide energy region including the p-wave resonance.$^7$ The neutron was polarized transversely by a polarized proton filter and was rotated to the longitudinal direction with an adiabatic passage method, using longitudinal field of 150-G solenoid, whose direction were changed every 2.5 sec. La targets of 2.5 cm in diameter were put in the neutron beam line at the position of 6.5 m from the neutron source. The target was cooled down to 40K~56K by a helium refrigerator in order to minimize the Doppler broadening of the resonance width. Capture $\gamma$-rays emitted from the target were detected by an annular BaF$_2$ scintillation counter. The threshold energy of $\gamma$-ray detection was about 1 MeV. Transmitted neutrons through the target were detected by a $^{10}$B-loaded scintillator. The asymmetry was found only in the p-wave resonance. The asymmetry in the continuum spectra in the energy region of 0.3~10 eV was less than $10^{-3}$. The asymmetry in the resonance peak ($c_{\gamma,p}$) was obtained after subtraction of the continuum spectra. The average value ($A_{L',\gamma}$) is (9.5±0.3)$\%$ in the resonance of 0.734 eV., which is consistent with the
previous result at KEK.  The neutron transmission asymmetry \( \varepsilon_n \) was observed for positive and negative helicity states. A longitudinal asymmetry in the total cross section for the p-wave resonance \( A_{L,n} \) is obtained from

\[
\varepsilon_n \text{ to be } 9.7\pm0.5\%, \text{ which is consistent with the result of the capture } \gamma\text{-ray measurement. However the value is inconsistent with the results at Dubna } (A_{L,n}=(7.3\pm0.5)\%). \]

The energy dependence of \( \varepsilon_n \) is roughly proportional to the Breit-Wigner resonance curve shown in Fig. 1. But small deviations from the resonance curve seem to exist. The small peaks seem also to be seen in the capture \( \gamma\text{-ray asymmetry. If these deviations are the effect of the d-wave resonance, the large asymmetry at 0.734 eV may be explained by the mixture of the p- and d-wave resonances. Further experimental study is desirable.}

§4 T-violation Test with Resonance Neutron Scattering

Possibility of using large enhancement of parity violating effects in neutron radiative capture reaction for the test of the time reversal symmetry violation has been searched for at KEK.

In the case of helicity asymmetry in the neutron radiative capture reaction on \(^{139}\text{La}\), the dynamic enhancement proportional to

\[
\langle p | H^{wk}_{p} | s \rangle/(E_p - E_s) \text{ is about } 10^3, \text{ where } E_p \text{ and } E_s \text{ are the p-wave and s-wave resonance energies, respectively, and } H^{wk}_{p} \text{ is the parity violating weak interaction Hamiltonian. In addition, the structural enhancement due to p-wave centrifugal barrier is about } 10^{2\lambda^3}, \text{ which is proportional to } \sqrt{\Gamma_s/\Gamma_p}, \text{ where } \Gamma_s \text{ and } \Gamma_p \text{ is the width of s- and p-wave resonances, respectively. Overall enhancement of } \sim 10^6 \text{ gives rise to the asymmetry of } \sim 10^{-1}, \text{ instead of typical parity violating asymmetries of } 10^{-7}. \text{ If there is the same resonance enhancement for the T-violation as for the P-violation, it will be hopeful to observe the T-violation effect in the neutron radiative capture reaction.}

The measurement of the asymmetry of the transmitted neutron in the reaction of transversely polarized neutron \((\vec{s})\) on the transversely polarized \(^{139}\text{La}\) nucleus \((\vec{l})\), whose spin direction is perpendicular to the neutron spin, is a simple example of the test of the T-violation. However, Stodolsky, Kabir, and Bunakov et al. pointed out the
following problem. In the first half of the target the real part of $\sigma \cdot \vec{I}$ interaction precesses about $\vec{I}$, so that in the second half the $\sigma \cdot \vec{p}$ interaction produces a transmission asymmetry. ($\vec{p}$ is the neutron momentum vector). Although both these interactions are time reversal invariant, overall effect looks as same as the time reversal violating interaction. They proposed several possible experiments of T-invariant test. One of them is the measurement of the difference between the probability of the helicity flip ($+ \rightarrow -$) and that of ($- \rightarrow +$) during the transmission in a transversely polarized target. It is a "macroscopic spin detailed balance". Another example is the measurement of the difference between the probability of the spin non-flip ($+ \rightarrow +$) of vertically polarized neutron during the transmission in the horizontally polarized (perpendicular to the beam) target and that of spin non-flip ($- \rightarrow -$) in the same target. It is as "backward run of the film". Both of them are possible experiments using a polarized $^{139}$La target.

Sizable polarization of $^{139}$La has been observed in the crystal of La$_2$Mg$_3$(NO$_3$)$_2$24H$_2$O with a little Nd$^{+++}$ by the dynamic method. However, the crystal is not useful for the La target, because the n-p scattering cross section is much larger than that of the n-La capture reaction at the p-wave resonance. The $^{139}$La in the crystal of La$_2$Mg$_3$(NO$_3$)$_2$24D$_2$O can also be polarized. But the background from other material than La is not negligible. Therefore, simpler crystal including La is now being tried to be polarized in a $^3$He-cryostat. The neutron spin rotation by the external field is minimized, as the polarization of $^{139}$La can be kept in lower field than that for the dynamic polarization, using a dilution refrigerator.

§5 UCN Source using Superfluid $^4$He for Spallation Neutron from Accelerator

The nature of CP-violation has been one of the most fundamental problems of elementary particle physics for more than 25 years. So far the CP-violating phenomena have been found only in kaon decay. On the other hand, the electric dipole moment (EDM) of the elementary particle would indicate both time reversal symmetry (T) violation and parity symmetry (P) violation, if it exists. It is believed that the neutron
is the most convenient particle to search for an EDM. The EDM of the neutron has been searched for after Purcell and Ramsey pointed out the importance of the measurement in 1950.\textsuperscript{9}

In the early stage of this kind of experiments, Ramsey et al. constructed a neutron beam magnetic resonance apparatus with a strong electric field applied in the region between the two separated oscillatory fields. Recent experiments on the neutron EDM involve ultra cold neutrons (UCN), whose velocity is less than 7 m/s and trapped by neutron total reflection at the surface of the trap. Two groups are engaged in the EDM searches with UCN using reactors. One of them is at Leningrad Institute of Nuclear Physics and the other is at ILL, Grenoble. Recent results show that the EDM of the neutron is $< 10^{-25}$ e·cm.\textsuperscript{10} In order to improve the values by 2 order of magnitude, it is necessary to produce higher density UCN.

Golub and Pendlebury proposed a method to cool the neutrons by super-fluid helium where produced phonons take away the heat of incoming cold neutrons.\textsuperscript{11} The advantage is that the density of the UCN increases with time, because the UCN are trapped and stored in a bottle. The method allows helium to be near the cold source. Installation of the helium converter near the cold neutron spallation source of a proton accelerator has an advantage over such a converter for neutrons from nuclear reactor.\textsuperscript{12} The source of spallation neutrons is small, and large solid angle of it can be covered by the helium converter. Such a UCN source was proposed as a facility of Japan Hadron Project (JHP). Thermal neutron flux of $7 \times 10^{12}$ cm$^{-2}$ s$^{-1}$ can be produced by primary proton beam of 1 GeV and 20 $\mu$A.\textsuperscript{13} It gives $5 \times 10^3$ cm$^{-3}$ s$^{-1}$ UCN which is 50 times stronger than that at Gatchina in Leningrad. I think, it might be also possible to get higher density by filling the bottle in which the UCN are stored only while the UCN flux is near its peak and then closing the bottle for the rest of the cycle. In order to test the feasibility of the method, a prototype of a bottle of helium II cooled by $^3$He cooling system was constructed at KEK (Fig. 2).\textsuperscript{14} Its length is about 3 m and the temperature of liquid reached 0.65K. The first test of slow neutron (12-15 Å) scattering in such low temperature
helium II using a cold neutron source at KEK-KENS facility started in May, 1989.

§ 6 Project of the Measurement of $K_{L}^{0} \to \pi^{0} e^{+} e^{-}$ Decay

The CP-violation occurs due to the complex phase in quark mixing angles in the standard model, which predicts the existence of the direct CP-violating amplitude ($\varepsilon'$) as well as the mass matrix term ($\varepsilon$). On the other hand, there is no direct amplitude in the super-weak model. Therefore, the measurement of $\varepsilon'$ is very important to check these models. The $K_{L}^{0} \to \pi^{0} e^{+} e^{-}$ decay is a CP-violating process if the $e^{+} e^{-}$ pair is produced through $Z^{0}$ or one photon. The standard model predicts that the contribution to this decay mode from $\varepsilon'$ is comparable or larger than that from the mass matrix term ($\varepsilon$). The branching ratio due to $\varepsilon'$ ($BR(\varepsilon'_{\pi^{0} e^{+} e^{-}})$) is expected to be $2 \times 10^{-11}$ from the values of Kobayashi-Maskawa angles, and the masses of $t$ and $c$ quarks. The $K_{L}^{0} \to \pi^{0} e^{+} e^{-}$ occurs also through CP-conserving two photon exchange process whose branching ratio is calculated to be $10^{-11} \sim 10^{-14}$.

The measurement of the branching ratio of $K_{L}^{0} \to \pi^{0} e^{+} e^{-}$ to the level of $10^{-10}$ was proposed at KEK.\(^{15}\) The setup of the experiment is shown in Fig. 3. It consists of a chamber system, an analyzing magnet, a gas Cherenkov counter, trigger/veto scintillators and an electro-magnetic calorimeter. As the electro-magnetic shower calorimeter, pure CsI scintillators (radiation length : 1.86 cm) will be used. They have fast and slow scintillation components. The fast component has a decay time of 10 nsec, while the slow one has that of about 1 nsec. The light output of the fast component is larger than that of the slow one. The energy resolution of $e^{+}$, $e^{-}$ and $\gamma$ at 1 GeV is about 2 %. The mass resolution of $\pi^{0}$ and $K^{0}$ of the detector system will be about 5-6 MeV.

The experiment will start in 1990.

The author is much indebted to Profs. Y. Masuda, N. Sasao, Y. Yamaguchi and H. Yoshiki for valuable discussions and giving him unpublished information.
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Fig. 1  The neutron transmission asymmetry ($\varepsilon_n$) around the p-wave resonance ($T=56$K). Solid line is the Breit-Wigner resonance curve.

Fig. 2  A bottle of helium II cooled by $^3$He cooling system for trapping and storing of UCN.
Fig. 3 Experimental set up of the measurement of $K^0_L \rightarrow \pi^+ e^+ e^-$. 
ON THE ELECTRIC DIPOLE MOMENT OF THE NEUTRON

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ABSTRACT

Experimental limits on the electric dipole moment of the neutron have placed stringent constraints on theories that attempt to explain CP violation. Preliminary results from an experiment at Grenoble on the electric dipole moment of the neutron will be presented.

INTRODUCTION

The origin of the $CP$ violation (where $C$ is charge conjugation and $P$ is parity), discovered in the $K^0$ system 25 years ago, remains a puzzle$^{1,2}$. Assuming that $CPT$ is conserved (where $T$ is time reversal symmetry), a $CP$ violating interaction will generate a permanent electric dipole moment (EDM) on elementary particles such as the neutron, as well as on non-degenerate collections of particles such as atoms$^3$. Historically, experimental limits on the EDM of the bare neutron have placed strong constraints on theoretical models that attempt to explain $CP$ violation (see ref.$^4$ for a review of the theory). In the standard model, the neutron EDM arises in second order of the weak interaction and is expected to be very small, of order $10^{-32}\,\text{e}\cdot\text{cm}$. Most extensions to the standard model allow the neutron EDM to be generated in first order of the weak interaction, making the expected value near $10^{-26}\,\text{e}\cdot\text{cm}$, close to the present experimental limit.

The electric dipole moment of an unpaired valence neutron in a nucleus can give the nucleus an EDM. Electromagnetic interactions between the nuclear EDM and atomic electrons can result in a permanent EDM on the atom. An atom can also acquire a permanent EDM from several sources independent of the neutron EDM$^5,6$. An intrinsic EDM on the electron can generate an atomic EDM in first order if the electron is unpaired, and through the hyperfine interaction if the electrons are paired. A $T$ violating interaction between nucleons and electrons can directly generate an atomic EDM. Finally, a $T$-violating nuclear force, for example from the QCD $\theta$ parameter, can give the nucleus and hence the atom an EDM. It is evident that EDM measurements on the bare neutron and atoms are complementary, and that both are necessary to isolate the source(s) of time reversal symmetry violation.

A new upper limit on the electric dipole moment of the neutron, $d_n$, has been measured recently by the Grenoble experiment. Preliminary results will be reported below. An upper limit on the EDM of the $^{199}\text{Hg}$ atom by the Seattle group$^7$ has been interpreted as having a sensitivity to $T$-violating interactions comparable to measurements on the bare neutron$^8,9$. Because the status of the $^{199}\text{Hg}$ experiment has been reported recently$^{10}$, no further mention will be made to it in the following discussion.
THE NEUTRON EDM EXPERIMENT

The Grenoble neutron EDM experiment is performed at the Institute Laue-Langevin High Flux Reactor. The measurement team includes: K.F. Smith, N. Crampin, J.M. Pendlebury, D.J. Richardson, and D. Shiers of the University of Sussex; K. Green, A.I. Kilvington, J. Moir, H.B. Prosper, and D. Thompson of the Rutherford-Appleton Laboratory; N.F. Ramsey of Harvard University; P. Ageron and W. Mampe of the Institut Laue-Langevin; A. Steyerl of the University of Rhode Island; and the authors. The experiment measures the nuclear magnetic resonance frequency of neutrons stored in a bottle. The neutron ground state, having spin $I = 1/2$, interacts with electromagnetic fields, $\vec{E}$ and $\vec{B}$, through the Hamiltonian:

$$H = -(d_n \vec{I} \cdot \vec{E} + \mu_n \vec{I} \cdot \vec{B})/I$$

where $d_n$ and $\mu_n$ are the electric and magnetic dipole moments of the neutron. We compare the neutron Larmor frequency (measured using Ramsey’s method of separated oscillatory fields) for parallel and for antiparallel magnetic and electric fields. It follows from Eq. (1) that the shift in Larmor frequency between the two field configurations is $\delta \omega_0 = -4d_n E/\hbar$; the minus sign is necessary because $\mu_n < 0$.

Ultra-cold neutrons (UCN) are neutrons with such a small energy that they are total internally reflected from the walls of a vessel (velocity $\leq 6$ m/sec)\textsuperscript{11}. Approximately 90 UCN/cm\textsuperscript{3} are provided by a neutron turbine\textsuperscript{12} and enter Figure 1, a schematic view of the experimental apparatus, from the left. The apparatus is similar to that used in an earlier neutron EDM search\textsuperscript{13} but benefits from an increase in UCN flux by two orders of magnitude. The UCN, polarized by transmission through a magnetically saturated 1 $\mu$m thick iron-cobalt foil, are guided to a 5 liter neutron bottle consisting of beryllium disk electrodes separated by a hollow 10 cm long BeO cylinder. A beryllium door in the ground potential electrode allows the bottle to be filled and emptied. The bottle is centered within five layers of mu-metal magnetic shields. A uniform 10 mGauss field is applied within the innermost magnetic shield. An electric field of 5 to 15 KV/cm is applied to the bottle by a high voltage connection to the beryllium disk opposite the neutron door. The electric field is reversed by reversal of the high voltage polarity.

After the bottle is filled with neutrons whose spins are along $\vec{B}$, and the bottle door is closed, a resonant $\approx 30$ Hz magnetic field is applied for 4 sec to turn the spins perpendicular to $\vec{B}$. The spins are allowed to precess freely for 70 sec (the escape time constant for neutrons in the bottle), before a second 4 sec pulse of 30 Hz field, coherent with the first, is applied. The resultant neutron spin state is analyzed by the iron-cobalt foil: after the door is opened, neutrons in the original spin state are transmitted through the foil and diverted to a detector. After a ten second wait, those neutrons remaining in the system are adiabatically spin flipped and counted. A minimum (maximum) in the original (flipped) spin state neutron count is found when the Larmor frequency is equal to the applied 30 Hz. The frequency of the oscillatory field is stepped to either side of the 7 mHz wide resonance on subsequent measurement cycles to determine the Larmor frequency. Typically, in
one measurement cycle, 20000 neutrons are counted for the two states combined. Including filling and emptying the bottle, each measurement cycle takes 124 s.

The electric field is governed by a 32 measurement cycle sequence: 8 cycles applied parallel to $\vec{B}$, 4 cycles off, 8 cycles antiparallel, and 4 cycles off, with 2 cycles taken to change each state. The leakage currents across the bottle is monitored and was less than 5 nA for most of the data presented here. The direction of the static magnetic field was reversed every two weeks.

![Neutron magnetic resonance apparatus](image)

**Figure 1.** Neutron magnetic resonance apparatus to measure the electric dipole moment of the neutron.

Magnetic field changes inside the mu-metal shield are monitored by three optically pumped rubidium magnetometers placed within 40 cm of the neutron bottle. The field at each magnetometer is averaged over the neutron storage time. A flux gate magnetometer monitors the magnetic field at the entrance to the hole where the high voltage cable enters the shields. A plot versus time of the Larmor frequency (in magnetic field units) as measured by the original spin state neutrons, the spin flipped neutrons and the field as measured by each of the three Rb magnetometers is shown in Figure 2 for an early two day data run. In Figure 2, the applied electric field follows the sequence given above. A neutron EDM would appear as a linear correlation between the resonant frequency of the neutrons and $E$, with no such correlation appearing on the magnetometer channels.

### DATA ANALYSIS

To determine the neutron resonance frequency, $\nu_n$, a first pass is made through
the data to extract the slope of the resonance curve. The slope corresponds to a neutron polarization of 0.8, and is found to be independent of the applied $E$. A second pass uses the slope and a combination of the counts for the two neutron spin states to yield a single frequency $\nu_n$ for each measurement cycle. Occasional jumps in the magnetic field occur due to movements of the reactor crane, relaxation stressing of the magnetic shields, or the switching of magnets on neighboring experiments.

A variety of digital filtering techniques, all in good agreement, have been used to extract the neutron frequency shifts that are linearly correlated with the applied electric field, $\Delta \nu_n = [\nu_n(E) - \nu_n(-E)]/2$, from the drifting background of the individual measurement sets (cycles around discrete jumps in the magnetic field are discarded). Identical analyses are performed for the three magnetometers yielding $\Delta \nu_{m1}$, $\Delta \nu_{m2}$, and $\Delta \nu_{m3}$. Frequency shifts quadratic in $E$ (voltage on/voltage off effects) and shifts between the zero electric field groups (due, for example, to the shields being magnetized by leakage currents or sparks) are also extracted.

The individual measurement set results can be combined into 15 groups, one

![Figure 2. A typical two day data run showing the neutron spin resonant frequency (converted to an equivalent magnetic field) for the original spin state neutrons (up) and spin flipped neutrons (down), and the field as measured by each of the three Rubidium magnetometers. An arbitrary offset has been added to separate the plots.](image)
for each six week reactor cycle of measurement time. For the neutron channel alone, before any corrections are applied, a weighted average of the 15 groups yields an uncorrected neutron EDM of \( \hat{d}_n = (-2.9 \pm 2.3) \times 10^{-26} \text{e} \cdot \text{cm} \) with a normalized \( \chi^2 = 2.9 \). The large value for \( \chi^2 \) indicates the presence of non-statistical (systematic) variations in the uncorrected EDMs.

The magnetometer readings have been used to correct for systematic errors in two independent ways. The first makes use of the observation that for perturbations external to the magnetic shields, the average value of the three magnetometers tracts the neutron resonance frequency to within 20%. For each measurement set, a corrected \( \Delta \nu_n \) is obtained by subtracting the average of \( \Delta \nu_{m1}, \Delta \nu_{m2} \) and \( \Delta \nu_{m3} \). The weighted average of the neutron results, after correcting by the averaged magnetometer results, yields \( d_{n1} = (-1.7 \pm 2.8) \times 10^{-26} \text{e} \cdot \text{cm} \) with normalized \( \chi^2 = 2.1 \). The reduction in \( \chi^2 \) for the corrected data comes primarily from the additional noise contributed by the magnetometers. Although this technique does not eliminate the non-statistical variations in the data, it reduces contributions to \( d_n \) from external magnetic field fluctuations.

The second magnetometer correction method allows for the possibility that systematic frequency shifts may be different for the neutrons and magnetometers (such as those due to magnetic fields from leakage currents across the storage bottle). For consecutive data sets between major alterations to the experimental apparatus (several reactor cycles), linear correlation coefficients between \( \Delta \nu_n \) and \( \Delta \nu_{m1}, \Delta \nu_{m2}, \) and \( \Delta \nu_{m3} \) were evaluated. A multiple linear regression routine\(^\text{14} \) was then used to fit \( \Delta \nu_n \) versus a term linear in the magnitude of \( E \) and linear terms in any \( \Delta \nu_{m3} \) if the \( mi \) correlation with \( \Delta \nu_n \) was statistically significant (greater than 80% probable). The weighted average of the four sets of data by this analysis yields a second corrected value for the neutron EDM, \( d_{n2} = (-3.3 \pm 3.5) \times 10^{-26} \text{e} \cdot \text{cm} \) with a normalized \( \chi^2 = 1.5 \). The larger error bar for \( d_{n2} \) compared with \( d_{n1} \) comes from the uncertainties generated by cross-correlations in the linear regression and from the use of a smaller data set (tighter cuts were used to determine which runs would be included in the correlation analysis).

For the data presented above, there was no frequency shift quadratic in the applied field. \( E^2: \Delta \nu_Q = |\nu(|E|) - \nu(0)|/E \) gives \( |d_Q| < 2 \times 10^{-26} \text{e} \cdot \text{cm} \), and no relative shift of the frequencies at \( E = 0: \Delta \nu_0 = |\nu(E = 0 \text{ after } +E) - \nu(E = 0 \text{ after } -E)|/E \) gives \( |d_0| < 4 \times 10^{-26} \text{e} \cdot \text{cm} \). The data analysis rejects measurement cycles on either side of a detected spark. When the cycles with sparks are included in the analysis, no significant changes in the results are found. Effects from small, undetected sparks are therefore expected to be negligible. The reversal of the direction of the static magnetic field has no statistically significant effect on the result.

Components of the neutron velocity, \( \vec{v}_n \), that circulate around the bottle circumference can generate a frequency shift proportional to \( \langle \vec{v}_n \times \vec{E} \rangle \). Such circulation may persist for times on the order of the bottle emptying time constant (3 sec). No evidence for such systematic effects was observed when the time interval
between closing the bottle door and applying the first Ramsey pulse was varied.

In summary, we see no evidence for systematic errors apart from the non-statistical variation in the neutron frequency shifts. The corrected neutron EDM values, \(d_{n1}\) and \(d_{n2}\) are in good agreement and differ little from the uncorrected value, \(d_n\). We conclude that at this stage, our preliminary result is:

\[
d_n = (-3 \pm 5) \times 10^{-26} \text{e} \cdot \text{cm}
\]

\[|d_n| < 13 \times 10^{-26} \text{e} \cdot \text{cm} \quad 90\% \text{ C.L.}
\]

The error bar comes from the regression analysis (times root \(\chi^2\)) rounded up to the nearest integer. The contribution to the above error from neutron counting statistics is \(1.9 \times 10^{-26} \text{e} \cdot \text{cm}\). Analysis is still in progress to try to determine the source of the non-statistical variation of the data. The above result can be compared to the upper limit on the neutron EDM found by the Gatchina experiment\(^{15}\): \(|d_n| < 26 \times 10^{-26} \text{e} \cdot \text{cm}\).

The present experiment is limited by how well the magnetic field within the neutron bottle is monitored by the spatially separated rubidium magnetometers. A larger neutron bottle that will contain polarized \(^{199}\text{Hg}\) as a magnetometer in coexistence with the neutrons is being developed to improve the experimental sensitivity.

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THE NEUTRON ELECTRIC DIPOLE MOMENT IN VARIOUS THEORIES OF CP VIOLATION

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ABSTRACT

There are many different theories of CP violation, all of which can fit the experimental information available from K decays. There have been many attempts to compute the neutron electric dipole moment in these theories, which have all concentrated on one particular mechanism for the generation of the neutron electric dipole moment. Here we consider both quark level and hadronic level processors, and argue that there is no double counting in simply adding these terms. We then survey the results obtained in the standard model, and in a number of non standard models.

INTRODUCTION - THE GENERATION OF THE NEUTRON ELECTRIC DIPOLE MOMENT

The simplest way to generate the electric dipole moment of the neutron is to simply generate the electric dipole moments of the quarks and to simply add them together. In principle, contributions of order $G_F^2$ are possible but in the standard model, these contributions are proportional to $|U_{q1}|^2$ and thus cannot introduce a CP violation phase. It is a remarkable result that in the standard model the simple order $G_F^2$ diagrams sum to zero\(^1\), and the lowest order contribution in the standard model is of order $G_F^2\alpha_s$.

The neutron electric dipole moment is then given in terms of the quark electric dipole moments by

$$D_n = \frac{1}{3}\left(4D_d - D_u\right)$$

A second mechanism which can generate an electric dipole moment of the neutron is the quark colour electric dipole moment. The replacement of the photon in any diagram contributing to the quark electric dipole moment by a gluon will transform a T and P violating interaction of the quark with the electric field, $-d_q\sigma\cdot E$, into a T and P violating interaction of the quark with the colour electric field, $-g_S\frac{\lambda^a}{2}\cdot\sigma\cdot E_a$, defining the quark colour electric dipole moment, $f^a_q$. This induces T and P violating components in the neutron wave function, and hence a neutron electric dipole moment which is given by

$$D_n = \frac{1}{3}e\left(\frac{4}{3}f_d + \frac{2}{3}f_u\right)$$

(2)
There are two further mechanisms which need to be considered. It is possible for exchange currents at the quark level to induce a neutron electric dipole moment. In practice the exchange current effect is reduced compared to the single quark electric dipole moment because it requires the two quarks to be close together. This effect thus only becomes potentially important in situations such as those of the standard model when the quark electric dipole moment contribution is already reduced.

![Meson loop graphs contributing to the neutron edm to first order in the weak interaction](image)

Figure 1: Meson loop graphs contributing to the neutron edm to first order in the weak interaction

One must also consider long range contribution to the neutron electric dipole moment. In a model in which these long range effects are represented through the coupling of mesons to baryons, meson loops (some of which are illustrated in figures 1 and 2) may contribute to the neutron electric dipole moment, as emphasised many years ago by Barton and White. Once again the qualitative difference between the standard model and other models is apparent - to obtain a non zero neutron electric dipole moment in the standard model one must go to contributions of second order in the weak interaction, i.e. of order $G_F^2$. Crewther et al. emphasised the logarithmic singularity in these meson loop diagrams in the chiral limit.
as the meson mass goes to zero. In practice these terms are singular but not divergent in this limit, behaving as $m^2 \ln m^2$.

![Meson loop graphs](image)

Figure 2: Meson loop graphs contributing to the neutron edm to second order in the weak interaction

With the information that one expects the neutron electric dipole moment in the standard model to be of order $G_F^2 \alpha_s$, whereas in other models it is expected to be of order $G_F$, one can make dimensional estimates of $D_n$:

$$D_n^{(SM)} \sim \frac{e}{2m_n} (G_F m^2) \eta_{+-} \sim 10^{-30} \text{ e cm} \quad (3)$$

$$D_n^{(other)} \sim \frac{e}{2m_n} (G_F m^2) \eta_{+-} \sim 10^{-24} \text{ e cm} \quad (4)$$

So we see that the present experimental limits (that $|D_n| < 2.6 \times 10^{-25}$ ecm with 90% confidence) are just on the threshold of an interesting stage in which they will be able to distinguish standard and non standard models of CP violation. It is also clear that it is now
important for theoretical calculations of the neutron electric dipole moment to be both reliable and believable if we are to achieve these goals.

We believe that the calculations reported in the literature do not meet these criteria because they have emphasised one of the possible mechanisms outlined above which contribute to the neutron electric dipole moment, without considering other possible contributions. It is important to ask whether these contributions are genuinely independent and additive. We believe they are, in that the quark electric dipole moment, quark colour electric dipole moment and quark exchange current effects are short distance effects, while the meson baryon loop are long distance effects. The separation between short distance and long distance effects is seen most clearly in the cloudy bag model, in which it is clear that both the proton baryon loops and the quark level terms will contribute in an additive way.

Now that we have some understanding of how the neutron electric dipole moment is generated by the underlying interactions we can discuss its value in particular theories.

**SUMMARY OF THEORETICAL RESULTS**

In a recent review, He, Pakvasa and I\(^5\) have made a systematic study of the electric dipole moment of the neutron electric dipole moment in various models of CP violation. We find that

(i) in the standard KM model with 3 families the neutron electric dipole moment is in the range \(1.4 \times 10^{-33} \leq |D_n| \leq 1.6 \times 10^{-31}\) e.cm,

(ii) the two Higgs doublet model has approximately the same value of \(D_n\) as the standard model,

(iii) \(D_n\) in the Weinberg model is predicted to satisfy \(|D_n| > 10^{-25}\) e.cm,

(iv) in a class of left-right symmetric models \(D_n\) is of the order of \(10^{-26}\pm 1\) e.cm,

(v) in supersymmetric models, \(D_n\) is of order \(10^{-22}\phi\) e.cm with \(\phi\) being the possible phase difference of the phases of gluino mass and the gluino-quark-squark mixing matrix,

(vi) the strong CP parameter \(\theta\) is found to be restricted by \(|\theta| < 10^{-9}\) using the present experimental limit on the neutron electric dipole moment.

I will now briefly discuss the first of these results briefly, referring to reference \(^5\) for further details, and for a discussion of cases (v) and (vi).

**THE STANDARD MODEL**

We have already seen that the neutron electric dipole moment arises in the standard model at the second order in the weak interaction. Indeed the quark electric dipole moment contributions are of order \(G_F^2 \alpha_s\), as emphasised by Shabalin\(^1\) and Khriplovich\(^6\), and are further suppressed because the quark electric dipole moment would vanish in a chirally symmetric world with massless quarks. The final result is that the quark dipole moments are of the order \(10^{-34}\) e.cm in the standard model. Quark level exchange currents also give contributions of this order.

Meson-baryon loops give significantly larger contributions. As we emphasised elsewhere\(^7\), when calculating CP violating effects at order \(G_F^2\) in a hadronic basis it is essential that the CP violating phase associated with each of the two interactions is included. If this is not done the results are not invariant under changes in the phase of the \(s\) quark wave function. We have included the octet baryons and mesons in the intermediate states,
calculating couplings using SU(3) where necessary. One also needs as an input the value of \( \frac{\text{Im} A_0}{\text{Re} A_0} \), which we calculate following the technique described in reference 8.

Subsequent to the publication of our letter we realised that there is also a contribution from neutral meson loops with the photon coupled to the baryon through a Pauli type of coupling. If we classify our results according to the type of photon coupling we obtain the contributions to \( D_n \) listed in table I.

<table>
<thead>
<tr>
<th>Photon coupling</th>
<th>Contribution to ( D_n ) in e cm</th>
</tr>
</thead>
<tbody>
<tr>
<td>meson and Dirac coupling to baryon</td>
<td>( 4.95 \times 10^{-29} s_2 s_3 s_8 )</td>
</tr>
<tr>
<td>Pauli coupling to charged baryon</td>
<td>( 1.95 \times 10^{-29} s_2 s_3 s_8 )</td>
</tr>
<tr>
<td>Pauli coupling to neutral baryons</td>
<td>( -2.7 \times 10^{-29} s_2 s_3 s_8 )</td>
</tr>
</tbody>
</table>

Note that the listed contribution from Pauli type coupling is that obtained on the assumption that the anomalous moment retains its full strength off mass shell. If this assumption is relaxed the contribution will reduce in magnitude. Noting this fact, and noting that \( s_2 s_3 s_8 \) is limited by

\[ 2 \times 10^{-4} \leq s_2 s_3 s_8 \leq 2 \times 10^{-3}, \quad (5) \]

we obtain our prediction for the magnitude of \( D_n \) in the standard model

\[ 1.4 \times 10^{-33} \text{ e cm} \leq |D_n| \leq 1.6 \times 10^{-31} \text{ e cm.} \quad (6) \]

It is worth remarking that \( D_n \) from the quark loops does not vanish in the chirally symmetric SU(3)\( \times \)SU(3) limit, as long as the limit \( m_\pi \rightarrow m_K \) is taken before the common meson mass is taken to zero, because in this limit the \( \ln \left( \frac{m_\pi^2}{m_K^2} \right) \) terms noted by Barton and White\(^2\) appear. However if the limit is taken through the meson masses tending independently to zero, then the electric dipole moment from the meson loops vanishes, just as does that from the quark electric dipole moments.

To complete our discussion of the standard model it is appropriate to comment on the calculations which include baryon resonances in the intermediate state, such as that of Gavela et al\(^9\). In reference 5 we have repeated their estimate, finding a contribution of the same order as that given above. However, we argued against including both terms on the grounds that baryon resonances are dual to the meson loop contributions (see 3), and that it is double
counting to include both. That each gives a contribution of the same order reinforces this argument, at least from my point of view.

Thus we see that the neutron electric dipole moment in the standard model with three generations is in the range of $1.6 \times 10^{-31}$ to $1.4 \times 10^{-33}$ ecm which is several orders of magnitude smaller than the present experimental upper bounds. With four generations, the calculated neutron EDM may be larger in magnitude. In this case there are six angles and three CP violating phases in the KM matrix. The CP violation in the $K^0 - \bar{K}^0$ system does not determine all the CP violating phases. Using allowed values from experiments, Barroso et al.\cite{10} calculated the neutron EDM with four generations for diagrams of Fig.(3.2) type and found that a factor of 20 enhancement with respect to the three generation model is possible. The same enhancement factor is also expected for the other diagrams. The experimental measurement of the neutron EDM at the level of $10^{-29}$ to $10^{-30}$ could be an indication of the presence of the fourth generation.

Finally, it should be remarked that the clear cut separation between the long distance and the short distance contributions to the neutron electric dipole moment is peculiar to the standard model, and is not repeated in most of other models we have considered.

**THE TWO HIGGS DOUBLET MODEL**

CP violation in the two Higgs doublet model\cite{11} can arise from different origins, CP violation from the phases in the KM matrix, or from spontaneous symmetry breaking. However, it can be shown that in the two Higgs doublet model, it is not possible to have spontaneous CP violation if one requires neutral flavor conservation at the tree level\cite{12}. We will therefore consider the case where CP violation arises only from the phases in $V_{KM}$. Consequently the calculation of the neutron electric dipole moment is similar to that in the standard model, with the new feature in the two Higgs doublet model that the charged physical Higgs gives additional contributions.

However it turns out that the new contributions from the Higgs graphs are significantly smaller (a few percent) than those in the standard model when the limits on couplings from an analysis of the experimental bound\cite{13} on $B \rightarrow K^* \gamma$, so that in practice the neutron electric dipole moment in the two Higgs doublet model may be taken to be the same as that in the standard model.

**THE WEINBERG MODEL OF CP VIOLATION**

A new factor arises when there is sufficient freedom in the Higgs sector to permit spontaneous CP violation. The minimal model of this type is the Weinberg model of three Higgs doublets\cite{14}. In this model the KM matrix is real and CP violation is due to complex vacuum expectation values of the Higgs fields. The relevant Lagrangian is

$$L_Y = 2^{3/4}G^{1/2} \sum \bar{U}[V_{KM}M_u \frac{1+\gamma}{2} X_i H_i^+] + M_d V_{KM} \frac{1+\gamma}{2} Y_i H_i^+]D + H.D$$

(7)

Here $X_i$ and $Y_i$ are related to the vacuum expectation values of the Higgs fields, $\text{Im} (X_i Y_i^*) = -\text{Im} (X_2 Y_2^*)$ and $H_{1,2}^+$ are the physical charged Higgs fields.
It is evident that this Lagrangian will induce a quark electric dipole moment at the one loop level through the mechanism of diagram similar to figure 4.1. The effective CP violating Lagrangian for the \( qqg \) and \( qq\gamma \) vertices are

\[
L_{qqg} = -\frac{G_F g_s}{16\sqrt{2}\pi^2} \sum_{k,L} \{ \text{Im}(X_L Y_L^*)V_{ik}V_{jk}\frac{df(x_k,\rho)}{dQ} \\
- \bar{q}_i m_j \sigma_{\mu \nu} e_\mu^a q_v (1 + \gamma_5) \frac{\lambda^a}{2} q_j \} + \text{H.C},
\]

\[
(8)
\]

\[
L_{qq\gamma} = -\frac{G_F e}{16\sqrt{2}\pi^2} \sum_{k,L} \{ \text{Im}(X_L Y_L^*)V_{ik}V_{jk}f(x_k,\rho)\bar{q}_i m_j \sigma_{\mu \nu} e_\mu^a q_v (1 + \gamma_5) q_j \} + \text{H.C},
\]

\[
(9)
\]

where \( L = 1,2 \) (the physical charged Higgs) and \( k \) sums over the internal quarks, \( x_k,\rho = (m_k/m_L)^2 \) and

\[
f(x) = \frac{x}{(1-x)^3} \left[ Q(3-4x+x^2+2 \ln x) - (1-x^2+2x \ln x) \right]
\]

At the one loop level it is clear that there will be quark electric dipole moments, quark colour electric dipole moments, and also CP violating \( K\Sigma\pi \) and \( \pi p n \) transitions. It is therefore important to examine all the possible contributions to the neutron electric dipole moment carefully. These contributions are displayed in Table 2.

Table 2. Contributions to the neutron electric dipole moment in the Weinberg Model

<table>
<thead>
<tr>
<th>Source</th>
<th>Contribution in ecm</th>
</tr>
</thead>
<tbody>
<tr>
<td>quark electric dipole moment</td>
<td>-1.9 ( \times 10^{-25} )</td>
</tr>
<tr>
<td>quark colour electric dipole moment</td>
<td>-0.32 ( \times 10^{-25} )</td>
</tr>
<tr>
<td>meson loop graphs with Dirac coupling</td>
<td>-11 ( \times 10^{-25} )</td>
</tr>
<tr>
<td>meson loop graphs with Pauli coupling</td>
<td>-0.44 ( \times 10^{-25} )</td>
</tr>
</tbody>
</table>

It is clear that the meson loop terms dominate. It is equally clear that the sum of the above terms exceeds the experimental upper bounds on \( D_n \). Furthermore, it has been pointed out\(^{[38,39]} \) that the exchange of neutral Higgs in this model violates CP due to the mixing between the real and imaginary parts of the neutral Higgs fields. The quark EDM from Fig.(5.1) has been evaluated in Ref.\(^{[38]} \) and found to be...
\[ D_q = \frac{Q_q e G_F}{\sqrt{2} \pi^2} m_q \sum \frac{X_i Y_i}{m_{H_i}^2} \left( \frac{m_H^2}{m_{H_i}^2} \right) \]  
\[ (11) \]

where \( Q_q \) is the electric charge of the q-quark, \( m_{H_i} \) the neutral Higgs masses, and \( X_i \) and \( Y_i \) are the mixings of the neutral Higgs particles. If one assumes that all the neutral Higgs have approximately the same masses and the average mixing \( <X'Y'> \) if of the same order as \( \text{Im}(X_1 Y_1^*) \), then this contribution to \( D_n \) is about \( 10^{-24} \text{ ecm} \). In Ref.\[39\], a different treatment of the neutral Higgs contribution to neutron EDM is carried out by evaluating the effective coupling for the scalar and pseudoscalar Higgs with nucleon and then using eq.(2.3.5). It is found that

\[ D_n = C_k n g_\sigma <X'Y'> m_N^2 \]  
\[ (12) \]

where \( C = 3.34 \times 10^{-3} \) from the loop integral, \( g_\sigma = (8/29)m_n(2G_F)^{1/2} \) is the scalar Higgs-nucleon coupling, and \( g_H = 2.5m_n(2G_F)^{1/2} \) is the pseudoscalar Higgs-nucleon coupling. If one assumes again that the mixing \( <X'Y'> \) is approximately the same as \( \text{Im}(X_1 Y_1^*) \), then \( D_n \) is about two to three orders of magnitude larger than the experimental bound. Since the exchange of the neutral Higgs particles conserves flavor, they do not contribute to the only observed CP violation in \( K^0-\bar{K^0} \) system, and thus no constraint can be put on the parameters. While it is possible that the assumption \( <X'Y'> \sim \text{Im}(X_1 Y_1^*) \) over estimates the neutral Higgs contribution and some cancelation may occur between the contributions from the neutral and the charged Higgses, it is unlikely that this will reduce the neutron EDM by an order of magnitude. If such happy cancellations do not occur, the Weinberg model for CP violation may be in trouble with the present experimental bound. If on the other hand the neutron EDM turns out to be of order \( \sim 10^{-25} \text{ ecm} \), the Weinberg model, with cancellations and an overestimate in eq.(5.14) could still be viable.

**THE LEFT RIGHT SYMMETRIC MODEL**

Left-right symmetric models\[16\] are based on the gauge group \( SU(2)_L \times SU(2)_R \times U(1)_{B-L} \). The gauge boson-quark interaction lagrangian is

\[ L_{W-F} = \frac{1}{2g_L^2} Q_L \gamma_\mu \overline{\tau} Q_L \overrightarrow{W}_L^\mu + \frac{1}{2g_R^2} Q_R \gamma_\mu \overline{\tau} Q_R \overrightarrow{W}_R^\mu + \frac{1}{2g^2} Q'_{\gamma_\mu} \overline{\tau} Q' B^\mu \]  
\[ (13) \]

where \( \overrightarrow{W}_L,R^\mu \) and \( B^\mu \) are the gauge bosons corresponding to the groups \( SU(2)_L,R \) and \( U(1)_{B-L} \) respectively; \( g_L,R \) and \( g' \) are the corresponding gauge coupling constants.

It is characteristic of the \( SU(2)_L \times SU(2)_R \) models that many additional parameters enter calculations of CP violation effects: these are the mixing angle, \( \xi \), of \( W_L \) and \( W_R \), and the phases \( \gamma, \delta_1 \) and \( \delta_2 \) which appear in the 2x2 mixing matrix for the right handed currents\[17\]. Explicitly the W mass eigenstates are
\[
\begin{pmatrix}
W_1 \\
W_2
\end{pmatrix} = 
\begin{pmatrix}
\cos \xi & \sin \xi \\
-\sin \xi & \cos \xi
\end{pmatrix}
\begin{pmatrix}
W_L \\
W_R
\end{pmatrix}
\]

and the right handed KM matrix is

\[
U_R = e^{i\gamma} \begin{pmatrix}
\cos \theta e^{-i\delta_2} & \sin \theta e^{-i\delta_1} \\
-\sin \theta e^{i\delta_1} & \cos \theta e^{i\delta_2}
\end{pmatrix}
\]

The phases in \( U_R \) give the principal contribution to \( K_1^0 - K_2^0 \) mixing, and \( \varepsilon'/\varepsilon \) and \( D_n \) both receive their leading contributions from \( W_L \cdot W_R \) mixing coupled with the phases in \( R_R \).

It has been calculated that\(^1\)

\[
\left| \frac{\varepsilon'}{\varepsilon} \right| = 276 \tan \xi \left[ \sin(\gamma-\delta_2) + \sin(\gamma-\delta_1) - 0.23 \left( \sin(\gamma+\delta_2) + \sin(\gamma+\delta_1) \right) \right]
\]

and

\[
|D_n| = \sin 2\xi \left\{ 4.5\sin(\gamma-\delta_2) + 74\sin(\gamma+\delta_1) - 1.1\sin(\gamma-\delta_1) + 15.8\sin(\gamma+\delta_2) \right\} \\
\times 10^{-23} \text{ e cm.}
\]

Unfortunately different combinations of phases enter the two CP violating observables, so it is not possible to use the CERN value for \( \varepsilon'/\varepsilon \) to constrain the dipole moment prediction without making further assumptions.

The simplest model of the phases is due to Chang\(^1\) using this model and using the 95\% confidence limits on \( \varepsilon'/\varepsilon \), viz

\[
5.5 \times 10^{-3} \geq \frac{\varepsilon'}{\varepsilon} \geq 1.1 \times 10^{-3}
\]

which gives the range of values

\[
1.9 \times 10^{-26} \text{ (ecm)} \geq |D_n^{\text{total}}| \geq 1.9 \times 10^{-27} \text{ (ecm)}
\]

There is the exceptional case\(^2\) when \( k'k \equiv 1 \) and \( 2\alpha \) lies near \( 3\pi/2 \); then \( \gamma \equiv \delta_2 \) and \( D_n \) can be as large as \( 10^{-25} \) ecm. Otherwise, this bound is respected.

The measurement of \( \varepsilon'/\varepsilon \) and \( D_n \) with slightly improved sensitivity will constrain pseudo-manifest left-right symmetric theories of CP violation tightly and can verify or rule out some models.

A variety of other models in the class have been constructed\(^3\) which can give rise to neutron electric dipole moments of the order of \( 10^{-25} \) ecm. These are reviewed in ref 5.

ACKNOWLEDGEMENTS

This paper describes work carried out in collaboration with Professor S Pakvasa, Professor S R Choudhury and Dr X-G He. It is a pleasure to express my gratitude for their contributions.
This work was supported in part by the Australian Research Council.

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ABSTRACT

The existing formalism used to describe spin observables in neutron transmission experiments is found to be inadequate. A suitable formalism is developed, through which time-reversal violating (and parity non-conserving) forward scattering amplitudes are identified, along with their corresponding spin observables. It is noted that new and more precise tests of $T$-symmetry are provided in transmission experiments and that such investigations are applicable more generally in nuclear and particle physics.

INTRODUCTION

Recent transmission experiments with slow neutrons have shown remarkable enhancements in two parity non-conserving (PNC) observables, the neutron spin rotation $A_z$ and the neutron longitudinal analyzing power, $A_z^L$. In the latter case, a value of $A_z$ close to 10% has been measured in the transmission of neutrons through $^{139}$La at the 0.734–eV resonance, and this nuclear enhancement of $10^5$ to $10^6$ is explained in terms of parity-mixed levels of $^{140}$La and p-wave barrier hindrance of the parity conserving transition.

The suggestion that nuclear effects could also provide enhancements in time-reversal violating (TRV) neutron transmission observables, which become accessible with polarized targets, has generated considerable interest and activity in the investigation of that possibility. Stodolsky and Kabir have developed a formalism to describe the spin aspects of neutron transmission, and they have suggested TRV observables to be measured. I have found some difficulties with their treatment, which I discuss in this paper. I then consider, again, the prospect for measurable TRV observables in transmission experiments.

DEVELOPMENT

Stodolsky and Kabir use the plane-wave neutron coherent amplitude in the target material,

$$M = e^{i(n-1)kd}, \text{ with } n-1 = \frac{2\pi}{k^2} \rho f(o),$$

$$f(o) = f_0 + f_x \hat{\sigma} \cdot \hat{s} + f_y \hat{\sigma} \cdot (k \times \hat{s}) + f_z \hat{\sigma} \cdot \hat{k},$$

where $n, k, d, \rho, f(o)$ are, respectively, the index of refraction, neutron wave number, distance of transmission in the target, density of target nuclei, and the neutron-nucleus forward elastic scattering matrix. It should be emphasized that only coherent non-spin-flip amplitudes contribute to the index of refraction, since scattering in which the target state is changed is
Incoherent.\(^5\) In (1a) \(\hat{\sigma}\) is the neutron spin operator and \(\hat{s}\) the target polarization, so \(f_z\) is a PNC term and \(f_y\) is both PNC and TRV\(^3\). With the coordinate system chosen (Fig. 1), the neutron direction \(\hat{k}\) and target polarization \(\hat{s}\) along the \(z\) and \(x\) axes, respectively, the forward scattering matrix becomes

\[
f(0) = f_0 + f_x \sigma_x + f_y \sigma_y + f_z \sigma_z.
\]

Then, taking

\[
C_j = \frac{4\pi}{k} \int f_j \rho d\,, \quad j = o, x, y, z
\]

\[
M = \exp \left[ i \frac{1}{2} \sum_j C_j \sigma_j \right].
\]

Since \(M\) is a two by two matrix in the neutron spin space, Stodolsky and Kabir take its most general form,

\[
M = \sum_j B_j \sigma_j, \quad j = o, x, y, z, \quad \sigma_0 = 1,
\]

and use (3) in the calculation of spin observables. For example, an analyzing-power component

\[
A_j = \frac{\text{Tr} \, M \, \sigma_j M^\dagger}{\text{Tr} \, M \, M^\dagger}.
\]

The transformation from the \(C_j\) of (2) to the \(B_j\) of (3) is not simple, and one loses the direct association and interpretation of the spin observables with the spin-dependent terms of the scattering matrix.

There are some clear difficulties with this procedure. By explicitly displaying the forward scattering matrix

\[
f(0) = \begin{pmatrix} f_0 + f_z & f_x - if_y \\ f_x + if_y & f_0 - f_z \end{pmatrix}
\]

it is seen that the TRV amplitude \(f_y\) appears in the off-diagonal neutron spin-flip terms. Since helicity is conserved in forward scattering, there must be a corresponding target-nucleus spin-flip. This, however, cannot contribute to the coherent neutron amplitude. Thus, in this treatment, a TRV observable is not available in coherent forward scattering.
The requirement of helicity conservation at $\theta = 0^\circ$ is even more restrictive when the forward scattering matrix is taken to be (1b). The equivalent condition is that $f(\omega)$ be invariant with respect to rotation about the z-axis, so

$$f(\omega) = f_0 + f_z \sigma_z$$

(6)

is the most general form. In this treatment, then, a TRV observable is excluded by conservation of angular momentum (z-component), whether the scattering is either coherent or incoherent.

Since (6) includes the PNC term $f_z$, it is of interest to calculate the PNC neutron coherent transmission observables in terms of $C_0$ and $C_z$ (thus $f_0$ and $f_z$) of eq. (2). Expressing the complex amplitude

$$C_j = \alpha_j + i\beta_j, \quad j = 0, z,$$

(7)
since $\sigma_0$ and $\sigma_z$ commute, (2) can be written as

$$M = R_0 D_0 R_z D_z$$

(8)

with

$$R_j = \exp \left[ \frac{i}{2} \alpha_j \sigma_j \right], \quad D_j = \exp \left[ -\frac{1}{2} \beta_j \sigma_j \right].$$

(9)

so $R_j (D_j)$ is a rotation (attenuation) matrix.

With $\sigma_j = \sigma_j^\dagger$,

$$R_j^\dagger = R_j^{-1}, \quad D_j^\dagger = D_j.$$

(10)

Also,

$$R_j = \cos \frac{\alpha_j}{2} + i \sin \frac{\alpha_j}{2} \sigma_j, \quad D_j = \cosh \frac{\beta_j}{2} - \sinh \frac{\beta_j}{2} \sigma_j.$$

(11)

Defining the transmission factor, through a target distance $d$,

$$\frac{I(d)}{I(0)} = T(d) = \frac{1}{2} Tr MM^\dagger,$$

(12a)

the analyzing power $A_z$ and the polarization transfer (rotation) coefficient $K_x^y$ are given by

$$TA_z = \frac{1}{2} Tr M \sigma_z M^\dagger$$

$$TK_x^y = \frac{1}{2} Tr M \sigma_x M^\dagger \sigma_y.$$

(12b)

Using (8) – (12),
\[ T = e^{i\beta_0} \cosh \beta_z \Rightarrow e^{i\beta_0} \]
\[ C_z \ll 1 \]
\[ A_z = -\tanh \beta_z \Rightarrow -\beta_z \]  
\[ (13) \]
\[ K_x^y = -K_y^x = \frac{-\sin \alpha_z}{\cosh \beta_z} \Rightarrow R_z(-\alpha_z). \]

Thus, one sees the known PNC differential attenuation of positive and negative helicity neutrons, i.e., \( A_z \), due to \( \beta_z \), the imaginary part of the amplitude \( C_z \), and the PNC neutron transverse spin rotation around the direction of propagation, due to \( \alpha_z \), the real part of \( C_z \). Since \( \beta_z \) is large (\( A_z \simeq 0.1 \)) at the 0.734-eV neutron-\( ^{139}\)La resonance, it would be most interesting to determine \( \alpha_z \) there via a measurement of the transverse spin rotation coefficient \( K_x^y \), assuming, of course, that the forward scattering is coherent at this energy (at \( 1-\text{eV} \lambda \simeq 3 \times 10^4 \text{ fm} \)).

Since target polarization is required in order to provide a TRV term in the forward scattering matrix, it is necessary for that matrix to encompass both the projectile and target spin-matrices. That is, an observable that involves only the projectile (target) polarization can be expressed in terms of the projectile (target) spin-matrix amplitudes alone, as in (13), but the combined spin-space amplitudes are required for an observable that involves both projectile and target polarizations. To investigate, then, the possibility of finding a TRV observable in transmission experiments, I consider in detail, as a prototype, the simplest case with spin-\( \frac{1}{2} \) projectile and target. There the 4x4 scattering matrix, including the same factor as in (2),
\[ F(o) = \frac{4\pi}{k} \rho d f(o) , \]  
(14)
can be constructed as the direct product of the 2x2 projectile and target spin matrices:
\[ F = F_1 \otimes F_2 , \quad \text{with} \]
\[ F_n = \sum_j C_{nj} \sigma_{nj} , \quad n = 1, 2 \]  
(15)
and \( n = 1 \) (2) labeling the projectile (target) matrix.

Then, a typical term in \( F \) is \( C_{1x} C_{2y} \sigma_{1x} \otimes \sigma_{2y} \), which I will abbreviate to \( C_{xy} \sigma_x \sigma_y \), so that the ordering of the subscripts provides the \( n = 1, 2 \) labeling. Recall from Fig. 1 that unit vectors along the coordinate axes are chosen to be
\[ \hat{z} = \hat{k} , \hat{x} = \hat{\sigma}_2 , \hat{y} = \hat{z} \times \hat{x} . \]  
(16)
Then, with \( \sigma_{1x} = \hat{\sigma}_1 \cdot \hat{x} \), etc., we have the following transformations under the \( P, T \) symmetry operations:
\[
\begin{array}{ccc}
\sigma_x & \sigma_y & \sigma_z \\
P & \sigma_x & -\sigma_y & -\sigma_z \\
T & \sigma_x & -\sigma_y & \sigma_z \\
\end{array}
\]  
(17)
The 16 $F$-matrix amplitudes,

$$C_{oo}, C_{ox}, C_{oy}, C_{oz}, C_{xo}, C_{xx}, C_{xy}, C_{xz}, C_{yo}, C_{yx}, C_{yz}, C_{zo}, C_{zx}, C_{zy}, C_{zz},$$

(18)
can then be classified according to their $P$ and/or $T$ symmetry. For example, the underlined in (18) are PNC amplitudes.

A particularly useful and more visual realization of these symmetries is obtained from the $F$-matrix itself,

$$F = \begin{pmatrix}
T & ++ & + & - & - \\
++ & F_{11} & F_{12} & F_{13} & F_{14} \\
+ & F_{21} & F_{22} & F_{23} & F_{24} \\
- & F_{31} & F_{32} & F_{33} & F_{34} \\
- & F_{41} & F_{42} & F_{43} & F_{44}
\end{pmatrix},$$

(19)

with the columns (rows) designated by the initial (final) helicities. Then, the amplitudes above the horizontal $P$-symmetry line are equal to the amplitudes below in inverse order,

$$(F_{11} \text{ to } F_{24}) = \pm (F_{44} \text{ to } F_{31}),$$

(20)

and, for elastic scattering, the amplitudes are similarly related across the diagonal $T$-symmetry line. Finally, imposing the condition of $\theta = 0^\circ$ helicity conservation, $F(\theta)$ becomes

$$F(\theta) = \begin{pmatrix}
F_{11} & 0 & 0 & 0 \\
0 & F_{22} & F_{23} & 0 \\
0 & F_{32} & F_{33} & 0 \\
0 & 0 & 0 & F_{44}
\end{pmatrix},$$

(21)

so that $R_z$-invariance alone reduces the number of amplitudes from 16 to 6. $P$-symmetry requires that

$$F_{11} = F_{44}, \ F_{22} = F_{33}, \ F_{23} = F_{32}$$

(22)

and the latter condition is required also by $T$-symmetry. Thus

$$F_{23} \neq F_{32}$$

(23)

is both PNC and TRV.
The corresponding requirements of \( R_z, P, \) and \( T \)-symmetries yield, from (15) and (17), the \( F(0) \)-matrix

\[
F(0) = C_{00} + C_{xx}(\hat{\sigma}_1 \cdot \hat{\sigma}_2) + C_{zz} \sigma_z \sigma_z ,
\]

so

\[
F(0) = C_{00} + C_{xx} \sigma_x \sigma_x + \sigma_y \sigma_y + C_{zz} \sigma_z \sigma_z ,
\]

This agrees with (22) in that only 3 independent amplitudes survive, and, specifically,

\[
F(0) = \begin{pmatrix}
C_{00} + C_{zz} & 0 & 0 & 0 \\
0 & C_{00} - C_{zz} & 2C_{xx} & 0 \\
0 & 2C_{xx} & C_{00} - C_{zz} & 0 \\
0 & 0 & 0 & C_{00} + C_{zz}
\end{pmatrix},
\]

in agreement with (21) and (22). Note, immediately that the \( C_{xx} \) term conserves helicity via a double spin-flip process, so it cannot contribute to the coherent neutron amplitude.

Now add the PNC terms from (18) that are \( R_z \)-invariant. These are

\[
C_{oz} \sigma_o \sigma_z + C_{zo} \sigma_z \sigma_o ,
\]

so

\[
F(0) = C_{00} + C_{oz} \sigma_o \sigma_z + C_{zo} \sigma_z \sigma_o + C_{zz} \sigma_z \sigma_z
\]

is the PNC coherent forward scattering matrix. Then proceeding analogously to (7) – (12), the result is

\[
A_{zo} = - \tanh \beta_{zo} \xrightarrow{C_{oz} \ll 1} -\beta_{zo} ,
\]

Thus, the statement immediately following (12), regarding the analyzing power, remains true with the simple substitutions \( A_{zo} , \beta_{zo} , C_{zo} \) for \( A_z , \beta_z , C_z \). Clearly, the target PNC analyzing power, i.e., using a polarized target and an unpolarized neutron beam, is given by

\[
A_{oz} = - \tanh \beta_{oz} \xrightarrow{C_{oz} \ll 1} -\beta_{oz} ,
\]

and this could be a very useful PNC observable in situations where a polarized target is more readily available than is a polarized beam.

Finally, the sixth \( R_z \)-invariant term to be considered is

\[
C_{xy}  (\hat{\sigma}_1 \times \hat{\sigma}_2)_z = C_{xy} (\sigma_x \sigma_y - \sigma_y \sigma_x) ,
\]

which, from (17), is seen to be both PNC and TRV. Thus, from (24), (26), and (29),

\[
F(0) = C_{00} + C_{oz} \sigma_o \sigma_z + C_{zo} \sigma_z \sigma_o + C_{xx} (\sigma_x \sigma_x + \sigma_y \sigma_y) + C_{zz} \sigma_z \sigma_z + C_{xy} (\sigma_x \sigma_y - \sigma_y \sigma_x)
\]
is the complete $R_x$-invariant PNC and TRV forward-scattering matrix. Again, the $C_{xy}$ term is a double spin-flip amplitude, so I conclude that in the $\frac{1}{2} + \frac{1}{2} \rightarrow \frac{1}{2} + \frac{1}{2}$ system, coherent forward scattering does not provide a TRV observable.

However, and perhaps more importantly, the term $C_{xy}$ in the forward scattering matrix suggests that a corresponding PNC, TRV observable is available in the more ordinary and widespread possibilities for incoherent transmission experiments in nuclear and particle physics. There one uses a treatment that features transmitted intensities rather than amplitudes; and the spin-dependent observables, i.e. the total cross-sections, are then related to the spin-dependent forward scattering amplitudes by the optical theorem.

The transmission factor (12a) is now

$$T(d) = \exp \left[ -\sigma_T p d \right] = \exp \left[ -\sigma \right] ,$$ (31)

where $\sigma_T$ is the unpolarized total cross-section, and thus $\sigma$ is a dimensionless "total cross-section" which includes the areal density factor $p d$. The corresponding spin-dependent cross sections are

$$\sigma_{jk} = \sigma(1 + p_j p_k A_{jk}) ; j, k = x, y, z ,$$ (32)

where $p_j (p_k)$ is the projectile (target) polarization along the $j(k)$ direction, and $A_{jk}$ is the corresponding (total cross-section) spin-correlation coefficient, which is essentially defined by this equation. Then, with $\sigma_{jk}(++) (\sigma_{jk}(+-))$ defined as the cross section for the pure spin states $p_j = p_k = 1 (p_j = -p_k = 1)$, we have

$$\sigma_{jk}(++) = \sigma_{jk}(-) = \sigma(1 + A_{jk}) ,$$
$$\sigma_{jk}(+-) = \sigma_{jk}(++) = \sigma(1 - A_{jk}) .$$ (33)

Using these spin dependent cross-sections, the corresponding transmission factors (31) are defined as

$$T_{jk} = \frac{1}{2} \left[ \exp \left[ -\sigma_{jk}(++) \right] + \exp \left[ -\sigma_{jk}(+-) \right] \right] $$ (34a)

and

$$\Delta T_{jk} = \frac{\exp \left[ -\sigma_{jk}(++) \right] - \exp \left[ -\sigma_{jk}(+-) \right]}{\exp \left[ -\sigma_{jk}(++) \right] + \exp \left[ -\sigma_{jk}(+-) \right]} $$ (34b)

Thus, $T_{jk}$ is the transmission factor for a completely polarized beam transmitted through an unpolarized target (and vice versa), while $\Delta T_{jk}$ is the transmission asymmetry of the polarized beam for opposite states of the target polarization. Using (33),

$$T_{jk} = e^{-\sigma} \cosh \sigma A_{jk} ,$$ (35a)
$$\Delta T_{jk} = - \tanh \sigma A_{jk} .$$ (35b)

We now use the spin-dependent form of the optical theorem\(^6\) to express $\sigma A_{jk}$ in terms of the imaginary part of the corresponding forward scattering amplitude. That is
\[
\sigma_T(\rho) = \frac{4\pi}{k} \text{Im} \ Tr [\rho f(\omega)],
\]  

where \(\rho_{jk}\) is the density matrix representing the initial polarizations and \(\sigma_T(p_j, p_k)\) is the corresponding total cross-section. The normalization \(Tr \rho = 1\) has been chosen. Then, in the established notation,

\[
\rho_{jk} = \frac{1}{4} (\sigma_o + p_j \sigma_j) \otimes (\sigma_o + p_k \sigma_k)
\]

so

\[
\rho_{jk} (+\pm) = \frac{1}{4} (1 + \sigma_j \sigma_o \pm \sigma_o \sigma_k \pm \sigma_j \sigma_k).
\]

With (14), (36) becomes

\[
\sigma_{jk} = \text{Im} \ Tr [\rho_{jk} F(\omega)],
\]

and with

\[
\sigma \Delta_{jk} = \frac{1}{2} [\sigma_{jk}(++) - \sigma_{jk}(--)],
\]

we have

\[
\sigma \Delta_{jk} = \frac{1}{4} \text{Im} \ Tr [(\sigma_o \sigma_k + \sigma_j \sigma_k) F(\omega)].
\]

Then noting that

\[
Tr (\sigma_j \sigma_k) (\sigma_j' \sigma_k') = 4 \delta_{jj'} \delta_{kk'}
\]

the only terms in \(F(\omega), (30)\), that survive in (41) are the terms \(C_{0k}\) and \(C_{jk}\), once the polarizations \(p_j, p_k\) have been selected. For our purpose here, to identify a TRV observable corresponding to the amplitude \(C_{xy}\), the appropriate choice is \(p_j = p_x, p_k = p_y\) (or vice versa), for which

\[
\sigma \Delta_{jk} = \text{Im} C_{xy} = \beta_{xy}.
\]

Finally, then, the corresponding transmission asymmetry,

\[
\Delta T_{xy} = -\tanh \beta_{xy} C_{xy} < 1 \xrightarrow{C_{xy} 

\text{is the observable to measure, with a vector (rank 1) polarized target, in the search for a combined PNC, TRV effect. This is also the simplest experimentally, for which high accuracy can be achieved.}

Since it is clear from the foregoing development that there is no purely TRV forward amplitude in the \(\frac{1}{2} + \frac{1}{2} \rightarrow \frac{1}{2} + \frac{1}{2}\) spin system, it is important to examine the suggestion\(^7\) for an additional \(T\)-odd, \(P\)-even term [presumably in \(f(\omega), (1a)\)] of the form \((\hat{\sigma} \cdot \hat{K} \times \hat{S}) (\hat{K} \cdot \hat{S})\), with target spin \(S \geq 1\) since \(S^2\) represents an alignment. It is clear that this term would not contribute to \(f_2\) in (1b), so it is excluded by \(R_z\)-invariance. It can, however, be provided by a term in the \(F\)-matrix (15) for a system with the spin structure \(\frac{1}{2} + 1 \rightarrow \frac{1}{2} + 1\). Since space limitations preclude a discussion here in the detail of that for the \(\frac{1}{2} + \frac{1}{2} \rightarrow \frac{1}{2} + \frac{1}{2}\) system, I will simply indicate the result, and the details will be available elsewhere.\(^8\) In (15), \(F_2\) now becomes a \(3 \times 3\) spin-1 matrix, with
\[ F_2 = \sum_j C_j P_j + \sum_{k l} C_{k l} P_{k l}, j = o, x, y, z; k, l = x, y, z, \quad (45) \]

with the second sum limited to five independent terms. The \( P_j (P_{k l}), j \neq 0, \) are the vector, rank 1 (tensor, rank 2) components of the spin-1 matrix-operator. Thus, the 9-term \( F_2 \) combined with the 4-term \( F_1 \) provides the required 36 terms of the \( F \)-matrix. The forward scattering matrix then is

\[
F(o) = \begin{pmatrix}
\frac{3}{2} & \frac{1}{2} & \frac{1}{2} & \frac{1}{2} & \frac{1}{2} & \frac{3}{2} \\
\frac{3}{2} & F_1 & 0 & 0 & 0 & 0 \\
\frac{1}{2} & 0 & F_2 & 0 & F_3 & 0 \\
\frac{1}{2} & 0 & 0 & F_4 & 0 & F_5 \\
\frac{1}{2} & 0 & F_5 & 0 & F_4 & 0 \\
\frac{1}{2} & 0 & 0 & F_3 & 0 & F_2 \\
\frac{1}{2} & 0 & 0 & 0 & 0 & F_1 
\end{pmatrix}, \quad (46)
\]

with the columns (rows) now labeled with the initial (final) channel helicities. \( P \)-symmetry has been invoked, so there remain only 5 independent amplitudes. From inspection, \( T \)-symmetry requires \( F_3 = F_5 \), so \( F_3 \neq F_5 \) will supply the \( TRV \) amplitude, and this condition is provided via the term

\[ C_{y,xz} (\sigma_x P_{yx} - \sigma_y P_{zx}) , \quad (47) \]

in a straightforward extension of the previous notation. Thus, \( C_{y,xz} \) is the \( TRV \) amplitude and the corresponding \( TRV \) observable is the transmission asymmetry

\[ \Delta T_{y,xz} = \tanh \beta_{y,xz} \frac{C_{y,xz} \ll 1}{\beta_{y,xz}} , \quad (48) \]

As the notation indicates, this corresponds to neutron polarization \( p_y \) in combination with the target tensor polarization \( p_{zx} \), i.e. alignment along the direction \( z = x \).

**SUMMARY**

It has been established that the complete spin-space matrix must be used in the description of \( TRV \) effects in transmission experiments, and it is seen that transmission observables can provide new and sensitive tests of \( T \)-symmetry. An especially important result is that the \( TRV/PNC \) observable is given directly by the imaginary part of the corresponding \( TRV/PNC \) forward scattering amplitude, in contrast to the situation that exists for the standard (non-forward) scattering experiment. There, \( T \)-symmetry tests can be accomplished only via comparison of two separate observables (e.g. \( P_y = A_y \)), and, as has been established, no single null test is possible. Now, \( \Delta T_{y,xz} \) (48) provides just such a null test of \( T \)-symmetry, and such an
experiment is capable of attaining an unprecedented precision, comparable to that reached in $P$-symmetry tests, where, for example, $A_z$ values of the order $10^{-7}$ have been measured.\(^{11}\)

Finally, although the motivation for the development of this spin formalism for transmission experiments was derived from the exciting results achieved with slow neutrons, this development clearly has application at higher energies in nuclear and particle physics with respect to attaining considerably higher precision in tests of $T$-symmetry.

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CHARGE SYMMETRY BREAKING IN np ELASTIC SCATTERING AT 183 MeV


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ABSTRACT

We report first results of an experiment recently completed at IUCF which is sensitive to charge-symmetry-breaking (CSB) effects in the np interaction. Specifically, we have measured the spin-dependent left-right asymmetries for np elastic scattering over a broad (60° c.m.) angular range by scattering polarized neutrons from polarized protons. Charge symmetry requires the neutron and proton analyzing powers $A_n(\theta)$ and $A_p(\theta)$ to be equal. Our preliminary result for their difference $\Delta A(\theta) = A_n(\theta) - A_p(\theta)$, averaged over the angular range $82.2° < \theta_{cm} < 116.1°$, is $(32.1 \pm 6.1 \pm 6) \times 10^{-4}$. The listed uncertainties are statistical and systematic, respectively. We also discuss aspects of the angular dependence of the CSB effects. Our results compare favorably with meson-exchange calculations which include the effect of the n-p mass difference on one-boson exchange and $\rho - \omega$ mixing.

INTRODUCTION

Isospin invariance of the strong interaction has been the subject of many experimental and theoretical investigations over the years. The interpretation of experimental tests has often been plagued by difficulties in properly subtracting effects with a purely electromagnetic origin. It is now generally accepted that the (Coulomb-corrected) spin-singlet nucleon-nucleon (NN) scattering lengths are not all equal, i.e., $a_{nn} \neq a_{np}$, indicating a clear violation of charge independence.\(^1\)\(^-\)\(^3\) However, until recently the evidence concerning charge symmetry (the weaker form of isospin invariance) has been inconclusive. A decade ago, a new generation of experiments was initiated to address the issue of charge symmetry violation definitively.\(^4\)\(^,\)\(^5\)

We report here the first results of an np elastic scattering measurement recently completed at IUCF, which is sensitive to charge-symmetry-breaking (CSB) effects in the strong interaction.

Formally, charge independence (CI) requires that the nuclear Hamiltonian be invariant under arbitrary rotations in isospin space. Hence in a given space-spin state, the strong potential must be the same for isorotplet substates (e.g., the three spin-singlet NN scattering lengths should all be equal after correction for electromagnetic contributions). Charge symmetry (CS), on the other hand, requires the invariance of the nuclear Hamiltonian with respect to a specific rotation, by 180° about the $T_2$ axis, which reverses the sign of the third component of isospin ($T_3 \rightarrow -T_3$). CS can remain valid even if CI is violated. Most previous tests of charge symmetry in the two-nucleon system have involved the search for differences between nn and pp potentials. In the classification scheme of Henley and Miller,\(^6\) experiments of this type (e.g., comparison of $a_{nn}$ to $a_{pp}$, or of $^3$H to $^3$He binding energies) are sensitive to class III CSB potentials. Such potentials have isovector terms symmetric under interchange of the two nucleons (i.e., $V_{III} \propto [r_3(1) + r_3(2)]$, and hence have no effect on the np system. However, charge symmetry has implications for the np system as well. The potentials that
give rise to CSB in the np system, referred to as class IV potentials, are antisymmetric under interchange: \( V_{IV} \propto [r_3(1) - r_3(2)]|\vec{\sigma}(1) - \vec{\sigma}(2)| \cdot \vec{L} \) or \([r_3(1) \times r_3(2)]_3|\vec{\sigma}(1) \times \vec{\sigma}(2)| \cdot \vec{L}\).

The experimental test we are about to describe involves spin observables in np elastic scattering. We have measured spin-dependent left-right asymmetries over a broad angular range for the elastic scattering of polarized neutrons by polarized protons. The measured quantities of interest here are \(A_n\), the analyzing power obtained by observing the asymmetry that reverses sign when the neutron (beam) spin is flipped, and \(A_p\), the corresponding analyzing power associated with the proton (target) spin. If charge symmetry (invariance under interchange of neutrons and protons) holds, these two analyzing powers must be equal (i.e., \(\Delta A \equiv A_n - A_p = 0\)) at all angles. Violations can arise through class IV potential terms, via the isospin mixing of \( T=0 \) and \( T=1 \) np scattering states with the same spin and parity: \(^1P_1\) and \(^3P_1\), \(^3D_2\) and \(^1D_2\), etc. (note that there can be no analogous mixing for S-waves).

The only previous experimental test of CSB due to class IV potentials was the TRIUMF measurement\(^7\) of \(\Delta A\) in np scattering at a laboratory bombarding energy of 477 MeV. They measured separately the zero-crossing angles for \(A_n(0)\) and \(A_p(0)\) and found them to differ in the c.m. frame by \(0.340^\circ \pm 0.162^\circ \pm 0.058^\circ\) (where the second error is systematic), which implies \(\Delta A = (47 \pm 22 \pm 8) \times 10^{-4}\) at the location of the zero-crossing (\(\theta \sigma \approx 71^\circ\) c.m.). We will hear about future CSB tests planned by that group later this morning. The experiment that we describe here, in addition to being an independent test of charge symmetry in the np system carried out at a different bombarding energy (183 MeV), has achieved a significantly greater statistical precision, and covers a much broader angular range than the TRIUMF data. Furthermore, the IUCF and TRIUMF measurements have different sensitivities to various CSB mechanisms, as will be discussed below. We mention in passing that a byproduct of our CSB experiment is an extremely high statistics measurement\(^8\) of the spin correlation parameter \(C_{NN}(\theta)\).

In the remainder of this presentation I will discuss: theoretical expectations for \(\Delta A(\theta)\), equipment and experimental procedures for the IUCF experiment, the analysis and results, and systematic errors. However, I would first like to make a few remarks to put this talk into perspective. The test of CSB in np scattering, while concentrating on a system that is particularly amenable to theoretical interpretation, without serious ambiguities arising from "trivial" electromagnetic effects, has required long and arduous experimental efforts. After many years of preparation, the actual IUCF production data acquisition required about 100 days of beam time, spread over a time period of just over one year, and was completed in September of 1988. In the ten months since, we have completed three rounds of replay of all the data, each requiring roughly a month of CPU-time on a VAX 8650, with progressively more sophisticated software. The last replay was completed only two weeks ago. In this sense our results must be termed preliminary. However, we feel that it is unlikely that any of our major conclusions will be altered substantially by our planned final round of replay, for which the major aims are to incorporate various further refinements in the analysis and to pin down certain systematic errors more quantitatively.

THEORETICAL EXPECTATIONS

Isospin mixing in the np system can be caused by both long- and short-range interactions. In the former category, there is a small direct electromagnetic (one-photon exchange) contribution that arises from the interaction of the neutron magnetic moment with the proton current. Violations in the short-range interaction are thought\(^9,10\) to arise fundamentally from the inferred mass difference between up and down quarks and from photon exchange among quarks. However, most current theoretical treatments of CSB in nuclear systems do
not attempt to use quark degrees of freedom directly; rather, they incorporate the latter effects indirectly, in a meson-exchange picture, by making use of measured manifestations (the n-p mass difference, matrix elements for isospin mixing of neutral mesons) of the quark-level CSB. Within this framework, many groups\(^{11-19}\) have calculated \(\Delta A(\theta)\) for np scattering at a variety of energies.

For np scattering at bombarding energies on the order of a few hundred MeV, the most important CSB mechanisms are: i) the electromagnetic spin-orbit coupling already mentioned; ii) the effect of the n-p mass difference on one-pion (and to a lesser extent, single-\(\rho\) exchange; and iii) isospin mixing of the \(\rho^0\) (\(1^{-},T=1\)) and \(\omega^0\) (\(1^{-},T=0\)) mesons. Other contributions involving multiple pion exchange or simultaneous exchange of mesons and photons are difficult to calculate\(^{6,20}\) but thought to be small.\(^{13,21}\) The n-p mass difference effects contribute to both class III and IV potentials described above. \(\rho - \omega\) mixing (with the mixing matrix element deduced from data\(^{22}\)) on the G-parity violating decay \(e^+e^- \rightarrow \omega^0 \rightarrow \pi^+\pi^-\) can also contribute to both classes of potentials, while a similar process involving pseudoscalar \(\eta - \pi\) meson mixing contributes only to class III, because of its spin structure.

\[
\Delta A \equiv \frac{A_n - A_p}{2} \div 100
\]

\[
\Delta A \quad \Delta A_{\rho} \quad \Delta A_{\omega} \quad \Delta A_{\gamma} \quad T_n = 183 \text{ MeV}
\]

\(\theta_{\text{cm}}\)

Fig. 1. CSB contributions to \(A_n(\theta) - A_p(\theta)\) from various meson-exchange processes, as calculated by Holzenkamp, Holinde, and Thomas\(^{14}\) employing the Bonn NN potential. The electromagnetic spin-orbit contribution arising from one-photon exchange (\(\Delta A_{\gamma}\)) is also indicated. The 183 MeV "SP88" phase shifts of Ref. 26 were used for the charge-symmetry-conserving amplitudes, and to calculate the average analyzing power \(A(\theta)\) also shown.

Typical calculations for various CSB contributions to \(\Delta A(\theta)\) are shown in Fig. 1. For the meson-exchange components, the calculations are based on those of Holzenkamp, Holinde, and Thomas\(^{14}\) who use the Bonn NN potential\(^{23}\) to describe the perturbing interaction and the distorted waves needed in evaluating the CSB scattering amplitude. We have scaled the \(\rho - \omega\) contribution to this amplitude to account for the most recent value\(^{24}\) of the \(\rho - \omega\) mixing matrix element \((-4520 \text{ MeV}^2 \text{ vs. } -3400 \text{ MeV}^2\) used in ref. 14). The curves in Fig. 1 have
been generated with the program SCORE\textsuperscript{25} by combining the theoretical CSB amplitude with isospin-conserving amplitudes based on Arndt (Spring 1988) np phase shifts.\textsuperscript{26}

There are several things to note about the predictions in Fig. 1. For any one of the individual CSB effects, $\Delta A$ is on the order of a few parts per thousand, which sets the level of accuracy needed in both the IUCF and TRIUMF measurements. Over the angular range of interest for the present experiment ($60^\circ \leq \theta \leq 120^\circ$), the nuclear $\Delta A$ terms are comparable to or larger in magnitude than the electromagnetic contribution. This is in sharp contrast to conventional class III CSB investigations, where Coulomb effects are often an order of magnitude larger than the strong-interaction signal one seeks. The $\Delta A(\theta)$ distributions are in general quite different in shape from that shown for the average np analyzing power $A(\theta)$. As we shall discuss later, this difference is crucial for a measurement of $\Delta A$ to succeed. In particular, in these calculations the $\rho - \omega$ mixing contribution peaks near the $A(\theta)$ zero-crossing angle $\theta_0$. At the energy of the TRIUMF measurement (477 MeV) the situation is quite different. There, calculations similar to those shown in Fig. 1 suggest that the $\rho - \omega$ term passes through zero very near $\theta_0$; hence, it is unlikely that a measurement at 477 MeV can be very sensitive to the $\rho - \omega$ contribution (for reasons explained later). Rather, such a measurement primarily probes the effect of the np mass difference.

The general features of the calculations noted above, namely, the order of magnitude of the nuclear effects and the shape of $\Delta A(\theta)$ for each contribution, apply to essentially all meson-exchange calculations that have been done. The use of different NN potentials can introduce small changes in $\Delta A(\theta)$ via the effects of distortions.\textsuperscript{18} However, the dominant ambiguity in the predictions is in the magnitude of the $\rho - \omega$ mixing contribution, stemming from uncertainties in the $\rho$-NN and $\omega$-NN coupling constants. The magnitude of this term varies by roughly a factor of three among different calculations, with the curve in Fig. 1 characteristic of the larger end of the range.

**EXPERIMENTAL EQUIPMENT AND MEASUREMENTS**

Our measurements were carried out at the Indiana University Cyclotron Facility (IUCF), using a specially constructed polarized neutron facility (PNF), polarized proton target (PPT), and particle-coincidence detection system which we now briefly describe.

A 200 MeV vertically polarized proton beam from the IUCF cyclotrons, incident from the left as depicted in Fig. 2, was bent downward by 10° and passed through a 20 cm long liquid deuterium (LD\textsubscript{2}) production target. Neutrons from the charge exchange reaction $^2\text{H}(p,n)^2\text{H}$ at $\theta_{lab} = 10^\circ$ passed through a dipole “sweep” magnet and were then collimated into a beam 5 cm wide $\times$ 7 cm high (at the PPT location) by a lead and steel channel embedded in a $\sim 3$ m thick steel and heavy concrete shielding wall. The resulting neutron beam energy distribution peaks at approximately 181 MeV with a FWHM of $\sim 15$ MeV. The neutron flux with 50 nA of incident primary proton beam was $5 \times 10^4$/cm$^2$ - sec. The dominant vertical component of the secondary neutron beam polarization (with a typical magnitude of 0.58) was reversed at regular intervals (30 sec.) by switching RF transitions at the polarized ion source. The sweep magnet served as a pre-collimator for neutrons and removed forward-going charged particles from the secondary beam. It also served the important function of precessing horizontal neutron spin components by $\pm 90^\circ$, allowing them to be effectively cancelled by averaging data taken with opposite sweep magnet polarities.

The primary proton beam polarization was monitored by polarimeters in both low- and high-energy cyclotron beam lines. The latter provided continuous monitoring of both vertical and sideways polarization components. The flux and polarization profile of the neutron beam was monitored continuously during acquisition by a high-energy neutron polarimeter located
Fig. 2. Side view of the IUCF polarized neutron facility (PNF) and CSB detector arrays. Polarized protons are bent 10° downwards and, after passing through a liquid deuterium production target, are focused into a Faraday cup below floor level. Feedback loops as indicated hold the beam position steady at the target and dump locations. Neutrons from the production target are collimated and may subsequently scatter in the polarized proton target (PPT), CSB flux monitor (FM), or neutron polarimeter. The "sweep" magnet removes charged particles and also precesses in-plane neutron spin components (see text).

downstream of the main CSB detector arrays, as indicated schematically in Fig. 2. In addition, an integral (relative) neutron flux measurement was obtained with a scintillator (FM) that completely intercepted the beam, just downstream of the PPT.

The PPT is based on the "spin refrigerator" method. Its operation is very similar to the device described in Ref. 29, the main difference being in the sample sizes used. The target material comprises single crystals of Yb-doped yttrium ethyl sulfate (YES): Y(C$_2$H$_5$SO$_4$)$_3$ . 9H$_2$O. The Yb ions in this crystal have a strongly anisotropic g-factor; they can be easily polarized at a temperature of ~ 0.5 K in a ~ 1.2 T polarizing field parallel to the crystal axis, and this polarization can then be efficiently transferred to the free protons by rotating the crystal axis until it is nearly perpendicular to the field. In practice, the sample was polarized by continuous rotation (at ~ 40 Hz) for 2-3 hours. Features of the target well suited to our experimental configuration were its modest requirements on polarizing field uniformity and cryogenics and its long spin relaxation times at quite low values of the holding field. The compact superconducting polarizing- and holding-field coils allowed large vertical and horizontal acceptance of the detection systems and easy adiabatic reversal of the holding field to flip the target spin. The target thickness was 1.0 g/cm$^2$, with lateral dimensions chosen to match the nominal size of the neutron beam (5 cm $\times$ 7 cm). The holding field applied to the (non-rotating) target during data acquisition was 590 G, and the polarization achieved, averaged over production running periods, was typically 0.42.

The main CSB detector arrays are shown schematically in Fig. 3. Neutron-proton elastic scattering events from the PPT were identified by detecting the scattered neutron and recoil proton in coincidence. Because YES has 9 times as many bound as (polarizable) free protons, the detectors must have position sensitivity to distinguish free from quasifree scattering by the angular correlation of the detected nucleons. The detector arrays used are left-right symmetric arms, each sensitive to both neutrons and protons, covering a wide angular range (24° to 62° in the lab). Each arm comprises a thin wedge-shaped timing and $\Delta$E scintillator
for the protons, followed by two pairs of x and y multi-wire proportional chambers (MWPC) with individual wire readout. The wire chambers determine the proton angle and permit ray-tracing to the point of origin at the PPT. Following the MWPC's are large position-sensitive neutron counters. These detectors were constructed by segmenting large vats of (mineral-oil-based) liquid scintillator into 96 cells separated by 0.1 cm thick highly polished aluminum walls. Each cell is 8 cm wide by 10 cm high at the front face of the detector, is 40 cm deep, and has its own photomultiplier tube mounted at the rear. The phototubes on the plastic and liquid scintillators could all be optically pulsed with a nitrogen laser system, enabling a variety of tests, calibrations, and diagnostics. During acquisition, software checking on an event-by-event basis continuously monitored the data for a wide variety of anomalies that might signal hardware problems, and alerted the user whenever the "worry level" for any one type of error was exceeded.

Fig. 3. Schematic top view of experimental setup showing configuration of CSB detector arrays. Polarized neutrons are incident from the left. A valid CSB event (np scattering from the polarized target), as defined by the electronics, has no signal in veto scintillators V1, V2 or in the flux monitor FM, but causes coincident signals in the start scintillator S and at least 3 of the 4 wire chambers on one arm, and in at least one cell of the liquid scintillator neutron detector on the opposite arm.

The data were taken in nominal 24 hour cycles (each producing \(~ 1.1 \times 10^6\) detected free np scattering events) during the course of 3 separate production running periods. During each cycle, the primary proton beam spin was flipped every \(~ 30\) seconds, and the PPT spin was flipped (a process that required about 40 seconds) by field reversal every 10 minutes. Approximately every two hours, the PNF sweep magnet polarity was reversed, effectively flipping horizontal neutron beam spin components (to cancel certain systematic errors, to be discussed later, that arise from in-plane spin correlations). Also every two hours, and interspersed with the above, the YES crystals were rotated by \(180^\circ\) (to cancel the spurious \(\Delta A\) that might arise from a small variation of target polarization with depth in the crystals). Approximately every 12 hours, the PPT was repolarized and its spin reversed with respect to the holding field (by adiabatic fast passage\(^30\)) in order to distinguish true spin-dependent from instrumental field-dependent asymmetries. In all, 41 cycles of production np scattering data were obtained, along with 25 cycles of assorted systematic error tests. During each cycle, 15-20% of the time was spent acquiring data with a "dummy" target (which simulated the YES crystals without hydrogen) for background subtraction.
ANALYSIS PROCEDURES

Events from free np scattering were distinguished from background processes (such as quasifree scattering from complex nuclei in the PPT) by a variety of software conditions placed on the experimental observables. After first subtracting accidental np coincidences (\(\lesssim 5\%\) of the real coincidences), we selected free-scattering events by the following sequence of cuts: a “beam spot” cut on the coordinates of the event origin at the target plane, determined from proton ray-tracing; a coarse cut on the incident neutron energy determined from the start time of the event with respect to the cyclotron RF signal; a cut on the azimuthal opening angle \((\phi_{\text{open}})\), or coplanarity, of the detected pair; two-dimensional windows on both \(\Delta E\) of the proton in the wedge scintillator and neutron detector time-of-flight vs. \(\theta_{\text{proton}}\); a two-dimensional cut on \(\theta_p\) vs. \(\phi_p\) to define the solid angle for the detected protons; and a cut on the polar opening angle \((\theta_{\text{open}})\) for the np pair. Both loose and tight versions of these cuts were applied during the replay. Here “loose” (“tight”) refers to a cut at a contour level of roughly 1% (5%) of the free-scattering peak yield after dummy subtraction.

Fig. 4. Opening-angle spectra for np coincidence events induced in the polarized proton target (PPT) and in a nearly hydrogen-free “dummy” target. The PPT spectrum is shown both before (RAW) and after (WITH CUTS) application of “loose” free-scattering conditions in the software. The data shown are from a portion of the CSB production running.

The results of such data sorting procedures are illustrated by the \(\theta_{\text{open}}\) spectra in Fig. 4. In this plot, the free np scattering events fall in the peak centered about the kinematically expected value \((\theta_{\text{open}} \simeq 87.3^\circ)\). The upper curve shows the “raw” opening angle spectrum, which includes most left-right coincidence events. Here the free-scattering peak sits on a broad background of roughly 40%. Applying the loose scattering conditions described above (except, of course, for the \(\theta_{\text{open}}\) cut itself) results in the curve labeled “with cuts,” where the background under the peak has been reduced to about 8% with essentially no loss of free-scattering events. As the background at this stage is thought to be primarily quasifree scattering, it is subtracted with the use of the “dummy” target. The lower curve in Fig. 4 shows the \(\theta_{\text{open}}\) spectrum for the dummy target, normalized relative to the other curves by the integrated neutron flux, with the same free-scattering cuts as used for the PPT. (The small free-scattering peak observed in the dummy spectrum reveals a small hydrogen contaminant, whose presence is inconsequential for the results we present here.) It is clear from Fig. 4 that subtraction of the dummy data will still leave a residual “background” on the order of 0.5%. Uncertainties in the treatment of this residual background represent a major remaining source of systematic error in our \(\Delta A\) result, as we discuss later.
Events which pass all the free-scattering conditions, including the final cut on $\theta_{\text{open}}$, are sorted into spectra as a function of left or right proton recoil angle. They are also sorted separately for the four beam and target spin combinations, giving eight independently measured yields per angle bin. From these we can calculate (among other observables) the neutron asymmetry $P_b A_n$ and corresponding proton asymmetry $P_t A_p$, where $P_b$ and $P_t$ represent the beam and target polarizations, respectively. In Fig. 5, we show these measured asymmetries as a function of center-of-mass angle. As expected on the scale of this figure, the two sets of data points appear virtually identical except for a normalization factor that reflects the inequality of the beam and target polarizations: $P_b / P_t \neq 1$.

If we had sufficiently precise independent measurements of the absolute polarizations $P_b$ and $P_t$, we could now extract the CSB observable $\Delta A(\theta) \equiv A_n(\theta) - A_p(\theta)$ directly from the asymmetries in Fig. 5. Unfortunately, there are at present no known techniques for such accurate ($< 0.1\%$) polarization measurements. Our approach is instead to extract values for $P_b$ and $P_t$ from the measured np scattering asymmetries themselves, and then to confine our attention to those components of $\Delta A(\theta)$ that manifest themselves as shape differences between $A_n(\theta)$ and $A_p(\theta)$. We will spell out the implications of this restriction in more detail later.

Extraction of the polarizations from the measured asymmetries reveals a very important potential source of systematic error in the analysis. The problem can be easily visualized upon examination of Fig. 6, where beam and target polarizations are shown in the top portion, and the zero-crossing angle ($\theta_o$) of the average analyzing power in the bottom panel, as a function of neutron bombarding energy near the nominal average value of 183 MeV. The plotted points are obtained by fitting our measured asymmetries $P_b A_n(\theta)$ and $P_t A_p(\theta)$ with Arndt’s energy-dependent phase-shift calculations\textsuperscript{26} for the analyzing power, scaled by an adjustable normalization factor and shifted by an adjustable angle offset; the method used for binning the neutron energy will be explained shortly. We note from Fig. 6 that $P_b$ and $\theta_o$ vary strongly with neutron bombarding energy, while $P_t$ is constant, as it should be. (The variation of $P_b$ is not qualitatively surprising, as it is well known\textsuperscript{27} that the effective polarization transfer in the $^2\text{H}(p, n)$ production reaction depends both on the incident proton energy and on the excitation energy within the recoiling two-proton system.) Since the beam polarization is larger at the higher energies, the neutron analyzing power $A_n(\theta)$ is measured at an effective average energy (weighted by $P_b$) that is somewhat higher than that for $A_p(\theta)$. Since the analyzing power itself is energy-dependent, as indicated by the variation of $\theta_o$, this effective energy difference leads to a spurious $\Delta A(\theta)$ if left uncorrected. (For the range $T_n = 170-193$ MeV, our results indicate that the energy difference between $A_n$ and $A_p$ measurements is $\sim 0.52$ MeV, corresponding to a zero-crossing angle shift in the lab frame of $0.076^\circ$; these numbers roughly double if we consider the range 158-193 MeV, which includes all of the data presented in Fig. 6.)

In order to minimize this false contribution to $\Delta A(\theta)$, it is essential to analyze the data in relatively narrow neutron energy ($T_n$) “slices.” The best information we have on $T_n$ on an event-by-event basis comes from the event arrival time with respect to the cyclotron RF signal ($t_{RF}$). Because of our finite time resolution ($\sim 0.65$ ns FWHM), a narrow bin in the $t_{RF}$ spectrum still incorporates contributions from a $T_n$-distribution of appreciable width, and therefore with a non-vanishing spurious contribution to $\Delta A(\theta)$. In order to correct for the small residual error in the results for each $t_{RF}$ bin, we need to know the shape of the contributing $T_n$-distribution, the variation of $P_b$ with neutron energy, and the energy derivative $\partial A(\theta)/\partial T_n$. To determine the first of these factors, we fit the observed $t_{RF}$ spectra by convoluting a model $T_n$-spectrum with a Gaussian time resolution function and a
Fig. 5. The measured asymmetries $P_bA_n$ and $P_tA_p$ for np elastic scattering at $T_n = 183$ MeV as a function of c.m. angle. The data from all CSB production runs has been collected in 2.4° proton lab angle bins, leading to statistical errors roughly one-tenth the size of the plotted symbols.

Fig. 6. Beam ($P_b$) and target ($P_t$) polarizations (top panel) and average analyzing power zero-crossing angle (bottom panel) as a function of neutron bombarding energy $T_n$. The plotted points have been extracted by fitting the measured asymmetries within neutron energy "slices," as described in the text.
simple analytical approximation for the energy-dependence of the event detection efficiency. The semiempirical energy spectrum model incorporates (with adjustable parameters) the expected variation of the $^2\text{H}(p,n)^2p$ production cross section with excitation energy of the $2p$ system, with bombarding energy (as the incident proton is degraded upon traversal of the LD$_2$ target), and with production angle within the acceptance of the neutron collimator. A typical fit to the charge-exchange peak region of the $t_{RF}$ spectrum is shown in Fig. 7.

Fig. 7. Observed neutron time-of-flight (to the PPT) spectrum with respect to the cyclotron RF, showing the pronounced charge-exchange peak in the production reaction. A model (see text) of the underlying neutron energy distribution gives the solid line fit to the measured spectrum. The range included in our analysis, corresponding to mean neutron energies between 170 and 193 MeV, is indicated by arrows.

The $P_b$, $P_t$, and $\theta_o$ values presented in Fig. 6 were thus extracted by analyzing the data in separate $t_{RF}$ bins, and then plotting each point at the centroid energy determined for that bin from our $t_{RF}$ spectrum model. The solid curve through the $P_b$ data points in the upper panel represents a fit to the measurements obtained by convoluting an analytical expression for $P_b(T_n)$ with the $T_n$-distribution deduced for each $t_{RF}$ bin. It is this empirical fit for $P_b(T_n)$ that we use to determine the correction to $\Delta A(\theta)$ for the remaining energy width of each "slice": correction $= -\partial A(\theta)/\partial T_n \cdot [< P_b T_n > / < P_b > - < T_n >]$, where the angular brackets denote averaging over the $T_n$-distribution for the "slice" of interest. For the energy derivative of the analyzing power, we use Arndt (Spring '88) phase shift calculations, represented by the solid curve for $\theta_o(T_n)$ in the lower panel of Fig. 6. While there is an appreciable overall offset in $\theta_o$ values between our measured asymmetries and the phase shift calculations, the data and calculations agree well for the energy derivative near 180 MeV. The calculated energy derivative, and hence the $\Delta A(\theta)$ correction, is reduced at angles away from $\theta_o$.

The data to be presented in the following sections are averaged over the highest 6 energy slices (between 170 and 193 MeV, with a centroid of 183 MeV) indicated in Fig. 6. A residual correction to $\Delta A(\theta)$ was made, as described above, for each energy slice before the results for different slices were averaged. The corrections vary from a few $\times 10^{-4}$ at the highest energies to almost $2 \times 10^{-3}$ for the lowest slice considered, with an average correction of $\sim 7 \times 10^{-4}$. If we were alternatively to treat this entire energy range as a single slice, we would have to apply a $\Delta A(\theta)$ correction of $\sim 16 \times 10^{-4}$. These two alternative procedures yield final corrected values of $\Delta A$ that agree with one another within $\sim 1 \times 10^{-4}$, providing strong support for the accuracy of our correction procedure.
RESULTS

We return now to the extraction of CSB observables from the data of Fig. 5, keeping in mind that we actually work with analogous data for individual neutron energy slices. Furthermore, we will concentrate for now on a restricted angular range, corresponding to the 8 angle bins in Fig. 5 bracketing the zero-crossing. This temporary restriction is imposed because quantitative analysis of systematic errors at the more extreme angles will require more time than we have had available since completing the third round of replay.

Recall that the absolute magnitudes of \( P_b \) and \( P_t \) are not known with sufficient accuracy to extract \( \Delta A(\theta) \) straightforwardly from the measured asymmetries \( P_b A_n(\theta) \) and \( P_t A_p(\theta) \). There are, however, several ways to proceed which do not require detailed knowledge of the polarizations. One approach is to compare the analyzing power zero-crossing angles of \( A_n(\theta) \) vs. \( A_p(\theta) \), extracted by fitting smooth curves to the \( P_b A_n \) and \( P_t A_p \) distributions. This was in fact the approach of the TRIUMF group.\(^7\)

We have chosen a somewhat more general approach, to provide additional sensitivity to shape differences between \( A_n(\theta) \) and \( A_p(\theta) \) at angles away from the zero-crossing. Since the measurements are made simultaneously over the full angular range, \( P_b \) and \( P_t \) take on the role of angle-independent normalization constants. Our experimental uncertainty (\(~\) a few \%) in the ratio \( P_b/P_t \) then makes us insensitive to terms in \( \Delta A(\theta) \) that would be manifested solely by a normalization difference between \( A_n(\theta) \) and \( A_p(\theta) \), namely, to any contribution to \( \Delta A \) proportional to the average analyzing power \( A(\theta) = [A_n(\theta) + A_p(\theta)]/2 \). That is, what we can strictly determine is the quantity

\[
\text{"} \Delta A(\theta) \text{"} = \Delta A_{\text{true}}(\theta) + c \cdot A(\theta), \tag{1}
\]

where \( c \) is an unknown constant whose value reflects the difference between our best evaluation of \( P_b/P_t \) and the actual value of this ratio. There are a number of equally justifiable alternative procedures for determining \( P_b/P_t \) from the measured asymmetries, and they will in general yield slightly different values for this ratio, and thus different "\( \Delta A(\theta) \)". For example, one procedure involves adjusting the polarization ratio until the error-weighted rms value of the extracted "\( \Delta A(\theta) \)" over some angle range is minimized; the resulting minimal "\( \Delta A(\theta) \)" is equivalent to that part of \( \Delta A_{\text{true}}(\theta) \) which is "orthogonal" to \( A(\theta) \) over the specified range.

Values for the "minimal" "\( \Delta A(\theta) \)" extracted via the above prescription, and averaged over energy slices, are displayed in Fig. 8 (with statistical error bars only) for the eight angle bins surrounding the \( A(\theta) \) zero-crossing (\( \theta_{\text{zero}} \approx 96^\circ \), indicated by the arrow). The use of alternative prescriptions to deduce \( P_b/P_t \) would lead in this figure to a "rotation" of the data points about a pivot located at the \( \Delta A \) value for \( \theta_0 \). Nonetheless, it is clear from Fig. 8 that the minimum rms value of "\( \Delta A \)" is already significantly different from zero, providing clear evidence for charge symmetry breaking. A quantitative comparison of the "\( \Delta A(\theta) \)" angular distribution with theory could be made by subjecting predicted curves to the same condition (minimizing the rms value of \( \Delta A + c A \)) that we have applied in extracting the data, but we will pursue an alternative path here.

From Eq. (1) we note that it is possible to extract, in a model-independent fashion, the mean value of \( \Delta A(\theta) \) over any ("magic") angular region for which the average analyzing power is zero: \( <A(\theta)> = 0 \). For such regions we have \( <\text{"} \Delta A(\theta) \text{"}> = <\Delta A_{\text{true}}(\theta)> \), and the ambiguity due to the polarization ratio uncertainty vanishes. We present our experimental average values for such "magic" angle ranges in Fig. 9 and Table I. Our primary result is averaged over the "total region" \( 82.2^\circ \leq \theta_{\text{cm}} \leq 116.1^\circ \): \( <\Delta A(\theta)>_{\text{total}} = (32.1 \pm 6.1 \pm 6) \times 10^{-4} \), where the first error is statistical and the second our present estimate of the systematic
Fig. 8. Measured values of the CSB observable $\Delta A(\theta)$ averaged over 6 neutron energy “slices” between 170 and 193 MeV. The “$\Delta A(\theta)$” values may differ from $\Delta A_{\text{true}}(\theta)$ by a constant times $A(\theta)$ (whose zero-crossing angle is indicated by the arrow), as a result of the experimental uncertainty in the polarization ratio $P_s/P_t$ (see text).

Fig. 9. Results for $\Delta A(\theta)$ averaged over “magic” angular regions straddling the zero-crossing of $A(\theta)$, for which $<A(\theta)> = 0$. The left panel presents the data for the range $82.2^\circ \leq \theta_{\text{cm}} \leq 116.1^\circ$. The same data are presented in the right panel, subdivided into “inner” and “outer” regions (see Table I), to give a coarse measure of the “curvature” of $\Delta A$. The error bars are statistical only. The calculations of Ref. 14, averaged over the same angular regions as the data, are indicated by the horizontal lines: the dotted line is the electromagnetic contribution; the dashed line includes in addition the effect of the np mass difference on one-boson exchange; the solid line is the full calculation, including $\rho - \omega$ mixing.

error (to be discussed in the next section). If we add the statistical and systematic errors in quadrature, our result differs from zero by nearly 4 standard deviations; more importantly, it differs from the value expected for the electromagnetic spin-orbit contribution alone (the dotted horizontal line labeled “$\gamma$” in Fig. 9) by 3 standard deviations. The other horizontal
lines in Fig. 9 represent the CSB calculations of Holzenkamp, Holinde, and Thomas\textsuperscript{14} (see Fig. 1) averaged over the same “magic” angle region as the data. Note the very good agreement between our result and the full Bonn potential prediction.

In the right-hand panel of Fig. 9, the “total” angle region considered above has been subdivided into “inner” and “outer” “magic” regions (the latter composed of two parts), and the same $\Delta A$ data included in the left-hand panel have been averaged separately for these two subregions. The exact angular limits of the “inner” and “outer” regions are given in Table I. The difference in value between $< \Delta A >_{\text{inner}}$ and $< \Delta A >_{\text{outer}}$ is essentially a measure of the “curvature” of $\Delta A(\theta)$ near $\theta_0$. (The ambiguity embodied in eq. (1) renders us insensitive to the slope of $\Delta A(\theta)$ near $\theta_0$). Figure 9 shows that the small curvature suggested by the data is again consistent with the full Bonn potential calculation.

<table>
<thead>
<tr>
<th>C.M. Angle Range</th>
<th>Experimental Value\textsuperscript{a}</th>
<th>Theoretical Predictions from Ref. 14 $\gamma$ alone $\pi, \rho$ - exch. $\rho - \omega$ mix</th>
</tr>
</thead>
<tbody>
<tr>
<td>“Total”</td>
<td>32.1 ± 6.1 (± 6)</td>
<td>6.5</td>
</tr>
<tr>
<td>(82.2° - 116.1°)</td>
<td></td>
<td>16.4</td>
</tr>
<tr>
<td></td>
<td></td>
<td>27.4</td>
</tr>
<tr>
<td>“Inner”</td>
<td>38.1 ± 8.4 (± 6)</td>
<td>7.2</td>
</tr>
<tr>
<td>(88.4° - 106.3°)</td>
<td></td>
<td>16.7</td>
</tr>
<tr>
<td></td>
<td></td>
<td>29.8</td>
</tr>
<tr>
<td>“Outer”</td>
<td>25.5 ± 8.8 (± 6)</td>
<td>5.7</td>
</tr>
<tr>
<td>(82.2° - 88.4°)</td>
<td></td>
<td>16.2</td>
</tr>
<tr>
<td></td>
<td></td>
<td>24.8</td>
</tr>
<tr>
<td>(106.3° - 116.1°)</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

\textsuperscript{a} Systematic errors are quoted in parentheses. The relative systematic error between “inner” and “outer” regions is probably smaller than $6 \times 10^{-4}$.

**SYSTEMATIC ERRORS**

Measurement of $\Delta A(\theta)$ to the desired level of precision is subject to many potential systematic errors, some of which may vary significantly with angle. In order to present first results at this Workshop we have concentrated in the analysis thus far on evaluating the systematic uncertainties for $< \Delta A >$ over the full “magic” region 82.2° - 116.1°. A list of the errors considered, and our present best estimates of their magnitude, is given in Table II. We believe that some of these uncertainties will be reduced by further analysis.

One of the largest systematic uncertainties is associated with the subtraction of background events. To the extent that the background arises from interactions of the neutron beam with bound (unpolarized) protons in the contaminant nuclei in the PPT, it will in general be characterized by non-zero $P_n A_n(\theta)$ but by $P_p A_p(\theta) = 0$. Thus, imperfect subtraction of these quasifree events via the dummy-target data can leave a residual spurious contribution to $\Delta A$. The dummy was constructed to simulate the non-hydrogenic content of the PPT (in both composition and thickness) as closely as possible, and the relative neutron fluxes for...
runs on the two targets were carefully monitored. It is nonetheless important for minimizing
the background subtraction error to develop an independent prescription to determine the
relative normalization of dummy vs. PPT data to a precision of a few percent.

<table>
<thead>
<tr>
<th>Source of Error</th>
<th>Error Estimate</th>
</tr>
</thead>
<tbody>
<tr>
<td>1) Background subtraction (dummy/PPT normalization)</td>
<td>3.0</td>
</tr>
<tr>
<td>2) Beam polarization difference for dummy vs. PPT</td>
<td>1.5</td>
</tr>
<tr>
<td>3) n-detector spin dependence</td>
<td>1.0</td>
</tr>
<tr>
<td>4) (P_b) variation over neutron energy spread</td>
<td>1.5</td>
</tr>
<tr>
<td>5) Choice of software cuts (&quot;loose&quot; vs. &quot;tight&quot;)</td>
<td>3.5</td>
</tr>
<tr>
<td>6) In-plane spin correlation effects</td>
<td>0.2</td>
</tr>
<tr>
<td>7) Proton bending in the PPT holding field</td>
<td>0.5</td>
</tr>
<tr>
<td>8) Field-dependent gain shifts</td>
<td>0.5</td>
</tr>
<tr>
<td>9) Spin-correlated beam motion</td>
<td>0.5</td>
</tr>
<tr>
<td>10) Effective angle-dependence of (P_b/P_t)</td>
<td>0.5</td>
</tr>
<tr>
<td>11) Correlated changes in (P_b/P_t) and other parameters</td>
<td>0.5</td>
</tr>
<tr>
<td>12) &quot;10-minute demons&quot;</td>
<td>0.5</td>
</tr>
<tr>
<td>13) Overall (\Delta A(\theta)) normalization uncertainty</td>
<td>1.6</td>
</tr>
</tbody>
</table>

Sum in quadrature = 5.6

The dummy normalization prescription we have used is illustrated in Fig. 10. The \(\theta_{open}\)
spectrum in the top panel reveals that, even after dummy subtraction, there are "wings" on
the free-scattering peak. Presumably, the two main sources of events in these "wings" are
imperfectly subtracted quasifree scattering and "misplaced" free scattering. Free-scattering
events may be "misplaced," for example, when the neutron interacts twice in the liquid scin­
tillator but produces sufficient light to surpass our detection threshold only in the second
interaction. The two contributions to the "wings" can be distinguished by their spin correla­
tion parameter \(C_{NN}\): quasifree scattering from unpolarized protons must have \(P_bP_tC_{NN} \equiv 0\),
while the "misplaced" events should be characterized approximately by the (nearly angle­
independent) free-scattering value \(C_{NN} \approx 0.65\). Thus, when the dummy normalization is
correct, the extracted \(C_{NN}\) should have roughly the same value in the wings as in the peak
of the \(\theta_{open}\) spectrum, while undersubtraction (oversubtraction) of quasifree events will yield
too small (large) a value of \(C_{NN}\) in the wings. Results for \(C_{NN}\) vs. \(\theta_{open}\) are shown in Fig.
10 for three different values of the effective dummy/PPT thickness ratio (\(t_D/t_P\)), illustrating
the strong sensitivity to small normalization changes. From these data we conclude that
\(t_D/t_P = 0.94 \pm 0.04\). [When we use this value for the ratio, the "wings" in the dummy­
subtracted \(\theta_{open}\) spectrum (top panel of Fig. 10) agree qualitatively in shape and magnitude
with Monte-Carlo simulations including the effects of neutron rescattering in the liquid scin­
tillator.] This 4% normalization uncertainty translates into the \(< \Delta A >\) systematic error of
\(3.0 \times 10^{-4}\) quoted for the background subtraction (item 1) in Table II.

Another error (item 2) allows for possible small differences in the average beam po­
larization between the dummy and PPT runs. The quoted error \((1.5 \times 10^{-4})\) in Table
II corresponds to a polarization difference of 0.02, the maximum value consistent with our
neutron polarimeter data.

There can also be systematic errors associated with misplaced free-scattering events if
they are removed by our cuts in a spin-dependent way. For example, our final cut on \(\theta_{open}\)
may preferentially remove neutrons which have been scattered toward the left rather than the
Fig. 10. Illustration of the procedure used to optimize the relative normalization of dummy to PPT data. The top panel shows the opening-angle spectrum for free np scattering events (within the range $30^\circ \leq \theta_{\text{lab}} \leq 50^\circ$) after dummy subtraction, with the arrows indicating the "loose" software gate on $\theta_{\text{open}}$ applied as the final free-scattering cut. The bottom panel shows the extracted values of the spin correlation parameter $C_{NN}$ as a function of $\theta_{\text{open}}$ for three choices of the effective dummy/PPT thickness ratio. When the relative normalization is chosen correctly, we expect $C_{NN}$ to be roughly constant at the average free-scattering value indicated by the horizontal lines.

right in their first interaction in the liquid scintillator. The survival rate of such misplaced events then depends on (poorly known) ${n}$-carbon analyzing powers and on the polarization of the free-scattered neutrons, which has different sensitivities to the beam and target spins. The associated systematic error, labeled by "n-detector spin-dependence" (item 3) in Table II, is estimated to be no worse than $1.0 \times 10^{-4}$ for "loose" cuts by investigating the dependence of the extracted $<\Delta A>$ on the $\theta_{\text{open}}$ gate used.

Another sizeable systematic error (item 4 in Table II) is associated with our corrections to $\Delta A$ (discussed earlier) to compensate for the beam polarization variation over the finite energy width of the neutron beam. In our $T_{n}$-"slice" analysis, we have limited the net correction to $<\Delta A>$ to $7.3 \times 10^{-4}$ for the angular range considered here. The systematic error quoted in Table II allows for a 20% uncertainty in this net correction. The results obtained with different widths of the energy slices agree with one another to well within this
estimated error.

The largest single entry in Table II (item 5) reflects the difference between the $\langle \Delta A \rangle$ results obtained with "loose" vs. "tight" free-scattering cuts at the present level of the analysis. The result we have quoted above, $\langle \Delta A \rangle_{\text{total}} = 32.1 \times 10^{-4}$, results from loose cuts; with tight cuts on all parameters we find $\langle \Delta A \rangle_{\text{total}} = 38.5 \times 10^{-4}$. Some of this difference may arise from statistical fluctuations associated with the removal of ~20% of the free-scattering events and ~70% of the background by the tight cuts: the statistical uncertainty in the difference between loose- and tight-cut results is $\pm 1.6 \times 10^{-4}$. Furthermore, some of the difference has presumably been accounted for already in systematic errors explicitly considered above: the sensitivities to dummy normalization errors and to misplaced free-scattering removal errors are, for example, very different for the loose- vs. tight-cut data samples. In our replay to date, we have carried out the "calibrations" relevant to estimating these latter errors (e.g., the $C_{NN}$ vs. $\theta_{\text{open}}$ analysis of Fig. 10), as well as the modeling of the neutron energy spectrum, only for events subjected to the loose cuts; we thus emphasize the corresponding $\langle \Delta A \rangle$ values for now. In order to allow, without excessive double-counting, for the possibility that the cut-dependence reveals some additional source of systematic error, we have included roughly half of the observed change (i.e., $3.5 \times 10^{-4}$) in $\langle \Delta A \rangle$ as a separate item in Table II.

Several of the entries in Table II represent sources of error that we have successfully "beaten down" by appropriate design of the equipment and procedures for the experiment. Our measured beam-spin- and target-spin-dependent left-right asymmetries can have small contaminating contributions from in-plane spin correlations (i.e., from $C_{LS}$ and $C_{SL}$, which reach 0.5 in magnitude for np scattering over parts of the angle range covered). The corresponding systematic error (item 6) in $\langle \Delta A \rangle$ has been kept to $\sim 2.0 \times 10^{-4}$ by minimizing the in-plane polarization components for both the target and the beam. For the target this was done by careful vertical alignment and monitoring of the holding field; for the beam we averaged equal amounts of data acquired with $+90^\circ$ and with $-90^\circ$ precession of horizontal components in the PNF sweep magnet. Similarly, we have minimized instrumental asymmetries associated with the PPT holding field by keeping the field strength relatively low (590 G) and by averaging equal amounts of data acquired with the target spin parallel (NRM mode) and antiparallel (AFP mode) to the field. Field-dependent asymmetries can arise both from the bending of detected protons (item 7 in Table II) and from induced gain (hence efficiency) shifts in the neutron detector phototubes (item 8). The bend angle ($\sim 0.5^\circ$ over the angle range covered) was corrected in software to an average accuracy of $\pm 0.01^\circ$, as monitored via the stability of the $\theta_{\text{open}}$ peak centroid with respect to field direction. Gain shifts were minimized by magnetic shielding of the phototubes, combined with compensation coils installed on the detector arms to offset effects of the (unclamped) holding field. The rms field-dependent gain shift for all cells, monitored via the pulse height of free-scattering protons in the liquid scintillator, was $\lesssim 0.2\%$. Residual field effects were further suppressed by a factor $\gtrsim 40$ by the NRM + AFP averaging; this was estimated by noting the change in the final $\langle \Delta A \rangle$ values when the software correction for the bending of the protons was purposely miscalibrated. The resultant systematic errors for field-induced bending and gain shifts are thus each no worse than $0.5 \times 10^{-4}$.

Other small systematic errors, evaluated from features of the data themselves (in some cases combined with Monte-Carlo simulations), arise from: 1) motion of the neutron beam position centroid correlated with the primary proton beam spin orientation (estimated to be $< \pm 0.03$ mm); 2) an effective angle-dependence of the $P_b/P_t$ ratio due, for example, to spatial polarization profiles coupled with angle-dependent changes in the fraction of the target volume...
sampled by the detectors (this is a more serious effect at the larger proton laboratory angles that have so far been omitted from our $\Delta A$ analysis); 3) accidental correlations among changes (e.g., from one production run to another) in $P_3$ or $P_4$ and changes in other parameters of the experiment (bombarding energy, left/right detection efficiency ratio, etc.), so that changing conditions are not averaged equally in the $P_3 A_n$ and $P_4 A_p$ measurements; 4) any short-term change in the left/right detection efficiency ratio that happens to occur for only one target spin orientation (referred to as a “10-minute demon” in Table II because the target spin was reversed every 10 minutes).

Finally, we note that the extracted values of $\Delta A$ are subject to an overall normalization uncertainty, estimated to be $\pm 5\%$, arising primarily from uncertainties in the absolute values of the polarizations. This scale uncertainty contributes a systematic error of $1.6 \times 10^{-4}$ to $<\Delta A>_{\text{total}}$. A few other known sources of systematic error have been omitted from Table II, either because their effects are expected to be appreciable only at angles outside the range we have considered here (e.g., the effect of the hydrogen contaminant in the dummy), or because they overlap strongly with included effects (e.g., the effects of a second interaction in the PPT presumably are among the differences considered between loose and tight cuts).

The sum in quadrature of the errors in Table II is $5.6 \times 10^{-4}$. Because we do not believe that our current evaluations of these effects permit quoting two significant digits, and in part to allow for possible neglected effects, we round upward to obtain our quoted net systematic error of $\pm 6 \times 10^{-4}$.

**DISCUSSION**

Our primary result, $<\Delta A>_{\text{total}} = (32.1 \pm 6.1 \pm 6) \times 10^{-4}$, is somewhat larger in magnitude than many theoretical predictions. For example, a calculation carried out by Beyer and Williams,$^{18}$ using the Reid soft-core potential for the distortions, when averaged over the same angle region ($82.2^\circ$ - $116.1^\circ$) as the data, yields (in units of $10^{-4}$) a value of 18.2 when $\rho - \omega$ mixing is included and 14.0 when it is not. The corresponding numbers for the Bonn potential calculation$^{14}$ (see Table I) are 27.4 and 16.4. As is clear from these numbers, the major difference among the various calculations is in the magnitude of the $\rho - \omega$ mixing contribution. This difference can in turn be attributed to different choices for the $\rho$-NN and $\omega$-NN coupling constants, which are not well determined in NN scattering analyses. The fitted coupling constants characteristic of the Bonn potential (e.g., Ref. 14 uses $g_{\omega}/4\pi = 0.77$ and $g_{\omega}/4\pi = 23$) are quite a bit larger than the more conventional values (0.55 and 8.1, respectively) assumed in the Reid calculation$^{18}$ cited above.

Our experimental CSB result clearly favors the calculations employing larger values of the vector meson coupling constants, but cannot rule out the smaller values. First, if one combines statistical and systematic uncertainties, our measurement is only 1.6 standard deviations away from the Reid potential calculation. In addition, there are known deficiencies in the calculations, not the least of which may be the neglect of additional graphs (e.g., simultaneous $\pi - \gamma$ exchange) which might lead to a slightly larger predicted $\Delta A(\theta)$, and thereby reduce the discrepancy with calculations using the smaller coupling constant values.

As discussed briefly above (see Fig. 9 and Table I) we have also made a crude measurement of the curvature of $\Delta A(\theta)$. In this case, the statistical precision is insufficient to distinguish among different calculations, but the results again agree very well with the Bonn potential calculation (horizontal lines in the right-hand panel of Fig. 9), where the predicted curvature arises predominantly from the $\rho - \omega$ mixing term.

One method for extending the theory–experiment comparison to finer angle bins than considered in Fig. 9 involves the alternative CSB observable$^{31}$ $X(\theta) \equiv$
Fig. 11. The measured CSB observable \( X \approx [\Delta A/A + \Delta P/P]/2 \) as a function of c.m. scattering angle. The horizontal solid line, indicating the expected value of \( X(0) \) in the absence of CSB, represents \( \Delta P/2P \approx 0.16 \), the value extracted by minimizing the rms value of \( \Delta A \). The other curves represent the theoretical predictions of Ref. 14 with (dashed curve) and without (dot-dashed curve) inclusion of \( \rho - \omega \) mixing. To allow for errors in the extracted value of \( \Delta P/P \), we have added an adjustable overall vertical offset to each curve, chosen to optimize its agreement with the measurements.

\[
(P_bA_n - P_tA_p)/(P_bA_n + P_tA_p) \approx [\Delta A(\theta)/A(\theta) + \Delta P/P]/2,
\]

where \( P \equiv (P_b + P_t)/2 \) and \( \Delta P \equiv P_b - P_t \). The nice feature of \( X(\theta) \) is that the uncertainty in the beam and target polarizations is manifested here, via the term \( \Delta P/P \), as an angle-independent offset in \( X(\theta) \) values. A disadvantage is that \( X(\theta) \) has a singularity at the analyzing power zero-crossing, and this singularity shifts in angle from one beam energy slice to another. To circumvent this problem, we have extracted \( X(\theta) \) for the 170 - 193 MeV region, treated as a single energy slice, but including the correction for the beam polarization variation over the width of this slice. The results are shown in Fig. 11.

The curves in Fig. 11 are again the theoretical predictions of Ref. 14, with (dashed curved) and without (dot-dashed curve) the \( \rho - \omega \) mixing contribution included. In order to allow for small errors in the extracted polarizations, a slight vertical offset has been added independently to each theoretical curve to optimize its agreement with the data. The results in Fig. 11 show again, from an alternative viewpoint, that inclusion of a strong \( \rho - \omega \) contribution helps significantly to describe our measurements.

The importance of \( \rho - \omega \) mixing has also been emphasized recently in the investigation of other CSB phenomena.\(^{10}\) It plays an essential role in quantitatively explaining observations of \( nn \) vs. \( pp \) scattering length differences and \(^3\)H vs. \(^3\)He nuclear binding energy differences. Furthermore, Blunden and Iqbal\(^{35}\) have shown that within meson-exchange models of CSB NN interactions, understanding of the A-dependence of mirror-nucleus binding energy...
differences (the Nolen-Schiffer anomaly[3]) requires a spin-orbit CSB potential (of class III), such as arises from $\rho - \omega$ mixing. Indeed, they are able to account for roughly 75% of the observed mass differences for all $A \leq 41$, provided that they use the large $\rho$-NN and $\omega$-NN coupling constants characteristic of the Bonn potential. It is thus apparent that the attempt to provide a consistent theoretical explanation for manifestations of CSB in nuclear physics, including the present result, is teaching us something important about the role of vector meson exchange in the NN interaction at ranges $\lesssim 1$ fm.

**SUMMARY**

We have very clearly observed charge symmetry violation in np elastic scattering over a broad angular range at 183 MeV. Our result for $\Delta A(\theta)$ averaged over the c.m. angular range $82.2^\circ$ to $116.1^\circ$ is $(32.1 \pm 6.1 \pm 6) \times 10^{-4}$. When statistical and systematic errors are added in quadrature, this result differs from zero (i.e., no CSB) by 4 standard deviations. More significantly, it differs from the value expected from the electromagnetic spin-orbit interaction alone by 3 standard deviations, representing the strongest and most clearcut experimental evidence for CSB in the nuclear force in any system to date. Comparison of $\Delta A(\theta)$ with meson-exchange calculations strongly supports the importance of the $\rho - \omega$ mixing contribution to CSB at 183 MeV. In particular, calculations based on the Bonn prescription for the NN potential, employing relatively large $\rho$-NN and $\omega$-NN coupling constants, reproduce our experimental results very well.

In order to present preliminary results at this Workshop we have focused on only a portion of our data. Before we can present results over our full angular range ($60^\circ \leq \theta_{\text{cm}} \leq 120^\circ$), we must still estimate systematic errors at the extreme angles. Another round of replay is needed in order to make several refinements to the analysis, and to finalize our estimates of certain systematic errors. Results of an independent analysis at Wisconsin will also have to be included. We expect to finish these tasks by early 1990.

**ACKNOWLEDGEMENTS**

The authors are extremely grateful to the members of the IUCF technical staff (especially D.C. DuPlantis and K. Solberg), without whose efforts the experiment would not have been possible. We thank also S. Fray, T. Bowyer, and G. Xu for helping with the data analysis. Many others, including B. Hichwa, P. Schwandt, C. Glover, W.K. Pitts, and many Hope College undergraduate students, contributed in important ways to the early stages of the experiment. Finally, we thank A.W. Thomas, A.G. Williams, and G.A. Miller for help in providing detailed theoretical results.

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Charge Symmetry Breaking in $np$ Elastic Scattering at TRIUMF

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Abstract

The effect of isospin-violating, charge symmetry breaking (CSB) terms in the $np$ interaction has been observed at TRIUMF by measuring the difference in the zero-crossing angles of the neutron and proton analyzing powers, $A_n$ and $A_p$, at a laboratory neutron energy of 477 MeV. The scattering asymmetries were measured in the center of mass angle range from $59^\circ$ to $80^\circ$ with a neutron beam incident on a polarizable proton target. To reduce systematic errors, interleaved measurements of $A_n$ and $A_p$ were made using the same beam and target (apart from their respective polarization states). The difference in zero-crossing angles was $0.340^\circ \pm 0.162^\circ$ ($\pm 0.058^\circ$), which yields $\Delta A = A_n - A_p = 0.0047 \pm 0.0022$ ($\pm 0.0008$) using $dA/d\theta_{cm} = -0.01382$ deg$^{-1}$. The second errors represent systematic effects. This result is in good agreement with recent theoretical calculations. A similar experiment is planned at 350 MeV to determine $\Delta A$ with a statistical error of $\pm 0.0005$.

I. INTRODUCTION

Isospin was the first internal symmetry postulated for the $NN$ force[1], suggested by the very small neutron–proton ($np$) mass difference. By analogy with spin in ordinary space, the nucleon could have "isospin", $T$, oriented up or down in "isospin space". For the $NN$ interaction to be isospin invariant (or charge independent), the Hamiltonian, $H$, must be invariant under arbitrary rotations in isospin space and the $nn$, $pp$, and $np$ forces must be identical. The electromagnetic interactions break charge independence (CIB) and violations on the order of the fine structure constant are expected.

Charge symmetry only requires that $H$ be invariant for $180^\circ$ rotations about the 2–axis in isospin space,

$$\left[ H, P_{CS} \right] = 0, \quad P_{CS} = e^{i\pi T_3}. \quad (1)$$

If $I$ is the isospin of the $NN$ system, then $I = T(1) + T(2)$. The charge symmetry operator reverses the sign of $I_3$, for example changing neutrons into protons and vice versa. For the $NN$ system the $nn$ and $pp$ forces should be identical, but there is no relation between them and the $np$ force. Charge symmetry also requires that isospin be conserved, so that for the $np$ system there is no mixing between the $I = 0$ and $I = 1$ states. In the classification of Henley and Miller[2], we are interested here in Class (IV) $NN$ forces which preserve neither charge symmetry, nor charge independence, nor isospin. Since isospin is not conserved there is mixing of the $I = 0$ and $I = 1$ states of the $np$ system. From the Pauli principle, a change in the isospin state of an $NN$ system necessarily involves a simultaneous transition between spin–singlet and triplet states. Therefore CSB will manifest itself in $np$ spin observables.
Charge symmetry breaking probes an untested aspect of existing theories. In particular CSB contributions from one boson exchange are sensitive to the low partial waves and therefore depend on the details of the short range behavior of the \(NN\) potential. Description of CSB may require explicit quark level characterization of the short distance strong interaction, where conventional nuclear theory breaks down. In the \(NN\) system, the small CSB effects due to one photon exchange, \(\rho-\omega\) meson mixing, and the \(n-p\) mass difference in one boson exchange have been extensively examined in a conventional meson theoretical framework[3–12].

Only two experimental measurements appear feasible given the size of the effects to be observed:

\[
2 \frac{d\sigma}{d\Omega} (A_{\text{oono}} - A_{\text{oono}}) = \text{Re} \, b^* f ,
\]

(2)

\[
2 \frac{d\sigma}{d\Omega} (A_{\text{ookk}} - A_{\text{ookk}}) = \text{Re} \, c^* f .
\]

(3)

Here a four subscript notation \(A_{\text{srbi}}\) is used to denote the polarization states of the scattered, recoil, beam and target particles respectively. In eqs (2) and (3) the index \(k\) refers to polarization along the beam direction, \(s\) is transverse to the beam direction in the scattering plane, and \(n\) is normal to the scattering plane.

The TRIUMF measurement at 477 MeV[13–20] is the first direct measurement of class (IV) forces. The difference between the neutron and proton analyzing powers, \(\Delta A \equiv A_{\text{oono}} - A_{\text{oono}} = A_n - A_p\), was determined at the angle, \(\theta_o\), where the average analyzing power crosses zero. For \(A_n\), the scattering asymmetry was measured using 477 MeV polarized neutrons incident on an unpolarized hydrogen target. For \(A_p\), the polarization states of the beam and target were interchanged. The value of \(\Delta A\) at the zero-crossing was determined indirectly from the difference between the zero-crossing angles of the analyzing power angular distributions. The principle of the measurement is shown schematically in Fig. 1. Phase shift predictions of \(dA/d\theta\) were used to relate the experimentally observed angle difference to \(\Delta A\). The experiment was done using left–right symmetric pairs of detection telescopes[14] to observe \(np\) elastic coincidences. The same equipment and detector configuration were used to measure both \(A_n\) and \(A_p\). The determination of the angle difference is then a null measurement where systematic errors common to both measurements cancel. The \(A_n\) and \(A_p\) measurements were interleaved

![Figure 1. The method employed to extract \(\Delta A = A_n - A_p\) from the difference in the \(A_n\) and \(A_p\) zero-crossing. \(dA/d\theta\) is determined from phase shift analyses.](image)
to eliminate the potential effects of long-term instabilities in the detection equipment and beam properties. Use of apparatus symmetric about the neutron beam direction, and elimination of geometric changes between the \( A_n \) and \( A_p \) phases of the experiment resulted in cancellation of most geometric systematic errors in \( \Delta A \) to at least first order. Further systematic errors not correlated with beam or target polarization reversal were cancelled to at least first order by frequently reversing the neutron and proton polarizations.

The choice of an angle range including \( \theta_0 \) was made to reduce systematic errors to a minimum. In particular, the \( A_n \) and \( A_p \) zero-crossing angles are independent of the beam and target polarizations. In any direct absolute measurement of \( \Delta A \), uncertainties in these latter parameters can dominate the expected CSB effect. If \( P_B \) and \( P_T \) are the beam and target polarizations, respectively, and \( A(\theta) \) is the charge symmetric analyzing power, then the difference between the observed scattering asymmetries can be written

\[
e_B - e_T = A(\theta)(P_B - P_T) + \frac{1}{2} \Delta A (P_B + P_T).
\]  

(4)

Absolute angle measurements were not required, however, all angular parameters were determined as accurately as possible to reduce possible systematic effects and to allow consistency checks to be made.

II. EXPERIMENTAL PROCEDURE

The experimental data-taking included four two-week periods using butanol and a final one-week period with graphite in the FST to determine the inelastic background. During successive four-day long cycles, data were collected with target polarization up and down and unpolarized neutron beam, and with polarized neutrons and an unpolarized target. The holding field was also used in both vertical directions. During the polarized beam portions of the experiment, the proton beam polarization was cycled at random under computer control between up, down and unpolarized states, the ratio of times being 3:3:1 for the different states.

Numerous on-line checks were made on-line to monitor the operation of the FST, position of the proton and neutron beams, the proton beam energy and polarization, and the reliability of the particle detection apparatus. An off-line analysis provided feedback on PHT drifts in the neutron detectors and wire chamber efficiencies were checked. High voltage adjustments to the neutron counters were made if gain drifts exceeded a few percent.

Experimental details have been given in previous publications[13–20]. A schematic diagram of the proton beam transport, the neutron beam production systems, and the experimental layout is given in Fig. 2. The neutron beam was created via the \( ^2H(\bar{p},n)2p \) reaction. Right handed coordinate systems are used with \( z \) along the incident proton direction, \( y \) vertically up, and \( x \) horizontal. Primes denote the final state neutron system. The most important systematic errors are summarized in Table I. The experiment the goal was to reduce each individual source of error in \( \Delta A \) to the level of \( 10^{-4} \) or smaller. The achieved levels are indicated in the last column.

A. Primary Proton Beam

The \( x \) and \( y \) positions of the primary proton beam at the LD\(_2\) target were monitored by
two split plate secondary emission monitors (SEM) which were used in a feedback system to stabilize the beam position. The position and angular uncertainties of the primary beam at the LD\(_2\) target were conservatively estimated to be \(\sigma_x < 0.15\) mm, \(\sigma_y < 0.15\) mm, \(\sigma_{\theta x} < 0.006^\circ\) and \(\sigma_{\theta y} < 0.004^\circ\). The influence of beam motion on systematic errors in \(\Delta A\) is \(< 10^{-4}\). No significant differences in the mean proton beam position for different beam polarization states were observed.

![Diagram](image.png)

**Figure 2.**
Schematic diagram of the proton beam transport, the neutron beam production systems, and the experimental layout. The diagram is not to scale.
Table I

Estimates of Systematic error contributions to $\Delta A$ at $A = 0$.

<table>
<thead>
<tr>
<th>Cause</th>
<th>Experimental Magnitude</th>
<th>Original goal</th>
<th>Actual Error</th>
</tr>
</thead>
<tbody>
<tr>
<td>1. Proton beam stability at LD$_2$</td>
<td>$\Delta x &lt; 0.2 \text{ mm}$</td>
<td>$&lt; 4 \times 10^{-4}$</td>
<td>$&lt; 10^{-4}$</td>
</tr>
<tr>
<td>2. Proton beam energy stability</td>
<td>$&lt; 0.2 \text{ MeV}$</td>
<td>$&lt; 2 \times 10^{-4}$</td>
<td>$&lt; 4 \times 10^{-5}$</td>
</tr>
<tr>
<td>3. Proton beam direction stability</td>
<td>$\Delta \theta &lt; 0.01^\circ$</td>
<td>$&lt; 4 \times 10^{-5}$</td>
<td>$&lt; 10^{-5}$</td>
</tr>
<tr>
<td>4. FST position stability</td>
<td>$\Delta r &lt; 0.25 \text{ mm}$</td>
<td>$&lt; 4 \times 10^{-4}$</td>
<td>$&lt; 10^{-4}$</td>
</tr>
<tr>
<td>5. Proton spin precession</td>
<td>$\Delta \phi &lt; 5^\circ$</td>
<td>$&lt; 10^{-4}$</td>
<td>$&lt; 10^{-4}$</td>
</tr>
<tr>
<td>6. Neutron spin precession</td>
<td>$\Delta \phi &lt; 5^\circ$</td>
<td>$&lt; 10^{-4}$</td>
<td>$&lt; 10^{-4}$</td>
</tr>
<tr>
<td></td>
<td>$P_n \sim 0.035$</td>
<td></td>
<td></td>
</tr>
<tr>
<td>7. Neutron beam stability</td>
<td>$\Delta x &lt; 1 \text{ mm, } \Delta \theta &lt; 0.005^\circ$</td>
<td>$2 \times 10^{-5}$</td>
<td>$2 \times 10^{-5}$</td>
</tr>
<tr>
<td>8. Multiple scattering in FST</td>
<td>$\sigma_B = 0.85^\circ$</td>
<td>$4 \times 10^{-5}$</td>
<td>$4 \times 10^{-5}$</td>
</tr>
<tr>
<td>9. Stability of FST field</td>
<td>$\sigma_B = 3 \text{ G}$</td>
<td></td>
<td>$&lt; 10^{-4}$</td>
</tr>
<tr>
<td>10. Unequal L/R scattering angles</td>
<td>$</td>
<td>\theta_L - \theta_R</td>
<td>&lt; 0.2^\circ$</td>
</tr>
<tr>
<td>11. $P^+ \neq P^-$ and $\theta_L \neq \theta_R$</td>
<td>$\Delta P &lt; 0.1$</td>
<td>$3 \times 10^{-4}$</td>
<td>$&lt; 10^{-5}$</td>
</tr>
<tr>
<td></td>
<td>$\Delta \theta &lt; 0.2^\circ$</td>
<td></td>
<td></td>
</tr>
<tr>
<td>12. Neutron efficiency stability</td>
<td>0.5% gain shift</td>
<td>$&lt; 2 \times 10^{-4}$</td>
<td>$&lt; 10^{-4}$</td>
</tr>
<tr>
<td>13. Quasifree background correction</td>
<td>1% background</td>
<td>$2 \times 10^{-4}$</td>
<td>$9 \times 10^{-4}$</td>
</tr>
<tr>
<td></td>
<td>30% uncertainty</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

The proton beam polarization was cycled between up, down and unpolarized states at the ion source and was measured continuously. Unpolarized beam was used to monitor instrumental asymmetries. The proton polarization ranged from a low of about 0.6 to a high of about 0.7. Its $y$ polarization component was precessed by $90^\circ$ with a superconducting solenoid into the $-x$ direction for production of the neutron beam.

It was necessary to maintain a constant proton beam energy between measurements of $A_n$ and $A_p$ because of the energy dependence of the $np$ c.m. zero-crossing angle; $d\theta_o/dE_n = 0.018^\circ/\text{MeV(lab)}$ at 477 MeV. The proton energy stability was monitored with a beam energy monitor (BEM)[20] which consisted of range counter assemblies located directly behind the forward detectors of the proton polarimeter. The left and right energy values were averaged to reduce kinematic effects associated with possible lateral movement of the beam on the BEM / polarimeter target. The monitor allowed the average beam energy to be checked with an accuracy of $\pm 0.1 \text{ MeV}$ in only 10 minutes of data taking. The largest spread in energy (i.e. width of the stopping distribution) during a two–hour run observed during the experiment was 0.31 MeV FWHM; the largest observed deviation of a single energy determination (two hour period) from
the global mean was 0.28 MeV. The average beam energies for the four major data taking periods were consistent within 0.20 MeV. The effect on $\Delta A$ was $\sim 10^{-5}$ so no corrections to the data were necessary.

B. Secondary Neutron Beam

The neutron beam was produced via the quasi-free $^2H(\bar{p},n)2p$ reaction. Fluctuations in the neutron beam energy due to LD$_2$ temperature variations were below ±30 keV. The neutron beam energy was measured to be 477 ± 2 MeV in a separate experiment. The width of the neutron energy distribution was estimated to be about 11–15 MeV FWHM$^{20,21}$ with a low energy tail from events for which the second final state proton is no longer simply a spectator. This tail decreases rapidly in intensity below the quasi—elastic peak to a level of a few percent of the peak intensity. The variation of the neutron beam polarization with neutron energy due to the quasi-free production mechanism when compounded with the small energy dependence of the $np$ analyzing power has an effect of about $10^{-4}$ on $\Delta A$.

The $x'$ and $y'$ neutron beam intensity profiles were determined by using two delay line wire chambers (DLCs) to reconstruct the origins of $p(n,p)n$ protons recoiling from a converter scintillator. The neutron beam had a uniformly illuminated area (<10% variation) of 56 mm by 40 mm at the FST which covered the entire volume of the target cell. The neutron beam shape remained stable throughout the experiment and the centroids were constant within approximately ±1 mm. The stability of the incident neutron beam direction was estimated to be ±0.005° which affects $\Delta A$ only at the $10^{-5}$ level.

Two dipole magnets (fields along the $+y'$ and $+x'$ directions) precessed the large $x'-z'$ component of neutron polarization into the vertical direction and the small $y'$ component (which does not change sign with the primary proton beam polarization) parallel to the neutron beam direction. A neutron beam polarimeter located downstream of the FST measured the $x'$ and $y'$ polarization components. The average $P_{y'}$ was 0.50 and ranged between 0.47 and 0.55. No evidence for any transverse ($x'$) component of the neutron beam polarization larger than ±0.01 was found during the experiment.

Parity conservation forbids contributions to a left–right asymmetry from longitudinal polarization components in scattering from an unpolarized target or one polarized normal to the scattering plane. The spin correlation parameters $A_{oozx}$ and $A_{ooxz}$ can influence the $A_p$ measurements if there are polarization components of the neutron beam and the FST in the $x'-z'$ plane. The latter spin correlation parameters are very small $^{23–24}$ ($A_{oozx}$ and $A_{ooxz}$ ≈ 0.03 - 0.04 at 477 MeV and 70° in the c.m.). When the uncertainties in the precession of the neutron polarization component due to $a_0y'00$ and possible target polarization components in the scattering plane are combined, the effects of $A_{oozx}$ and $A_{ooxz}$ influence the $\Delta A$ measurement by $<10^{-4}$.

C. The Frozen Spin Target

The frozen spin polarized hydrogen target (FST)$^{15–17}$ had a 4.0 cm diameter, 4.5 cm high cylindrical target cell filled with butanol beads. The target volume and the location of the target canister was determined from X-ray radiographs taken before and after each major data-
taking period. The position was reproducible to better than 2 mm. Material above and below the inner target canister was kept to a minimum to reduce possible background sources. The neutron beam illuminated the entire inner target canister but not the sides of the outer aluminum vacuum shell.

The holding field of 0.257 T was produced by lowering the superconducting solenoid out of the path of the neutron beam and energizing a conventional solenoid above the target cell. The superposition of the fields from the two solenoids provided a vertical field (within ±1.0°), reproducible and stable to ±0.5 mT over a typical three week running period. $A_n$ and $A_p$ measurements with the holding field vertically up and down were averaged to reduce the possibility of systematic errors. The proton target polarization could be parallel or anti-parallel to the holding field direction and all permutations of field and polarization direction were used.

Target proton polarizations were measured with a nuclear magnetic resonance (NMR) system. The average target polarization was ~0.77, varying from 0.90 to 0.50 for individual runs. Comparing the slopes of the $A_n$ and $A_p$ analyzing power data indicated that the target polarizations were free of major systematic errors at the 4% level.

D. The Proton Detection Apparatus

A schematic diagram of the experimental detection apparatus is shown in Fig. 3. The proton detection assemblies were mounted on rigid booms located at 52.00° ± 0.02° to the neutron beam axis. Each proton detection telescope subtended a laboratory angle range of approximately 11° straddling $\theta_0$. Each boom supported four 58 cm by 58 cm (active area) DLCs for proton track reconstruction, time-of-flight (TOF) counters for momentum determination, and a range counter assembly to help separate elastic np recoil protons from high and low energy background.

The range counter assembly consisted of a $\Delta E$ counter, a wedge–shaped brass energy degrader which reduced the energies of elastically scattered protons to a common value independent of their recoil angle, the E-counter, a thin brass absorber and a veto counter. The proton TOF was measured between a small scintillator "start" counter (pTOF) near the FST and a larger "stop" counter (E-counter) 3.4 m away from the FST. Corrections to the calculated proton energies were also made for energy losses in the FST and the wedge–shaped absorber.

E. Neutron Detection Apparatus

The neutron detection assemblies were centered at 32.00° ± 0.02° relative to the neutron beam axis and subtended an 11.4° angular range. Each neutron detector included two stacks of seven optically isolated NE110 scintillator bars 1.05 m wide, 0.15 m deep and 0.15 m high. Three overlapping scintillation counters were placed in front of the neutron array in order to veto charged particles. Small 70 mm by 64 mm "button" counters were located behind the center of each scintillator bar of the back array.

The neutron array pulse height (PHT) signals for the passing protons which penetrated to the button counters were used to calibrate the gain and timing characteristics of the scintillator array PMTs on a run by run basis. A counter (nTOF) similar to the pTOF was used to select
charged particles coming from the FST region. The PHT centroids were maintained at their nominal values by periodic adjustments of the high voltages. Drifts in gain between consecutive two hour periods for each of the 56 PMTs could be determined to within about ±0.3%. The centroids of the button proton positions were determined to within ±0.4 mm during each two hour data taking run which allowed corrections for timing drifts to be made.

Figure 3.

Schematic diagram of the layout of the experimental detection apparatus showing the positions of the neutron and proton booms relative to the frozen spin target.

The hardware discrimination threshold was in a region where the PHT distribution is changing rapidly. In order to check for possible systematic errors associated with gain changes in the neutron array PMTs or variations in the hardware discrimination thresholds, asymmetry zero-crossing angles were calculated as a function of several software cuts on the PHTs. No systematic effects were found even without software discrimination (ADC channel zero as the lower limit).
Uncertainties in the $y'$ and $z'$ neutron interaction coordinates both were limited to ±7.5 cm by the 15 cm height and depth of the neutron scintillator bars. The position resolution along the bars was estimated to be 30 mm FWHM from the difference of the observed positions of individual button events in the front and back bars.

Neutron time–of–flight information was determined from the time of the scattering at the FST calculated using $p$TOF information and the average of the timing signals from the struck neutron scintillator bar. The effects of timing drifts on the extracted particle energies were monitored using $np$ elastic events. The deviations of the proton and neutron energies from the values expected by kinematics (derived from the angle information) were determined on a run by run basis. Timing corrections were made to keep the centroids of the $np$ elastic neutron energy deviations for each bar smaller than 2 MeV. All these considerations were to minimize possibilities of systematic differences between the $A_n$ and $A_p$ data.

F. Selection of Neutron–Proton Elastic Events

For $np$ elastic scattering in the laboratory, assuming the incident energy to be known, only two kinematical variables (including one azimuthal angle $\phi_p$ or $\phi_n$) need to be observed to calculate all the parameters of an event. Since all of the following observables were measured: $T_p$, $T_n$, $\theta_p$, $\theta_n$, $\phi_p$ and $\phi_n$, each event was four times kinematically over–determined. This over-determination allowed $np$ elastic events to be distinguished from inelastic background. $\chi^2$ variables were defined for the following kinematic parameters: the sum of the neutron and proton kinetic energies, $T_{sum} = T_p + T_n$; the $x'$ component of the sum of the particle momenta, $P_{x'} = (p_p + p_n)x'$; the $np$ opening angle, $\theta_S = \theta_p + \theta_n$; and the $np$ azimuthal coplanarity, $\Delta \phi = \phi_n - \phi_p - 180^\circ$.

A typical $T_{sum}$ histogram for elastic events is illustrated in Fig. 4(a). The 14 MeV FWHM spread in the incident neutron energy distribution, the spread of proton energy losses in the FST, and the resolution of the scattered neutron TOF were the main contributors to uncertainties in $T_{sum}$. The tail below the peak is due to a low energy component in the incident neutron energy distribution and the $(n,np)$ background not rejected by cuts on the other kinematic parameters. An example of $P_{x'}$ is illustrated in Fig. 4(b). For all elastic scattering, the total transverse momentum of the scattered $np$ pair should be zero. Constraints on $P_{x'}$ eliminate only $(n,np)$ events. The $np$ opening angle error, $\Delta \theta_S$ is illustrated in Fig. 4(c). The width of the elastic scattering distribution is due primarily to Coulomb multiple scattering of the protons in the FST. Constraints based on $\Delta \theta_S$ eliminated $(n,np)$ events, double scattered neutrons and elastic scattering events originating from the low energy neutrons in the beam. The coplanarity is shown in Fig. 4(d). The width of $\Delta \phi$ is dominated by proton multiple scattering and the size of the neutron scintillator bars. Only $(n,np)$ events and double scattered neutrons are eliminated using this observable. Based on the width of the coplanarity distribution, a standard deviation of about 0.75° was expected for $\theta_S$. The observed width of 0.85° is consistent with an effective energy spread in the neutron beam of about 10 – 20 MeV.

Assuming a nominal neutron beam energy of 477 MeV, the statistical quality of each event was evaluated by defining $\chi^2$ variables for the four observables presented in Fig. 4. Several different combinations of constraints on the $\chi^2$ variables were used in the analysis to verify that the
Figure 4.
Kinematic variables used for $\chi^2$ tests. Events require $\chi^2_1 < 5$ for the other kinematic parameters. The solid curves are gaussian fits. (a) $np$ kinetic energy sum. (b) $np$ horizontal momentum balance. (c) Error in $np$ opening angle. (d) $np$ coplanarity.
The experimental result was not dependent on the value and/or type of cut. Alternatively, constraints were applied to the sum of the four $\chi^2$ variables. The $\chi^2$ cuts were selected to accept a large percentage of $np$ events while reducing quasi-elastic background to a level where it had minimal effect on the asymmetry calculation. No significant dependence on the details of the $\chi^2$ cut was found. This is consistent with the effect of background contamination being small, in agreement with other considerations that will be described later.

G. Asymmetry Calculation

The principle of extracting $\Delta A = A_n - A_p$ from the difference between the zero-crossing angles is illustrated in Fig. 1. The constant of proportionality is the c.m. slope of the analyzing power

$$A_n - A_p = -\Delta \theta \frac{dA}{d\theta}$$

Events were sorted into histograms with 0.25° (lab) wide bins for $\theta_p$ and $\theta_n$, the left–right scattering asymmetries calculated on a bin–by–bin basis, and the slopes of the asymmetries and the zero–crossing angles determined with a linear least squares fit. Simulations showed that no systematic error is introduced by ignoring deviations from linearity in this limited angle range. The results were expressed in terms of the equivalent c.m. neutron scattering angle. All major runs were analysed separately to avoid systematic errors. Separate fits (a total of 48) were made for $A_n$ and $A_p$ data for each FST field direction before the results could be combined. The distribution of confidence levels for the fits was flat so their quality was considered satisfactory.

The analyzing power zero–crossing angle can be determined from the scattering asymmetry distribution as a function of either $\theta_n$ or $\theta_p$. The analysis concentrated on the proton angle result because measurement of neutron angles is susceptible to a number of additional systematic and random errors (discussed later). However, the agreement was satisfactory.

The holding field caused a deflection of the recoil protons by approximately 1.2°. As a result the angle range observed for left and right events was shifted by $\sim 2.5°$. The common angle range is referred to as the "overlap" region. Two different procedures to calculate the scattering asymmetries were possible for the overlap region. The mean difference in $\Delta \theta$ for the two methods was only 0.004°. There was no indication of any uncorrected systematic effect that affected the left and right events differently.

H. Background Subtraction

The most serious potential source of error in the measurement of $\Delta A$ arises from the presence of $(n, np)$ events from the non–hydrogenous materials in the path of the neutron beam. Inelastic processes influence the $A_n$ measurement if the quasi–elastic scattering processes, for example $^{12}C(n, np)$, $^{3,4}He(n, np)$, and $^{16}O(n, np)$, have non–zero asymmetries at the elastic scattering zero–crossing angle. An unpolarized neutron beam impinging on unpolarized target nuclei does not produce any left–right scattering asymmetry. Therefore background events only decrease the slope of the observed $\varepsilon_p = A_pP_T$ asymmetry distribution (appearing to reduce the effective FST polarization, $P_T$), but do not affect the actual determination of $\theta_0$. Estimates of background effects were made by replacing the butanol target beads with graphite beads during the
final data taking period to create a dummy target. This removed most of the hydrogen from within
the target but retained the relative mass ratios of the various background sources. It was assumed
that the effects of carbon and oxygen in the FST are equivalent in producing background events.

The number of \((n,np)\) events from target material and support structure was minimized
by the design of the apparatus. The sides of the FST vacuum walls were outside the limits of the
collimated neutron beam. The acceptance of the pTOF – E counter combination excluded that part
of the vacuum wall intercepted by the beam when entering or exiting the target. Material
immediately above and below the target container was also minimized to reduce background.
Constraints on kinematic observables and the FST target region further reduced quasi–elastic
scattering contamination.

The background changes the true \(A_n\) such that

\[
A_n(\text{true}) = \frac{A_n(\text{measured}) - r_b A_b(\text{measured})}{1 - r_b},
\]

where \(r_b\) is the ratio of background rate to the total observed event rate and \(A_b = \varepsilon_\nu / P_{\text{Beam}}\) is the
background effective analyzing power with a polarized neutron beam. \(P_{\text{Beam}}\) is the neutron beam
polarization. The factor \(r_b A_b\) should remain < \(10^{-3}\) if corrected background effects on \(\Delta A\) are to
remain negligible compared to the experimental statistical accuracy.

It was impossible to determine reliably the contamination fraction \(r_b\) from the butanol
data alone. Even at large \(\chi^2\), background events cannot be distinguished unambiguously from \(np\)
elastic scattering events. Non–Gaussian proton multiple scattering, neutron rescattering and the
low energy tail in the incident neutron energy distribution cause the observed event \(\chi^2\) distributions
to deviate from the ideal case.

Ideally, in the background run all hydrogen should have been removed from the FST so
as to observe only the contribution from inelastic events. Unfortunately small amounts of hydrogen
remained within the graphite target and \(np\) elastic scattering events dominate the data sample
because of the large \(np\) elastic scattering differential cross section. The primary sources of
hydrogen in the graphite target were a resistor located in the base of the target canister (used as a
thermometer), and the mylar superinsulation around the target walls. The epoxy resistor contained
approximately 11 mg of hydrogen. To eliminate the majority of the \(np\) elastic events, only events
originating from the top half of the target were used to estimate the background.

The most accurate estimates of the background fraction were obtained by scaling the
background in the graphite target run to correct for the differences in the amount of non-
hydrogenous material in the graphite target and the butanol filled target, and to normalize the
relative integrated neutron beam fluxes. Estimates were made separately for both holding field
directions and for both \(left\) and \(right\) events. The estimated contamination of inelastic events in
the butanol target data for the \(\chi^2 < 5\) constraint was found to be 0.010±0.003.

The graphite target asymmetries were evaluated using the same \(\chi^2\) conditions as for the
butanol target data. There was no significant dependence on the conditions chosen and the
background asymmetry was consistent with zero within the experimental accuracy. The correction
for the background asymmetry was obtained from the \(\chi^2 < 5\) data.
Using the definition of $\Delta \theta$ and the fact that $A_{n}(np\ \text{elastic})$ in eq. 6 is zero by choice at $\theta_{op}(A_{n})$, the corrected $\Delta \theta_{p}$ for $np$ elastic scattering can be written

$$\Delta \theta_{p}(\text{true}) = \Delta \theta_{p}(\text{measured}) + \frac{r_{b}A_{b}}{(1 - r_{b})\frac{dA}{d\theta}},$$

where $\Delta \theta_{p}(\text{measured})$ is the effect with the butanol filled target, $A_{b}$ is the measured analyzing power for the graphite target data at the zero-crossing angle and $r_{b}$ is the ratio of background to detected events satisfying all cuts applied to the $np$ elastic scattering data.

At the c.m. zero-crossing angle $\theta_{0}(A_{n})$, the proton angle based analyzing power for the $(n,np)$ background was found to be $\epsilon_{b}(\theta_{op}) = 0.004 \pm 0.085$. This results in a correction of $\Delta \theta_{p}(\text{measured}) = 0.340^\circ \pm 0.147^\circ$ to $\Delta \theta_{p}(\text{true}) = 0.340^\circ \pm 0.162^\circ$.

The systematic errors arising from the analysis are summarized in Table II. Geometrical biases caused by the apparatus do not affect $\Delta \theta$ because the effects cancel for the CSB difference $\theta_{o}(A_{n}) - \theta_{o}(A_{p})$ provided that the biases do not change between the $A_{n}$ and $A_{p}$ parts of the experiment. Although a systematic error entry for the background correction is included separately in the table, contributions from this source are also included implicitly in the tests associated with the $\chi^{2}$ cuts since the background fraction changes by about a factor of five for the conditions chosen.

Systematic error estimates were determined by using the sample standard deviation for $\Delta \theta_{p}$ for the groups of individual conditions investigated. The total systematic error estimate is obtained by combining the individual conditions in quadrature. This is realistic because there is no evidence of systematic trends when the parameters are varied, and all fluctuations were of a statistical nature. "Theoretical" estimates of contributions to systematic errors in $\Delta A$ obtained from the observed beam and target properties, and geometrical imperfections are all much smaller than the statistical error and are not included.

As a final check against errors, two independent analyses were made and compared extensively for all but the final stages of extracting $\Delta A$. The average disagreement was only $1/3$ of the statistical error. This is considered to be very reasonable considering the known differences in procedures and the resulting event samples.

III. CONCLUSIONS

The background corrected experimental values of $\Delta \theta$ and $\Delta A = A_{n} - A_{p}$ at the analyzing power zero–crossing angle were determined from the observed values of $\theta_{o}$:

$$\Delta \theta = \theta_{o}(A_{n}) - \theta_{o}(A_{p}) = +0.340^\circ \pm 0.162^\circ \pm (0.058^\circ)$$

for the proton scattering angles. The first error quoted is the statistical uncertainty and the second is the estimate of systematic uncertainty.

The slope of the analyzing power at the zero–crossing angle was determined from a linear fit to various phase shift predictions for a $\pm 1^\circ$ region about the predicted crossover angle. The Saclay–Geneva analysis[23] gives $\frac{dA}{d\theta} = -0.01382$ deg$^{-1}$ and a crossover angle of 69.44$^\circ$ while the Arndt et al.[24] phase shift analysis (SP88) gives $\frac{dA}{d\theta} = -0.01254$ deg$^{-1}$ and a crossover angle of 71.69$^\circ$. 
Table II
Summary of the experimental $\Delta\theta_p$ and $\Delta\theta_n$ systematic errors. The error estimate is obtained from the sample standard deviation of all different tests for that item.

<table>
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<th>Item</th>
<th>$\Delta\theta_p$</th>
<th>$\Delta\theta_n$</th>
</tr>
</thead>
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<tr>
<td>Fitting procedures</td>
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<td>0.009°</td>
</tr>
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<td>$\chi^2$ constraints</td>
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<td>0.028°</td>
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<tr>
<td>Neutron Scintillator Bar Discrimination Threshold</td>
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<td>0.033°</td>
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<td>Background Subtraction</td>
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<td>0.020°</td>
</tr>
<tr>
<td>Independent analyses</td>
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<td>0.050°</td>
</tr>
<tr>
<td>Neutron Scintillator Bar TDC Length</td>
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<td>0.010°</td>
</tr>
<tr>
<td>Fluctuations in Neutron Scintillator Bar Timing</td>
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<td>0.040°</td>
</tr>
<tr>
<td>TDC non-linearities</td>
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</tr>
<tr>
<td>Total a)</td>
<td>0.058°</td>
<td>0.100°</td>
</tr>
</tbody>
</table>

a) Summed in quadrature

New measurements at TRIUMF at 425 MeV indicate that the Saclay–Geneva phase shift analysis gives a slightly better prediction for the analyzing power. This is also confirmed by the better agreement between our observed zero-crossing angle and analyzing power slopes and that prediction. Therefore the Saclay–Geneva value was used to obtain the final results for $A_A$. Note however, that at 188 MeV, the energy of the IUCF experiment, the Saclay–Geneva predictions do not appear to give a reasonable value of the zero-crossing angle.

Using the Saclay–Geneva value, $\frac{dA}{d\theta} = -0.01382 \text{ deg}^{-1}$, $A_n - A_p$ at $\theta = \theta_0$ is

$$\Delta A(\theta_0) = +0.0047 \pm 0.0022 (\pm 0.0008),$$

(9)

based on the laboratory proton scattering angles. A scale error of about 10% for the present uncertainties in the phase shift solutions is appropriate. This should diminish considerably in the next iteration of fitting.

CSB theoretical calculations are discussed in other contributions to this workshop. In brief, the model dependence of the strong distorting waves is small for calculations based on one boson exchange. There is one exception to this: Beyer and Williams[11] have shown that the coordinate space Bonn potential prediction does not agree with the present experiment. Since the primary differences in the models is in the short range behavior, there would appear to be some deficiency at short range in the coordinate space representation of the Bonn potential. Beyer and Williams also show that the Virginia–Mainz hybrid quark–meson potential yields reasonable predictions. The contributions of the dominant CSB mechanisms in a system of six valence quarks appear to be negligible[6,9]. Typical predictions are compared with the experimental result at 477 MeV in Fig. 5.

The experimental result at 477 MeV is the first measurement of a class IV charge symmetry breaking effect in the free $NN$ interaction. The value observed confirms the extension of the conventional OBE potentials to the previously unmeasured class IV $NN$ interaction. Work
is currently underway to measure the CSB parameter $\Delta A = A_n - A_p$ at lower incident neutron beam energies, 350 MeV at TRIUMF[25] and 188 MeV at IUCF[26]. Measurements which determine $\Delta A$ to $\pm 0.0005$ are required to extract information regarding the individual terms contributing to the charge symmetry breaking interaction. The $\rho - \omega$ term is of particular interest, since it relates directly to the OBE description of the $NN$ interaction at short distances.

Figure 5.
Comparison of the 477 MeV TRIUMF result with the calculations of Refs. 6 and 8.

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References

CALCULATIONS OF CHARGE-SYMMETRY BREAKING IN n-p ELASTIC SCATTERING

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ABSTRACT

Charge-symmetry breaking of nuclear forces can be observed in neutron-proton elastic scattering. The major contributions arise from the neutron-proton mass difference in one-pion and one-rho exchanges, from the neutron anomalous magnetic moment in one-photon exchange, and from rho-omega meson mixing. Predictions are compared for a number of different models and are found to agree well with both the existing TRIUMF measurement at 477 MeV and with the new IUCF measurement first reported at this meeting.

INTRODUCTION

Isospin invariance has proved to be a powerful tool for developing an understanding of the strong interactions, even though from the beginning it was recognized as only an approximate symmetry. As with other symmetries, an enormous amount can be learned from testing our understanding of symmetry violations against experimental measurements. We are primarily concerned here with low and intermediate energy phenomena, where hadrons are the appropriate basis for a description of the strong interactions. At the quark level isospin invariance is broken by the up-down quark masses and by electromagnetic and weak interactions. At the hadron level obvious sources of isospin violation are: electromagnetic (em) and weak effects, mass differences within isospin multiplets (e.g., $M_p \neq M_n$ and $m_{\pi^\pm} \neq m_{\pi^0}$), and mixing of mesons with different isospin (e.g., $\pi^0 \leftrightarrow \eta$ and $\rho^0 \leftrightarrow \omega$). Charge-independence breaking (CIB) is another way of saying isospin violation. Charge-symmetry breaking (CSB) is a particular kind of CIB where a system fails to be invariant under reflections in the $T_3 = 0$ plane, ($T \equiv$ total isospin, $T_3 \equiv$ 3rd component).

$N$-$N$ forces can be divided into four classes according to their isospin invariance, where only class III and IV potentials give rise to CSB. In the $np$ system only class IV forces can contribute to CSB. The simplest such forces are

\begin{align*}
V^{IV} &= (r_1 \times r_2) \cdot (\sigma_1 - \sigma_2) \cdot L\nu (r), \\
V^{IV} &= (r_1 \times r_2) \cdot (\sigma_1 \times \sigma_2) \cdot L\nu (r),
\end{align*}

where $L$ is the orbital angular momentum operator in the center-of-mass frame and $r = r_1 - r_2$ is the internucleon separation. Here we are concerned with the phenomenon of $n$-$p$ elastic scattering, which is of particular interest since the absence of Coulomb forces makes it possible to establish the existence of CSB unambiguously. CSB in the $np$ system requires mixing of $T = 0$ and $T = 1$ states and, hence, also of $S = 0$ and $S = 1$ states. Then only $J^L$ states can contribute to CSB in $n$-$p$ elastic scattering and hence to $A_A$. The calculations are performed using the distorted-wave Born approximation (DWBA), where

\begin{equation}
V = V_{CS} + V^{IV}
\end{equation}
and
\[ \delta T_{fi} = \int d^3 r \psi_f^{(-)^\dagger}(r)V^{IV}(r)\psi_i^{(+)}(r). \] (3)

The initial and final distorted waves (\(\psi_i\) and \(\psi_f\)) are calculated from some charge-symmetric model of the N-N interaction (V\(_{CS}\)).

A recent TRIUMF experiment\(^2\) has found evidence of CSB in this system by measuring a difference between neutron and proton analyzing powers in n-p elastic scattering at 477 MeV at the zero-crossing point of the analyzing power. Their nonzero result is a direct measurement of CSB, since it shows that the system is not invariant under the transformation \(n \leftrightarrow p\). The difference in \(n\) and \(p\) analyzing powers is \(\Delta A(\theta) = A_n(\theta) - A_p(\theta)\), where \(A_n(\theta)\) and \(A_p(\theta)\) are the neutron and proton analyzing powers, respectively. Denoting the zero-crossing angle of the analyzing power in the center-of-mass (c.m.) system by \(\theta_{c.m.}\) we define \(\Delta A \equiv \Delta A(\theta_{c.m.})\). The result at 477 MeV is \(\Delta A = [47 \pm 22(\pm 8)] \times 10^{-4}\). An experiment at IUCF has just been finished and the preliminary results\(^3\) have first been announced at this meeting. The result reported was that the average \(\Delta A(\theta)\) between 80° and 100° was \(32 \pm 6(\pm 67)\) \times 10\(^{-4}\).

CSB CONTRIBUTIONS

The bar phase-shift \(L-S\) representation of the \(N-N\) scattering matrix\(^4\) can be extended to include a new parameter \(\gamma_J\) (the spin singlet-triplet mixing angle). This can then be related to the CSB amplitude

\[ f(k, \theta) = \frac{i}{2k} \sum_{J=1}^{\infty} (2J + 1) \sin(2\gamma_J) \exp(i\delta_J + i\bar{\delta}_{JJ})d_{J0}^J(\theta). \] (4)

where the \(d_{J0}^J\) are the Wigner functions and \(\delta_J\) and \(\bar{\delta}_{JJ}\) are the singlet and uncoupled triplet bar phase shifts, respectively. Use of the DWBA and using \(\sin(2\gamma_J) \approx 2\gamma_J\) (since \(\gamma_J\) is very small) allows \(V^{IV}\) to be related to \(f(k, \theta)\) and hence to \(\gamma_J\). This gives

\[ \gamma_J = -2E_T k \sqrt{J(J+1)} \int_0^\infty dr r^2 R_J(r)g(r)R_{JJ}(r), \] (5)

where \(E_T \equiv \text{(total) c.m. energy} \equiv 2E\) with \(E\) the energy per nucleon in the c.m. frame and where

\[ g(r) \equiv v(r) \text{ and } (1)^J w(r) \] (6)

for Eqs. (1a) and (1b), respectively. The distorting effects of the strong interaction are included through the radial wave functions \(R_J(r)\) and \(R_{JJ}(r)\), which are the \(^1P_1\) and \(^3P_1\) partial waves, respectively. The analyzing power difference is given by

\[ \Delta A(\theta) = A_n(\theta) - A_p(\theta) = 2 \text{Re}(b^* f) \sigma_0 \] (7)

where \(\sigma_0\) is the unpolarized cross section and where \(b\) is one of the usual charge-symmetric invariant amplitudes making up the scattering matrix.

The main sources of CSB are shown in Fig. 1 and are: a) e.m. interaction between proton charge and neutron magnetic moment. b) \(p-n\) mass difference in \(\pi\)
and $\rho$ exchanges, c) $\rho\omega$ mixing, and d) and e) are due to the $p-n$ mass difference in two-pion exchange. Estimates of quark effects were negligibly small, as were two-pion effects. So only a) → c) contribute significantly. Mixed $\pi\gamma$ exchanges have not yet been calculated but are expected to be very small and $\pi^0-\eta$ mixing gives no class IV contribution. The contributions from $\gamma$ and mixed $\rho\omega$ exchanges are of the Eq. (1a) type, while $\pi$-exchange has the form of Eq. (1b). Both class IV forces are present in the $\rho$-exchange. It is tedious but otherwise straightforward to derive these CSB potentials from the nonrelativistic reductions of the invariant amplitudes, (see e.g., Refs. 5-8). In recent work $^7,^8$ relativistic corrections were included in the derivation of the potentials (with the exception of the $\rho-\omega$ mixing contribution). For a more complete discussion of these CSB contributions see Refs. 7-8.

![Diagram of charge-symmetry breaking processes](image)

**Fig. 1.** The charge-symmetry breaking processes considered (with the exception of quark effects). Shown are (a) one-photon, (b) one-pion or one-rho, (c) mixed rho-omega, and (d) uncrossed and (e) crossed two-pion exchanges. The crosses indicate the CSB vertex function arising from the neutron-proton mass difference. The cross hatching refers to the usual subtraction procedure in (d). Only (a), (b), and (c) have been found to have significant effects.

Other approaches have also been considered. For example a relativistic calculation $^9$ of the CSB contributions from the $n-p$ mass difference and $\rho-\omega$ mixing has been carried out in the momentum space Bonn potential. These agree well with the results presented here. The effects of $\sigma$, $\omega$, and $A_1$ meson exchanges have also been calculated $^{10}$ and found to be very small. A purely relativistic approach using a covariant representation of the on-shell $N-N$ scattering amplitudes was also attempted $^{11}$, and while it can be shown to be formally equivalent $^{12}$ there is ambiguity due to an arbitrary off-shell extrapolation of the $T$ matrix and an effectively soft pion form factor.$^{10}$

**NUMERICAL RESULTS**

The calculations were carried out in the following way. To minimize theoretical uncertainties empirical phase shifts$^{13}$ were used to determine the charge-symmetric
amplitudes. The distorted waves in Eq. (5) used to calculate the $\gamma_J$ were obtained from the particular model of the $N-N$ interaction being considered. The coupling constants, form factors, and masses used in the CSB potentials (i.e., in $v(r)$ and $w(r)$ of Eq. (6)) were taken from the model where appropriate or alternatively typical values from low-energy analyses were used. The model dependence in these calculations enters only through the $\gamma_J$'s of Eq. (5).

The models considered here are the Reid soft-core (RSC), Paris, coordinate space-Bonn (OBEPR), and Virginia-Mainz (VMZ) potentials. The VMZ potential is a hybrid quark-meson model. For the RSC and Paris potentials, which are not one-boson exchange models, the parameters used in the calculation of $V^{IV}$ were taken from low-energy analyses as previously discussed. For the Bonn (OBEPR) and VMZ potentials the parameters were taken from the models, with the exception of the $\rho\omega$ mixing strength which was taken from a recent analysis of $e^+e^- \rightarrow \pi^+\pi^-$ data,

$$<\rho^0|H|\omega> = -4520 \text{ MeV}^2.$$  \hspace{1cm} (8)

Note that this differs by a factor of $\simeq 1.3$ from the mixing matrix element used in earlier work \cite{7-9}.

$$<\rho^0|H|\omega> = -3400 \text{ MeV}^2.$$ This means that the contributions to $\Delta A$ from $\rho\omega$ are scaled up by $\simeq 1.3$ from these.

Table I. Comparison of the different model predictions of $\Delta A (\times 10^4)$ at the three energies of interest. Also shown are the two experimental results, the laboratory kinetic energy $E_{lab}$, the c.m. angle (in deg) at which the analyzing power goes to zero $\theta_{c.m.}$, and the separate contributions from $\pi$, $\gamma$, $\rho$, $\omega$, and $\rho$ exchanges. The different nucleon-nucleon potentials used are the Reid soft-core (RSC), Paris, coordinate-space Bonn (OBEPR), and Virginia-Mainz (VMZ) potentials. The VMZ potential is a hybrid quark-meson potential and the two versions considered here have a six-quark core radius of $b = 1$ fm and 1.2 fm, respectively.

<table>
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<th>Potential</th>
<th>$\pi$</th>
<th>$\gamma$</th>
<th>$\rho\omega$</th>
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</table>

Note: $45 \pm 22(\pm 8)$
The results are shown in Table I. We see that the predictions are in good agreement with each other and with the two experimental results with one exception. They also agree well with the relativistic Bonn model calculations of Ref. 9. The exception is the coordinate space Bonn potential (OBEPR) from which it seems reasonable to conclude that there is some deficiency in its description of the short-range behavior (see Ref. 8 for more details). In Fig. 2 and 3 the angular contributions to $\Delta A(\theta)$ from the one-pion, one-rho, e.m., and mixed $\rho$-$\omega$ exchanges are shown at 188 MeV and 477 MeV respectively.

![Graph](image_url)

Fig. 2. The contributions to $\Delta A(\theta)$ at 188 MeV from the $\pi$, $\gamma$, $\rho$-$\omega$, and $\rho$ exchanges. The curve for $\rho\omega$ should be scaled up by $\approx 1.3$ for $<\rho^0|H|\omega> = -4520$ MeV$^2$, (c.f., $-3400$ MeV$^2$).

![Graph](image_url)

Fig. 3. The contributions to $\Delta A(\theta)$ at 477 MeV. As in Fig. 2 the $\rho\omega$ curve should be scaled up by $\approx 1.3$. 
CONCLUSIONS

While there is naturally some small model-dependence of the predictions the agreement with the two existing experimental measurements is quite good. Early discrepancies between different theoretical studies have now been satisfactorily resolved. It is satisfying that the 188 MeV prediction was made 2 years before the experimental result from IUCF was known. It is reasonable to conclude at this time that the meson exchange picture is adequate to explain the data. Clearly the meson exchange picture must give way to the quark picture at sufficiently short distances, however the data have not yet forced us to invoke explicit quark degrees of freedom. On the other hand it is interesting that the hybrid quark-meson VMZ model also agrees well with the data.

The long-term theoretical goal must be to find some quark-level description of CSB and the $N-N$ force which neatly marries on to the meson exchange picture at intermediate and long distances. In the meantime new experimental data provide a powerful stimulus to and test of such theories.

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CHARGE SYMMETRY BREAKING EFFECTS
IN NEUTRON-PROTON SCATTERING
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ABSTRACT

In most theoretical models nearly all of the charge symmetry breaking observable \( \Delta A = A_n - A_p \) measured recently at TRIUMF and IUCF at the zero crossing angle of the average analyzing power can be explained as the \( np \) mass difference effect in pion exchange. More interesting heavy meson exchanges contribute to the angular distribution mainly outside that angle. A moderate contribution arises from two pion exchanges involving \( N\Delta \) intermediate states.

INTRODUCTION

Symmetries have had a twofold role in physics. Firstly, they lead to tremendous simplifications in otherwise formidable problems, secondly, they and their breakings give revealing glimpses and insight at the interactions. Both ways studies of symmetries offer great opportunity in tackling basic problems of physics. For example, without any symmetries arbitrary spin-\( \frac{1}{2} \) on spin-\( \frac{1}{2} \) scattering would have 16 amplitudes, and an amplitude analysis would be hopeless. However, the parity and time reversal invariances reduce the number of independent amplitudes to 6. For identical fermions this number is further reduced to 5. If the proton and neutron are regarded as states of the same particle species in an abstract isospin space, then the generalized Pauli principle unambiguously relates the symmetry of the isospin state with that of the spin and space. If these states cannot mix, then for both the antisymmetric \( T = 0 \) and symmetric \( T = 1 \) isospin combination of the two nucleon system there are only 5 independent spin amplitudes. In strong interactions the value of isospin is conserved to good accuracy. The concept of isospin and its conservation have proved to be a very useful simplification in understanding the strong interaction of nucleons.

However, rotations and reflections in the isospin space are not exact symmetries. Of course, the proton and neutron do not behave in the same way under the electromagnetic interaction. Even for the strong interaction the symmetry is not exact, because different charge states of the same particle species (e.g. proton and neutron) can have different masses. This means that neither for an arbitrary rotation of the isospin

\[
[H, \hat{T}] = 0 \quad \text{(charge independence)},
\]

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nor for a rotation by 180°

\[
[H, P_{CS}] = 0, \quad P_{CS} = e^{i\pi T_2} \quad \text{(charge symmetry)},
\]  

are exactly valid. The operator \( P_{CS} \) essentially reverses the "z-axis" of the isospin space, for nucleons it changes protons into neutrons and vice versa. At the hadronic level the reasons for isospin symmetry breaking can be traced to mass differences between different charge states of mesons and baryons, to electromagnetic interaction effects and to the fact that mesons of different isospins mix \((\pi^0 \text{ and } \eta \text{ or } \rho \text{ and } \omega)\). Earlier also the mass differences and meson mixings could have been argued to be of electromagnetic origin. However, in the quark model the \( u \) and \( d \) quark mass difference can be fundamentally independent of the electromagnetic interaction, so isospin symmetry breaking may well occur even in the pure strong interaction. At the quark level an obvious reason for isospin symmetry breaking is just this \( u \) and \( d \) quark mass difference as well as the Coulomb interaction between quarks.

One reason for the recent interest in isospin breaking, in particular charge symmetry breaking (CSB), is the hope that its detailed understanding could tell something about the underlying hadron structure in terms of quarks or solitons etc. One should be able to gain information about the strong interaction, which is complementary to the isospin symmetric case. For example, the meson exchange model of nuclear forces can be tested in a new environment where it was not originally fitted. At the very least new constraints on e.g. meson-nucleon couplings could be expected, since meson exchanges appear in different combinations. On the other hand, deviations from the meson exchange picture could be signatures of a deeper mechanism. A particularly interesting candidate for a "smoking gun" is the spectacularly great theoretical difference between the \( \rho \omega \) mixing effect and CSB quark-gluon calculations.

In the wake of the recent TRIUMF experiment on CSB in \( np \) scattering much of the recent work concentrates on spin observables in this scattering and in this talk I shall concentrate on the \( np \) system. For earlier and more comprehensive reviews see Refs. 1 and 2.

**CSB OBSERVABLES AND AMPLITUDES**

The most transparently CSB interaction is of the form (class III in the terminology of Henley and Miller)

\[
V_{CSB}^{III} = (\vec{\tau}_1 + \vec{\tau}_2)_3 \, v(r).
\]

This changes sign in the inversion of the \( \tau_3 \) axis. Here \( v(r) \) has a similar spin-space structure as the isospin symmetric \( NN \) interaction. This charge dependence causes a difference between the \( pp \) and \( nn \) scattering lengths and affects binding energy differences in mirror nuclei. Due to the presence of the Coulomb interaction extracting reliable information on the strong interaction is quite difficult and model dependent, although considerable progress has been
made. Since class III forces are proportional to the total isospin operator, they conserve the isospin and also vanish in the neutron-proton interaction ($T_3 = 0$).

In the $np$ system the nonzero CSB interaction compatible with space inversion and time reversal invariances and with the generalized Pauli principle are class IV forces of the forms

$$V_{\text{CSB}}^{\text{IV}} = (\vec{\tau}_1 \times \vec{\tau}_2)_3 \vec{\sigma}_1 \times \vec{\sigma}_2 \cdot \vec{L} v(r),$$

$$V_{\text{CSB}}^{\text{IV}} = (\vec{\tau}_1 - \vec{\tau}_2)_3 (\vec{\sigma}_1 - \vec{\sigma}_2) \cdot \vec{L} w(r).$$

Both of these necessarily change the spin and isospin. Therefore they must vanish for $pp$ and $nn$ states and cannot directly affect e.g. mirror nucleus energy differences. It is also clear that the first form requires an exchange of a charged meson (notably $\pi^\pm$), whereas the second needs a neutral one (most importantly a photon or $\rho\omega$ meson mixing). In partial waves the difference between these two forms is a phase factor $(-1)^J$. This causes the CSB observables from charged and neutral meson exchanges to be qualitatively different. Because of spin mixing and conservation of parity and angular momentum, it is immediately clear that these forces can only contribute to the mixing of spin singlets with triplets which are not tensor coupled (i.e. $J = L$). This mixing is now a sixth independent amplitude in $np$ scattering.

Under Lorentz invariance, parity conservation and time reversal invariance, the scattering matrix of two spin-$\frac{1}{2}$ particles can be expressed in terms of six amplitudes as

$$M = \frac{1}{2} [(a + b) + (a - b)\sigma_1n\sigma_2n + (c + d)\sigma_1m\sigma_2m + (c - d)\sigma_1l\sigma_2l + e(\sigma_1n + \sigma_2n) + f(\sigma_1n - \sigma_2n)].$$

The CSB amplitude $f$ mixes the triplet and singlet states, whereas $a...e$ are the usual isospin conserving amplitudes. The directions in eq. (6) are defined in terms of the initial and final momenta as

$$\hat{i} = \frac{\vec{k}_f + \vec{k}_i}{|\vec{k}_f + \vec{k}_i|}, \quad \hat{m} = \frac{\vec{k}_f - \vec{k}_i}{|\vec{k}_f - \vec{k}_i|}, \quad \hat{n} = \frac{\vec{k}_i \times \vec{k}_f}{|\vec{k}_i \times \vec{k}_f|}.$$  

In partial wave presentation $f$ is conventionally expanded in terms of mixing angles $\gamma_J$ similar to the tensor coupling as

$$f = f_{\text{em}}^J + \frac{i}{2k} \sum_{L=1}^{\infty} \frac{2L + 1}{\sqrt{L(L + 1)}} \sin 2\gamma_J e^{i(\delta_J + \delta_{JJ})} P_L^J(\cos \theta).$$

with $P_L^J(\cos \theta) = \sin \theta P_L^J(\cos \theta)$ the associated Legendre polynomial. To make the series converge the long ranged magnetic interaction amplitude $f_{\text{em}}^J$ is normally separated as an

---

1Note that the difference with the $np$ scattering length in the same singlet state is mainly due to the class II isotensor interaction $\propto (3\tau_3\tau_2 - 1)$ arising from meson mass differences, in particular from $m_{\pi^\pm} - m_{\rho}$. Although this is charge dependent, it is invariant in the reversal of the $\tau_3$ axis and therefore charge symmetric.
exact plane wave amplitude. The same can be done with OPE to improve the convergence. The parameter $\gamma_j$ is defined similarly to the tensor coupling by

$$S = \begin{pmatrix} \cos 2\gamma_j e^{2i\delta_j} & i \sin 2\gamma_j e^{i(\delta_j+\delta_{jj})} \\ i \sin 2\gamma_j e^{i(\delta_j+\delta_{jj})} & \cos 2\gamma_j e^{2i\delta_{jj}} \end{pmatrix}.$$  \hspace{1cm} (9)

The isospin violating part is the nondiagonal elements of the S-matrix. These are so called ‘bar’ phase shifts, although the bar is omitted.

The simplest CSB observables are the differences of the neutron and proton analyzing powers in polarized scattering

$$\Delta A = A_n - A_p = 2 \text{Re} (b^* f)/\sigma_o ,$$ \hspace{1cm} (10)

$$A_{ks} - A_{sk} = 2 \text{Im} (e^* f)/\sigma_o .$$ \hspace{1cm} (11)

These are nonzero, if and only if the neutron and proton are distinguishable. Typically they are very small and difficult to obtain experimentally. The recent TRIUMF experiment\textsuperscript{6} measured $\Delta A$ at 477 MeV for the angle 72° at which $A$ crosses through zero, with the result

$$\Delta A(72°) = 0.0047 \pm 0.0022 \pm 0.0008 ,$$ \hspace{1cm} (12)

thereby for the first time establishing the existence of this new class of nuclear force. The restriction to the zero crossover angle is necessary to eliminate the major systematic errors, but as will be seen later, unfortunately misses some most interesting piece of CSB in np scattering. A new measurement at 350 MeV extending somewhat off the zero crossing angle has been proposed at TRIUMF\textsuperscript{7} and another from IUCF at 183 MeV is preliminarily reported in this workshop.\textsuperscript{8}

**THEORY OF CSB np SCATTERING**

Considerable amount of theoretical work has been done to explain or predict the TRIUMF result by nonrelativistic DWBA calculations,\textsuperscript{9-13} a relativistic t-matrix approach,\textsuperscript{14} or quark clusters.\textsuperscript{15} All DWBA calculations agree that at the zero-crossing angle the effect of the np mass difference in the OPE is dominant and enough to explain the data point obtained at TRIUMF. They use some more or less phenomenological potential to produce nucleon wave function distortions, whereupon the isospin breaking amplitude is treated perturbatively. (Ref. 12, in fact, uses an exact coupled channels method, but for the small isospin mixing this is unnecessary luxury.)

The CSB OPE contribution arises in the Dirac formalism in the following way. The pseudoscalar pion-nucleon coupling is essentially

$$\Gamma_{NN\pi} = \left( 1 - \frac{-\vec{p} \cdot \vec{\sigma}}{(E + M)(1 - \delta \tau_3)} \right) \gamma_5 \tau_i \left( 1 - \frac{\vec{p} \cdot \vec{\sigma}}{(E + M)(1 - \delta \tau_3)} \right) \phi_i ,$$ \hspace{1cm} (13)
Fig. 1. $\Delta A$ for two energies. Solid line: full calculation of ref. 12 including all one meson exchange and electromagnetic effects; dash-dot: photon exchange subtracted; dashed: the np mass difference effect in OPE. The difference between the dash-dot and dashed curves is the heavy meson contribution, mostly $\rho\omega$ mixing. The data are from refs. 6 and 8.

if the np mass difference is taken into account and $E \approx M$ assumed. Keeping the terms of first order in $\delta = (M_n - M_p)/(M_n + M_p)$ and using properties of the Pauli matrices one gets the coupling in the form

$$\Gamma_{NN\pi} \approx -\frac{(p' - p) \cdot \vec{\sigma}}{E + M} \vec{r} \cdot \vec{\phi} - \frac{(p' - p) \cdot \vec{\sigma}}{E + M} \phi_3 \delta - \frac{1}{E + M} (\vec{p}' + \vec{p}) \cdot \vec{\sigma} \cdot (\vec{r} \times \vec{\phi})_3 \delta .$$

(14)

In pion exchange between two nucleons the second term leads to a class III interaction, the third one combines to give a nonrelativistic CSB potential of type of eq. (4) as

$$V(OPE) = \frac{f^2}{4\pi} \frac{1}{\mu^2} \left\{ \nabla \cdot \vec{\sigma}_1 \ \nabla \cdot \vec{\sigma}_2 \left( \frac{e^{-\mu r}}{r} \right) \left[ \vec{r}_1 \cdot \vec{r}_2 + (\vec{r}_1 + \vec{r}_2)_3 \delta \right] \right. $$

$$-2\delta (\vec{r}_1 \times \vec{r}_2)_3 \ \vec{r}_1 \times \vec{r}_2 \cdot \vec{L} \ \frac{d}{dr} \left( \frac{e^{-\mu r}}{r} \right) \right\} .$$

(15)

For charged mesons isospin violation arises from the noncommutativity of the isospin matrices in the nucleon mass and in the $\vec{r} \cdot \vec{\phi}$, whereas neutral pion exchange contributes only to class III. The effect of CSB OPE is shown by the dashed curve in Fig. 1. Note that the spin-space structure of the class III interaction is the same as the standard isoscalar part and does not give rise to new spin observables as class IV does. In a similar way the mass difference term arises in other meson exchanges. However, the algebra for vector mesons is more complicated, and one has to be careful with the isospin operators. For example, the Gordon decomposition requires a subtle modification for the $\rho$ meson, when the nucleon mass difference is not neglected.\textsuperscript{12}
Fig. 2. The \( np \) mass difference effect in \( \Delta A \) arising from \( \rho, \sigma, \omega \) and \( a_1 \) exchanges at 477 MeV. (Ref. 12)

In general other CSB contributions than OPE are rather small at the zero-crossing angle and except for perhaps \( \rho \) exchange and certain two pion exchange effects can be mostly neglected. It is interesting to note that at intermediate energies the \( \Delta A \) arising from neutral exchanges goes through zero at nearly the same angle as \( A \) itself does. However, as seen in Fig. 1 \( \rho \omega \) mixing and the magnetic interaction of the proton with the neutron spin can be quite as important as OPE outside the zero-crossing point and even at this angle for the lowest energies. The former appears in \( np \) scattering as an isospin breaking potential

\[
V_{\rho\omega} = f_\omega g_\rho - f_\rho g_\omega \frac{1}{2M^2} \frac{\langle \omega|H|\rho \rangle}{m_\omega^2 - m_\rho^2} (\vec{r}_1 - \vec{r}_2)^3 (\vec{\sigma}_1 - \vec{\sigma}_2) \cdot \vec{L} \frac{1}{2} \frac{d}{dr} \left( \frac{e^{-m_\rho r} - e^{-m_\omega r}}{r} \right),
\]

with vector and tensor coupling constants \( g_i \) and \( f_i \). This interaction is of the form (5) and is therefore qualitatively different from OPE. The mixing matrix element \( \langle \omega|H|\rho \rangle = -0.00452 \pm 0.0006 \text{ GeV}^2 \) is well determined\(^\text{19} \) and also the \( \rho NN \) coupling is known from other sources,\(^\text{20} \) so that good data on \( \Delta A(\theta) \) would constrain the \( \omega NN \) coupling independently of usual \( NN \) scattering fits.

The contributions (from ref. 12) shown in Fig. 1 for two energies consist mainly of these three sources. Reference 12 evaluates CSB due to mesons of arbitrary Lorentz and isospin structure and finds other contributions to be small. As seen from Fig. 2, except for the \( np \) mass difference effect in \( \rho \) exchange, they can be safely neglected. All the way through ref. 12 and Figs. 1 and 2 the standard Reid soft core potential\(^\text{17} \) is used for the isospin obeying interaction and the Bonn OBEP\(^\text{18} \) parameters for the CSB one. The agreement with the existing experiments is perfect and it is noteworthy that at the lower energy all the above important mechanisms contribute about equally.

Different isospin conserving "diagonal" potentials, give rather similar results, as long
as they are fitted to $NN$ phase shifts. A survey of this dependence has been made in Ref. 13 and is reviewed by Williams in this workshop. Unfortunately the pure effect of the diagonal potential variations is not extracted using the same CSB interaction, so the origin of the differences found in Ref. 13 is not always easy to trace to either of these components. OPE with its well known coupling constant is an exception, and it is fairly independent of the isospin conserving potential except for the coordinate space representation of the Bonn potential\(^\text{18}\) (OBEP), which gives an unacceptably large $\Delta A$.

As a test of meson exchange models nearly any CSB interaction with a sensible OPE passes the test with flying banners. This is due to the fact that at the zero-crossing angle the badly known heavy meson couplings are unimportant, whereas OPE is well known, so that the prediction should be reliable. Only at the IUCF energy the $\rho\omega$ mixing with poorly known coupling constants becomes appreciable. The IUCF result appears to favour the strong couplings of the $\rho$ and $\omega$ as used e.g. in the Bonn potential. However, to make any strong conclusion, a careful study of two pion and pion-photon exchange contributions is necessary. Ref. 12 shows that even $\Delta A(\text{OPE})$ can be varied by 30\% by varying the dipole form factor within reasonable limits. The results of Fig.1 were obtained for the Bonn cut-off $\Lambda_\pi = 1300$ MeV, but variations are not excluded by experiment. In this sense a more accurate value of $\Delta A$ at higher energies could be a way to constrain the pion-nucleon form factor. The use of a very soft form factor $\Lambda_\pi = 550$ MeV in the off-shell extension of the t-matrix may be one reason why the relativistic calculation of Ref. 14 gives only $\Delta A(\text{OPE}) = 10 \times 10^{-4}$, much lower than the nonrelativistic calculations. However, a nonrelativistic potential calculation with this value of $\Lambda_\pi$ still gives twice as large a value for $\Delta A(\text{OPE})$.\(^\text{12}\) A more likely reason for the deviation of the t-matrix method is that it parametrizes the upper components of the experimental $NN$ scattering matrix in terms of five invariants. However, this is too much of a constraint in the relativistic formalism which requires also an explicit treatment of the lower components.\(^\text{16}\) In conclusion, the TRIUMF and IUCF experiments confirm the theoretical predictions of CSB in the neutron-proton system. However, much more accurate data are necessary to discriminate between most models and to really pin down anything that we don’t already know. Even then $\Delta A$ constrained at the crossover angle seems of limited use beyond establishing class IV CSB in OPE.

A particularly interesting feature in CSB mechanisms is that meson mixing is presumably due to the up and down quark mass difference proportional to $(m_u - m_d)(u\bar{u} - d\bar{d})$. If, however, quark exchanges take over at short distances instead of vector meson exchanges, this kind of term may not arise. The gluon exchange model of Ref. 15 gives a contribution to $\Delta A$ two orders of magnitude smaller than the experimental results (see Fig. 7, the dash-dot curve). So there appears to be a qualitative difference between vector meson exchanges and quark models, which is not the case for isospin respecting forces. Therefore, angular distribution results for $\Delta A$ would be of great interest. Even the slope of $\Delta A$ at the zero-crossing
Fig. 3. Two pion exchange mechanisms for CSB. Any vertex can be CSB. Normal box diagrams with nucleon intermediate states are contained in the initial and final state correlations.

angle should be a useful constraint. Data at low energies also can constrain $\rho\omega$ mixing, but they should be exceedingly accurate. The IUCF result in Fig. 1 with only statistical error appears to require all the above major mechanisms to contribute at 183 MeV indicating also a significant $\rho\omega$ mixing. However, the precaution of still missing hadron level mechanisms should be kept in mind. Also at the present time any quark-gluon models of the $NN$ interaction can at best be considered as qualitative. For example, in the $\pi^0 pp$ and $\pi^0 nn$ coupling constant difference even the sign arising from different quark models seems uncertain.\(^{21}\) Furthermore, it should be noted that the mixing of the pseudoscalar mesons $\pi^0$ and $\eta$ gives rise only to a CSB force of class III and does not directly contribute to the $np$ system.

The extreme forward direction is dominated by the long ranged magnetic interaction which causes a forward divergence. This results in spectacular differences between the neutron and proton analyzing powers for very small angles $\theta \leq 2^\circ.2,5$ Of course, the electromagnetic component of CSB is in principle as well known as OPE.

A largely neglected uncertainty in CSB $np$ scattering is the effect of $N\Delta$ admixture and related two pion exchanges and inelasticities as shown in Fig.3. Similarly to the $\pi NN$ vertex case, a CSB $\pi N\Delta$ coupling is possible due to the mass splitting of the $\Delta$'s. This is necessarily of the charged cross product form

$$\Gamma_{\pi N\Delta} \propto -i \frac{(\vec{p}' + \vec{p}) \cdot \vec{S}}{E + M} (T \times \phi)_{3\delta},$$

with $\vec{S}$ ($\vec{T}$) the transition spin (isospin) operator. This vertex produces a CSB $NN \rightarrow N\Delta$ transition potential similar to the class IV OPE of eq. (4). Also the neutral $\pi NN$ vertex can cause CSB transitions via a potential similar to the class III term in eq. (15) with $\tau$'s replaced by $T$'s.

Reference 22 indicates a non-negligible contribution outside the zero-crossing point above TRIUMF energies. Again at the crossover angle the effect is small. This apparently reflects the fact that the two pion exchange dominantly behaves like a scalar ("$\sigma$ meson"),
which produces a class IV interaction of form (5). This in turn acts much like $\rho\omega$-mixing, as was pointed out earlier. In addition to the mass differences within baryon multiplets ref. 22 uses also $\pi\pi$-mixing as an isospin breaking transition potential $NN \rightarrow N\Delta$. Isobar components are intimately related to pionic inelasticities of the $NN$ system. Fig. 4 shows at 650 MeV the effect of first including inelastic $N\Delta$ components by coupled channels and then switching the inelasticity off. At higher energies these effects are not negligible. The isobar components may have a significant effect even at low energies by changing short range correlations and also by contributing to the overall CSB interaction directly. Since ref. 22 considers only box diagrams (Fig. 3a) neglecting crossed boxes and does not ensure phase equivalence, this effect (the difference between the solid and dashed curves) remains somewhat uncertain. The effect of actual inelasticity is very likely less uncertain, since crossed boxes are not expected to contribute significantly to that. The dotted curve is obtained from the dashed curve by simply making the inelastic imaginary parts of the $\gamma_j$'s to zero. Above the $\Delta$ threshold this effect is nonnegligible outside the zero crossing angle, although it does not cause a qualitative change in $\Delta A$. However, in the mixing parameters the imaginary part is as important as the real part.

To improve over ref. 22 in the treatment of the $\Delta$ contribution to CSB one should include also a residual interaction arising from the crossed boxes (Fig. 3b). This part has been calculated as a two pion exchange $NN$ potential at zero energy to be added to the coupled channels calculation and the results are shown in Figs. 5 and 6. The additional crossed box contribution is small as well as those from Figs. 3c and 3d, so that in this respect ref. 22 should be rather good. However, the "class III" transition potential box cancels much of $\gamma_1$ (solid vs. dashed curves). This interaction was included in ref. 22. As
Fig. 5. The mixing parameters $\gamma_1$ and $\gamma_2$ with coupled channels class IV cross product transition potential (dashed) and the same by a two pion exchange potential (dotted). The solid curve is the sum of the dashed and residual effects of Figs. 4b-d and "class III" transition potential box.

Fig. 6. As Fig. 5 for $\Delta A$. The dash-dot curve is the coupled channel result with a CSB photon exchange transition potential.

A consistency check of the validity of the method is made first between direct boxes calculated by the coupled channels and as a two pion potential evaluated at zero energy (dashed vs. dotted curves). The good agreement between the two especially far below the $\Delta$ threshold supports the accuracy of the method. Since the crossed boxes should not show
Fig. 7. $\Delta A$(OPE) for the Reid potential (solid) and coupled channels modified Reid (dashed) with $\Lambda_\pi = 800$ MeV and the latter with $\Lambda_\pi = 1200$ MeV (dotted). The dash-dot curve is the negative of the gluon exchange result of ref. 15.

such strong energy dependence and threshold effects as the direct ones the same potential can be applied for them also at higher energies. Overall, the effect in $\Delta A$ at 477 MeV is about 15-20% of the dominant OPE contribution, of the order of the mass difference effect in $\rho$ exchange. Furthermore, the dash-dot curve shows a preliminary result of photon exchange in the transition potential.

It was pointed out earlier that (with a possible exception of the Bonn OBEPR) the names of the diagonal $NN$ potentials are of little consequence for CSB as long as they are phase equivalent. The inclusion of $N\Delta$ components by coupled channels may change essentially the short range part of the $NN$ wave function by absorbing some probability. However, Fig. 7 shows that this effect is not large on $\Delta A$. The solid curve is calculated using the pure Reid soft core potential, whereas the dashed is a coupled channels calculation with a modified Reid potential (to remove double counting), which reproduces the experimental phases well. For these a rather soft dipole form factor was used in the vertices with $\Lambda_\pi = 800$ MeV. For comparison, the dotted curve is the coupled channels result with $\Lambda_\pi = 1200$ MeV, more in line with the Bonn potential.

Finally, it is worth noting that the constraint of isospin breaking being possible only in the tensor uncoupled states with $J = L$ is valid solely for spin and isospin $\frac{1}{2}$ particles. Once isobar admixtures are allowed, also tensor coupled states can have CSB. In this case isospin can be changed either by transition potentials of isotensor type (class II) or by interactions resembling class III interactions as well as class IV. Consequently, for example the deuteron could have small isospin breaking $N\Delta$ components of the order of one per mille in the wave
Unfortunately it is difficult to imagine a situation where this could be seen in the two nucleon system, but it might well be an important effect in the isospin breaking reaction \( d + d \rightarrow ^4\text{He} + \pi^0 \), for which there have been extensive experimental searches.

**OPEN QUESTIONS**

In summary, theoretically CSB interactions offer a possibility to use and test familiar concepts from both nuclear and particle physics in a new context, in many ways complementary to traditional isospin symmetric interactions. This involves a mixture of well known meson dynamics and still unanswered questions. Firstly it has been shown that the recent TRIUMF\(^6\) and IUCF\(^8\) measurements of \( \Delta A = A_n - A_p \) at the angle where \( A \) crosses zero can be easily and reliably explained by the \( np \) mass difference effect in OPE. However, even in this different form factors can vary the CSB effect by about 30%. Outside this angle more uncertainties arise. Some relevant meson couplings are badly known, in particular the \( \omega NN \) coupling in \( \rho\omega \)-mixing. Also the "\( \sigma \) meson", largely simulating two pion exchange with crossed pions and/or \( N\Delta \) intermediate states, should apparently be resolved into these more microscopic ingredients with distinct isospin structures. This leads to the role of isobars and two pion exchanges in general. The calculation of their contribution to \( \Delta A \) in ref. 22 can be improved by including crossed boxes, \( \rho\omega \)-mixing and the magnetic interaction. Another quite long ranged mechanism never considered in class IV forces is crossed \( \gamma\pi \) exchange, which has been shown to be important for class III.\(^{25}\) So already at the hadron level much interesting physics awaits.

At the deeper, quark level one could question and test the validity of the meson exchange picture at short distances. Generally, quark-gluon and meson exchanges appear to give qualitatively similar results for the isospin respecting interaction. However, CSB \( \rho\omega \)-mixing and quark-gluon exchange may be qualitatively different.\(^{15}\) One might question the origin of CSB. How much is to be attributed to quark mass differences, how much to electromagnetic effects? In early theoretical work much attention was paid to EM effects at hadron level. For example \( \omega \) mesons could virtually annihilate into a photon and reemerge as a \( \rho \) meson, or an \( \omega \) could exchange a \( \delta \) meson with a sea photon and change into a \( \rho \) (tadpole contribution). These seemed to be enough to explain the contemporary data on meson mixing (for a review see e.g. ref. 26). Nowadays much of the mixing could be explained in terms of the \( u \) and \( d \) quark mass difference. How can these different mechanisms be combined so as not to overcount? The interrelations are not trivial. For example, it has been suggested that the \( np \) mass difference would arise from \( \rho\omega \) mixing without a need for up and down quark mass differences.\(^{27}\)

It is rather amusing that a 50% effect in the current quark masses reduces to less than 1% effect in baryon masses and to nearly negligible effects in interactions. This may be understandable, since constituent masses are much larger than current masses. However, here
apparently a small perturbative strong interaction component with distinct isospin properties
is generated in the otherwise nonperturbative realm of low energy quarks and gluons. It would
be of great interest and importance to see whether and how this could be used to obtain new
insight to QCD itself.

As seen above, there are interesting and tractable problems in theory. Experimentally
in \( np \) system there are only two data points for CSB.\textsuperscript{6,8} The effects are small and their separa-
ration from the background and overcoming systematic errors may be an experimentalist’s
nightmare. In principle accurate data on \( \Delta A \) even at the zero crossing angle at a low and
an intermediate energy can help resolve the \( \pi N N \) form factor on one hand and distinguish
the \( \rho \omega \) mixing contribution vs. quark-gluon model on the other. In the latter case angular
dependence would be definitely a better signature, but at present experimentally unavailable.
Formidable efforts are required on the experimental side in order that CSB would not remain
a theorist’s problem.

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FORWARD - BACKWARD ASYMMETRY IN THE REACTION np + dπ^0

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ABSTRACT

Two experiments are being designed to measure the forward-backward asymmetry in the center of mass differential cross sections for the np + dπ^0 reaction and thus provide additional tests of the charge symmetry breaking part of the NN interaction. One of these experiments is proposed at 477 MeV at TRIUMF and the other at selected energies between 600 and 800 MeV at LAMPF. The experimental challenges in measuring these cross section differences to the order of a few tenths of a percent are described. In addition, preliminary results at 800 MeV are presented from a Monte Carlo code written to simulate the LAMPF experiment.

INTRODUCTION

The first unambiguous violation of charge symmetry breaking (CSB) in the NN interaction was observed at TRIUMF in the recent measurement at 477 MeV of ΔA(0) = A_n(0) - A_p(0) at the zero crossing angle of the n-p analyzing power. The TRIUMF collaboration is planning another measurement of this quantity this fall with a 350 MeV neutron beam. A second experiment has recently been completed at 183 MeV at IUCF in which ΔA(0) or a quantity directly related to it has been measured for a range of angles around and including the zero crossing angle.

In order to test the predictions of CSB effects it is necessary to extend the range of energies over which experiments are carried out and also to measure different observables sensitive to the CSB interactions. Recent theoretical predictions have shown that the forward-backward asymmetry in the np + dπ^0 reaction, defined in the center of momentum frame (c.m.) as

\[ \frac{\sigma(\theta^+) - \sigma(180^° - \theta^+)}{\sigma(\theta^+)} \]

Fig. 1. Fractional forward-backward asymmetry plotted as a function of energy for three c.m. scattering angles (Ref. 4)
may be as large as 1%. The energy dependence as predicted at three c.m. angles is shown in Fig. 1. In these calculations the two dominant CSB mechanisms are found to be (1) the n-p mass difference in one pion exchange and (2) [$\eta - \pi^0$] mixing.

\[
A_{FB}(\theta) = \frac{\sigma(\theta^*) - \sigma(\pi - \theta^*)}{0.5[\sigma(\theta^*) + \sigma(\pi - \theta^*)]} \tag{1}
\]

THE EXPERIMENTS

The most extensive previous efforts to measure $A_{FB}$ in the $np + d^0$ reaction have been reported by Bartlett et al.\textsuperscript{5} and Hollas at al.\textsuperscript{6}. Neither of these experiments were carried out with sufficient precision to measure $A_{FB}(\theta)$ but rather reported asymmetries in terms of the integrated forward cross section minus the integrated backward cross section, with this difference divided by the total cross section. The Bartlett experiment used a continuum neutron beam and reported ten values of the asymmetry between 300 MeV and 650 MeV with an overall average $A_{fb} = (1.0 \pm 1.9)\%$. The Hollas experiment used the 795 MeV neutron beam at LAMPF and reported the value $A_{fb} = (-0.3 \pm 1.0)\%$.

The TRIUMF and LAMPF experiments which are in the planning stages seek to achieve a precision near $\pm 0.1\%$. The TRIUMF experiment will use a 477 MeV neutron beam produced at $9^0$ with 497 MeV protons incident on a LD2 target. The deuteron spectrometer will be set up on the neutron beam line in order to concentrate on the region near $0^0$ and $180^0$ in the c.m. The LAMPF experiment proposal is presently being configured with the Medium Resolution Spectrometer (MRS) now being installed in Area B at LAMPF. The MRS has a
momentum acceptance $\Delta p/p = \pm 20\%$ and a momentum resolution $\delta p/p = \pm 0.2\%$. The momentum acceptance of MRS precludes making simultaneous measurements of the forward and backward cross sections for c.m. angles less than about 30° at 50 MeV and about 50° at 800 MeV. The kinematics at 800 MeV are shown in Fig. 2 for the np + d$n^0$, np + d$\gamma$ and np + np reactions. The rectangular box in Fig. 2 indicates the maximum acceptance of the spectrometer.

As already implied, to minimize the systematic errors associated with the incident beam flux and target density variations these experiments are designed to measure the forward and backward yields simultaneously. The ratio of the c.m. cross sections can then be written as

$$\frac{\sigma(\phi^*)}{\sigma(\pi\phi^*)} = \frac{Y(\theta_F)}{Y(\theta_B)} \frac{J_{BE_Bd\theta_B}}{J_{Fe_Fd\theta_F}}$$

(2)

where the $Y$'s represent the observed deuteron yields obtained in the laboratory solid angles $d\Omega_B/F$ at the lab angles $\theta_B/F$ corresponding to the c.m. angles $(\pi-\phi^*)/\phi^*$. The $\varepsilon_{B/F}$ represent the respective efficiencies for detecting the deuterons emerging at the conjugate angles and include 1) detector efficiencies which may be momentum dependent and 2) corrections for losses due to deuteron breakup which are momentum dependent and are not well known. The $J$'s represent the Jacobians for the two lab to c.m. cross section transformations. The ratio of the Jacobians is both energy and angle dependent and, together with the deuteron breakup uncertainties, presents the most serious systematic errors in these experiments.

SYSTEMATIC ERRORS ASSOCIATED WITH $J(Tn,\theta)$

The ratio of the forward and backward angle Jacobians must be known to $\pm 0.1\%$ to achieve the desired accuracy in the c.m. cross sections. This requires precise knowledge of the neutron beam energy and the absolute calibration of the spectrometer angle. At 477 MeV the Jacobian ratio (F/B) changes by 0.2% per MeV beam energy change near 0°. Thus the central beam energy must be known to 0.25 MeV and its stability must be comparable. In positioning the spectrometer at 0° this experiment will be monitoring the neutron beam direction and consequently the absolute scattering angle is determined.

The variation of the Jacobian ratio with angle is a much more difficult problem in the proposed experiment at LAMPF. As the lab angle moves away from 0° the Jacobian changes with increasing rapidity. At 800 MeV the Jacobian at $\theta^* = 130°$ changes by 5% with a 0.1° change of the lab angle. The present effort is directed at simulating the experiment with the Monte Carlo calculation in order to study the behavior of the "effective Jacobians" and their ratio. Since the beam direction cannot be monitored concurrent with measurements in the MRS, three measurements are proposed to determine both the beam energy and the absolute spectrometer angle.

1) Determining the crossing angle where the momentum of the np elastic protons equals that of the np + d$n^0$ deuterons.

2) Determining the maximum lab angle at which the d$n^0$ deuterons are
observed. The sensitivity of this measurement is 0.012°/MeV.

3) With a π° spectrometer at 90° lab triggering on events having the two decay gamma rays lying in a plane normal to the d-π° reaction plane the pion polar scattering angle is determined. Achieving 1° accuracy for the pion angle results in 0.13° accuracy for the polar scattering angle of the coincident deuterons. The resulting sensitivity is dΩ/dTn = 0.005°/MeV.

MONTE CARLO

A Monte Carlo program has been written to study the combined effects of the energy spread of the neutron beam, multiple scattering, finite geometry, losses due to deuteron breakup, and detector resolution. Preliminary results have been obtained for the case of Tn = 800 MeV and θ* = 58°. This is an angle where the predicted fractional asymmetry is large but so are the difficulties in reducing the systematic errors associated with the Jacobian ratio.

The c.m. angular distribution used as input for the calculations is the parameterization of Richard-Serre at 794 MeV for the even powers of cosθ* plus a term in cosθ* which has the approximate shape of the predicted asymmetry for c.m. angles larger than 45° and has an amplitude such that Aπ(45°) = 1.0%.

The two multi-wire drift chambers (MWDC) which determine the deuteron scattering angle are placed with an 80 cm separation in the region between the LH2 target and the MRS entrance. A Gaussian resolution function with .3 mm (FWHM) is used to randomize the hits in both x and y coordinates in each chamber. The angular resolution of the system is determined by computing dθ for each event where dθ is the difference between the true lab scattering angle, determined from the picked c.m. angle, and the measured scattering angle determined by the MWDC coordinates. The full width at the base of the distribution of dθ for all deuterons entering the spectrometer is 0.12° when detector resolution is the only effect. Finite geometry effects increase this number to 0.28° and the additional effects of multiple scattering further increase it to 0.31°. The finite geometry assumes an 0.6 mm diameter proton beam incident on a 10 cm LD2 neutron production target with a 5 cm diameter LH2 target located 9.0 m downstream at the MRS pivot. At this time a monoenergetic neutron beam is assumed in the code. Taking lab angle bins of 0.3° width, the momentum resolution (full width at base) for the bin centered on θ* = 58° develops as follows:

   a) The 0.3° acceptance            Δp = 11 MeV/c,
   b) Full Monte Carlo                 Δp = 15 MeV/c,
   c) Incl. 12 MeV/c beam spread          Δp = 19 MeV/c.

The central momentum in this bin is 1461 MeV/c and the deuterons from np + dy are centered at about 1500 MeV/c in this bin. The design resolution of the MRS is 2.1 MeV/c at this momentum so there should be negligible background from the dy deuterons.

Corresponding numbers for the conjugate backward angle at θ* = 122° are:
a) The 0.3° acceptance $\Delta p = 23$ MeV/c,
b) Full Monte Carlo $\Delta p = 35$ MeV/c,
c) Incl. 12 MeV/c beam spread $\Delta p = 36$ MeV/c.

In this case the central momentum is 1072 MeV/c and the dy deuterons are cleanly resolved at 1000 MeV/c.

Initial calculations of the effective Jacobian have been completed. With roughly one million events in each of the two lab bins which correspond to 2° bins in the c.m. we find

\[ J_{\text{eff}}(58^\circ) = 30.91 \pm 0.4, \quad \text{and} \]
\[ J_{\text{eff}}(122^\circ) = 44.54 \pm 0.06, \]

where the errors are statistical. A summary of these and additional calculations is shown in Fig. 3 together with curves showing the kinematic Jacobians as a function of c.m. angle at 800 MeV.

Fig. 3. Results of the Monte Carlo calculation of $J_{\text{eff}}(\theta^*)$ using 2° bins in the c.m. (horizontal bars). Solid curves are the kinematic values of $J(\theta^*)$ and $J(\pi-\theta^*)$ at $T_n = 800$ MeV.
CONCLUSIONS

The proposed experiments represent an effort to extend the experimental evidence for CSB interactions to additional observables and over an extended energy range. These data should be of considerable help to the theorists in refining their calculations of CSB effects. The TRIUMF measurement at c.m. angles near 0° and 180° would seem to have every chance of achieving the desired accuracy in determining the front/back asymmetry in the differential cross section. The LAMPF proposal is an attempt to extend the angular range of these measurements and must deal with much larger potential systematic errors in dealing with the Jacobians. It is hoped that some running of the pp + dπ+ reaction will be possible during the period this summer and fall when the MRS is being brought into operation. The Monte Carlo calculations are being continued and input to the code could be refined by comparing it to the dπ+ data. The use of MRS for this experiment may not be feasible and it may be necessary to use a lower dispersion magnet and restrict the 600 to 800 MeV measurements to angles near 0°.

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CHARGE SYMMETRY BREAKING IN FEW-NUCLEON SYSTEMS

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ABSTRACT

The investigation of charge symmetry breaking effects in few-nucleon systems is reviewed. The conditions for charge symmetry in mirror reactions with light nuclei are discussed. Examples of results of measurements of the differential cross section and polarization observables of mirror reactions and their analysis are presented.

INTRODUCTION

One of the basic assumptions in nuclear physics is the charge independence of the nucleon-nucleon (NN) interaction. This symmetry postulates that the protons and neutrons are identical except for their isospins. The isospin symmetry, i.e. invariance under rotations in the isospin space, has the special case of the rotations of 180°, which interchange the proton and neutron states and is called charge symmetry. However, nucleons are not completely identical and therefore these symmetries must be broken by small amounts due to the differences in the meson and nucleon masses and the electromagnetic effects. At the quark level, the charge symmetry operation interchanges d and u quarks. Since the mass difference of these two quarks \( \Delta m = m_d - m_u \) is about 3 to 4 MeV and since d and u quarks have electrical charges \(-1/3 \) e and \(+2/3 \) e, respectively, the charge symmetry is intrinsically broken. In the comparison of the pp and nn systems the charge symmetry breaking appears trivially through the presence of the Coulomb interaction. From the experimental point of view, the investigation of these two systems in order to extract the strong interaction part of the charge symmetry breaking is difficult. For this goal few-nucleon systems with \( A > 3 \) are more suitable for tests of the charge symmetry. In particular, this symmetry violation can be explored in mirror reactions. When exact charge symmetry holds the differential cross sections and the polarization observables of mirror reactions are expected to be identical. Since, however, the Coulomb interaction as well as the nucleon mass difference breaks this exact symmetry these simple conclusions from such experiments may not be straightforward. In very light nuclei, Coulomb effects are expected to be small and therefore mirror reactions with few-nucleon systems are favorable processes for such an investigation.

It is suggested that detailed model calculations based on the nucleon mass difference and or the isoscalar-isovector-meson mixing, which is dominated by \( \rho - \omega \) and \( \pi - \eta \) mixing as well as the electromagnetic interactions lead to a understanding of the effects caused by the interchange of the u and d quarks.

A classification of charge dependent forces has been given by Henley and Miller according to the isospin dependence which is based on nuclear and meson degrees of freedom. The sensitivity of different tests of charge symmetry has been reviewed by Nefkens and Slaus by grouping various systems and reactions by their dependence on the d-u quark mass differences.
MIRROR REACTIONS

Mirror reactions in very light nuclei have, besides the small Coulomb effects, the advantages that the systems are treatable in the context of realistic model calculations and systems can be found, which are free from resonances over larger excitation regions or have well separated resonances with different widths. Further, the tiny charge symmetry breaking effects can be enhanced in few-nucleon systems due to inference effects.

Favorite reactions for tests of charge symmetry are the mirror reactions given in Table I and the corresponding pairs of inverse reactions.

<table>
<thead>
<tr>
<th>Reaction</th>
<th>Observables</th>
<th>Ref.</th>
</tr>
</thead>
<tbody>
<tr>
<td>n-d p-d scattering</td>
<td>σ, Ay</td>
<td>5</td>
</tr>
<tr>
<td>$^3\text{He}(n,\gamma)^4\text{He}$</td>
<td>σ, Ay</td>
<td>6,7</td>
</tr>
<tr>
<td>$^3\text{H}(p,\gamma)^4\text{He}$</td>
<td>σ, Ay</td>
<td>8-11</td>
</tr>
<tr>
<td>$^2\text{H}(d,n)^3\text{He}$</td>
<td>σ, iT, T_{21}, T_{22}</td>
<td>12,13,17-20,22-24</td>
</tr>
<tr>
<td>$^2\text{H}(d,p)^3\text{H}$</td>
<td>σ, iT, T_{21}, T_{22}</td>
<td>12,13,17-20,22-24</td>
</tr>
</tbody>
</table>

THE MIRROR REACTIONS $^2\text{H}(d,p)^3\text{H}$ AND $^2\text{H}(d,n)^3\text{He}$

In these reactions corresponding observables are expected to be identical. The Q-values of these reactions, however, are different, +3.269 MeV for $^2\text{H}(d,n)^3\text{He}$ and +4.033 MeV for $^2\text{H}(d,p)^3\text{H}$ due to the Coulomb displacement energy. Thus, deuteron laboratory bombarding energies about 1.5 MeV higher in the reaction $^2\text{H}(d,n)^3\text{He}$ lead to the same center-of-mass neutron and proton energies in the n+$^3\text{He}$ and the p+$^3\text{H}$ exit channels.

In comparison of observables it has been proposed that the laboratory energy scale of the incident deuterons of the $^2\text{H}(d,p)^3\text{H}$ reaction should be shifted down by 1.5 MeV in order to bring the nucleons in the exit channels to the same energy. This method is considered as a possible way to take the Coulomb effects into account, but of course such an ad hoc shift can only be considered as a rough approximation. To take proper account of the Coulomb interaction requires a detailed understanding of the reaction mechanism.

Another possibility to judge the influence of Coulomb effects is the study of these mirror reactions over a wide energy range of the incident deuterons. The importance of Coulomb effects is expected to decrease with increasing energy.

We have carried out an extended experimental program in order to investigate the discrepancies of the analyzing power. High precision data of the polarization observables iT_{11}, T_{20}, T_{21} and T_{22} of the two mirror reactions have been measured between 1.0 and 24 MeV^{9,11}. Energy steps of 1.5 MeV have been chosen such that the results can be compared at the same energies in the entrance or exit channel and the behaviour of the deviation can be studied over a wide energy range. Data of the $^2\text{H}(d,n)^3\text{He}$ reaction at lower energy (1.5 to 4 MeV) from the Ohio state group^{10} have been used to compare to our low energy $^2\text{H}(d,p)^3\text{H}$ data.

Samples of the experimental results of both reactions are compared in Fig. 1. The solid curves represent Legendre polynomial fits to the $^2\text{H}(d,p)^3\text{H}$ reaction, the dashed curves are the results of the analysis of the (d,n) reaction. In one particular plot only the experimental results of one reaction are shown, since in many cases the corresponding mirror reaction data are very close or overlapping each other, so that an unclear representation would result.

In Fig. 1 comparisons are made at the same incident deuteron energies. The vector analyzing power iT_{11} shows only small differences over practically the whole energy range.
In contrast to this behaviour large discrepancies are observed in Fig. 1 for the even tensor component T_{20}.

Comparisons of the same quantities with the ^2\text{H}(d,p)^3\text{H} results shifted down in energy by 1.5 MeV are performed too. The even component T_{20} is then in much better agreement, particularly at the higher energies. However, strong deviations are observed for the odd component iT_{11} over the whole energy range.

Fig. 1. Comparison of the analyzing powers iT_{11} and T_{20} of the reactions ^2\text{H}(d,p)^3\text{H} and ^2\text{H}(d,n)^3\text{He} at the same deuteron energy. The dots are the experimental results of the ^2\text{H}(d,p)^3\text{H} reaction, the solid curves are fits to these data. The dashed curves are fits to the data of the mirror reaction ^2\text{H}(d,n)^3\text{He}.

For a quantitative review of the situation, it is advantageous to define the following average as a measure of the deviation

\[ D_{av} = \int_{\theta_1}^{\theta_2} D_{kq}(\theta)d\theta/ (\theta_2 - \theta_1) \]  

(1)

with

\[ D_{kq} = |T_{kq}(d,p) - T_{kq}(d,n)| \]  

(2)

The angular range of the comparison is chosen to be the same for all energies namely \( \theta_1 = 100^\circ \) and \( \theta_2 = 160^\circ \). This is a compromise for the different angular ranges of the data at different energy values.
Fig. 2. Energy dependence of the average deviations $D_{av}$ of the analyzing powers $iT_{11}$, $T_{20}$, $T_{21}$ and $T_{22}$ for the mirror reactions compared at the same entrance energies.

energies. This angular region is well determined for all energies by the present data. The results of this calculation between 1.5 and 24 MeV for the same entrance channel energy are presented in Fig. 2. The crosses are the results comparing the $(d,n)$ data of ref. 10 with our low energy $(d,p)$ data. The dot size indicates the effect of the experimental uncertainty.

Fig. 3. Energy dependence of the average deviations $D_{av}$ of the analyzing powers $iT_{11}$, $T_{20}$, $T_{21}$ and $T_{22}$ for the mirror reactions compared at the same exit channel energies.
on the calculation; the error bars shown with the crosses reflect the larger uncertainty at lower energies due to the larger statistical errors and lower number of angles measured of the $^2$H(d,n)$^3$He reaction. At $E_d=4$ MeV the results are in fair agreement from both (d,n) data sets. The solid lines are drawn to guide the eye.

It is interesting to calculate also the deviation $D_{av}$ for the energies shifted by 1.5 MeV, comparing in this way the observables at the same exit channel energy. These results are shown in Fig. 3. The symbols have the same meaning as in Fig. 2. In general stronger deviations for odd analyzing power components and a decrease of the disagreement in the even components is clearly exposed by this representation. It is also clear from Figs. 2 and 3 that for some components the behaviour below 4 MeV is different from that at higher energies.

As a result of this comparison, one finds clear evidence for deviations in all four analyzing powers. This fact is the more astonishing as the disagreement extends smoothly over an energy range from 1.5 to 24 MeV, a factor in energy of more than ten.

The deviations observed in the comparison of different kinds of observables are large and complex. This observation applies in particular to the low energy results, where, however, they might be understood as an interference between nuclear and Coulomb scattering. At higher energy evidence for a weak energy dependence is observed when compared at the same entrance channel energy. However, at the same exit channel energy the quantity $D_{av}$ is constant in the limits of the uncertainties of the results between 4 and 24 MeV. This feature can be considered as a success of the approximate Coulomb correction in the exit channels. Although the polarization effects in the d-d elastic channel are small, their variation over the energy range in question is large and therefore the constancy of the differences as a function of energy can hardly be compatible with such a strong energy dependent behaviour in the entrance channel. For these reasons charge symmetry violation has to be considered seriously as a possible explanation for the behaviour of the polarization observables of these mirror reactions. It is obvious that these arguments are not a proof in a strict sense, but they support the hypothesis of non-conservation of charge symmetry. Clearly no final conclusion can be drawn until detailed reaction calculations which include the Coulomb force are available. These calculations can decide whether the present strong indication of the non-conservation of charge symmetry of nuclear forces or a combination of strong nuclear forces with charge symmetry property and electromagnetic forces are the cause of the experimental results.

**THE REACTION $^4$He(d,$^3$He)$^3$H**

This reaction is characterized by the two trinucleon particles in the exit channel. The consequences for the differential cross section due to particle symmetry in the exit channel of the reaction type a(b,c)c' have been discussed by Barshay and Temmer. In this notation the particles c and c' are the charge-symmetry members of an isospin doublet. Simonius has extended the considerations to spin observables.

Charge symmetry conservation demands that the angular distributions of the cross section $\sigma=T_{oo}$ and the $T_{2q}$ with even q are symmetric in respect to $\theta_{cm}=90^\circ$ whereas the component $iT_{11}$ and $T_{21}$ are antisymmetric in the c.m. system i.e.

$$\sigma(\theta) = \sigma(\pi - \theta)$$

and

$$T_{kq}(\theta) = (-1)^q T_{kq}(\pi - \theta)$$
This relation implies that there is no interference between positive and negative parity channels in the symmetry limit.

We have measured the cross section and the polarization observables of the $^4\text{He}(d,^3\text{He})^3\text{H}$ reaction with the aid of the polarized deuteron beam from the PSI injector cyclotron at six energies in the energy range between 30 and 50 MeV\textsuperscript{12}. Six telescopes, each composed of three detectors, are used at the same time. They are installed in the scattering chamber symmetrically with respect to the incident beam. The first two detectors of each telescope are used as a $\Delta E$-E telescope to discriminate $^3\text{He}$ against $^4\text{He}$, whereas the second and third detectors act as a $\Delta E$-E telescope to discriminate $^3\text{H}$ against protons and deuterons. In this way the $^3\text{He}$ and $^3\text{H}$ are measured at the same time at three different angles with the same geometrical conditions. The measured trinucleon particles correspond to the forward or backward angular distribution in the c.m. system with an overlap near 90°.

Samples of the results of the cross section and the vector- and tensor analyzing powers in cartesian representation at 35.15 MeV are shown in Fig. 4. While violation of the symmetry rule in the cross section and the vector analyzing power $iT_{11}$ has been observed already earlier (see e.g. refs. 17 and 18) this effect has been observed for the first time also in the tensor components. As can be seen from Fig. 4 the deviations are enhanced in these components and therefore the investigation is facilitated.

An analysis of the present large amount of data in terms of Legendre polynomials is interesting since the resulting small number of coefficients can be considered as a bridge between the theoretical and experimental investigations. Further the result of such an analysis constitutes an overall consistency test of the data. Finally the coefficients with odd L values reflect the charge symmetry breaking effects. The curves in Fig. 4 are the results of the Legendre polynomial analysis. In the analysis values of coefficients of L(odd) significantly different from zero are found for all observables but in general the odd L-values are an order of magnitude smaller than the even L-value coefficients. An energy dependent structure is observed for these coefficients over the measured energy range.

For a quantitative comparison of the deviations of the different observables the following symmetry functions are calculated

\begin{align*}
W(\theta) &= [\sigma(\theta) - \sigma(\pi - \theta)]/[\sigma(\theta) + \sigma(\pi - \theta)] \quad (5) \\
D_{11}(\theta) &= iT_{11}(\theta) + iT_{11}(\pi - \theta) \quad (6) \\
D_{2q}(\theta) &= T_{2q}(\theta) - (-1)^q T_{2q}(\pi - \theta) \quad (7)
\end{align*}

The determination of these functions are free from asymmetry errors in the angular setting in the laboratory system since the simultaneous measurement of the $^3\text{He}$ and $^3\text{H}$ particles at a particular angle $\theta_{\text{Lab}}$ corresponds already to the symmetry angles $\theta$ and $(\pi - \theta)$ in the c.m. system. The change of the angle $\theta_{\text{cm}}$ due to the mass difference between $^3\text{He}$ and $^3\text{H}$ is negligibly small, namely only 0.01°.

Samples of the results are shown in Fig. 5. In most cases the dot size is larger than the statistical error of the results. The curves are calculated from the Legendre polynomial fits. For the cross section deviation function $W(\theta)$ also the results of ref. 19 (crosses), ref. 17 (open circles), and ref. 20 (83 MeV, open circles) are shown.
Fig. 4. Differential cross section, the vector analyzing power $A_y$ and the tensor analyzing powers $A_{yy}, A_{xz}$ and $A_{xx}$ of the reaction $^4\text{He}(^3\text{He})^3\text{H}$ at $E_d=35.15$ MeV.
Fig. 5. The symmetry functions $W(\theta)$ of the cross section between 32 and 83 MeV and the function $D_{20}(\theta)$ from the tensor analyzing power $T_{20}=-(A_{xx}+A_{yy})/\sqrt{2}$ between 32.1 and 49.7 MeV.
It is already obvious from Fig. 4 that the symmetry conditions (3) and (4) are not fulfilled. In particular, in the tensor components, large deviations occur. Deviations from the expected symmetry behaviour are observed for all functions (5) to (7) and at all measured energies. These deviations show strong variations as a function of energy and reach values of up to 0.2. While the values of the functions $W(\theta)$ and $D_{11}(\theta)$ are decreasing with increasing energy, for the $D_{2q}$ this tendency is less pronounced. The complex structure and their energy dependence suggest that the observed charge symmetry breaking effects are not due to simple Coulomb interactions but a complicated interplay between the Coulomb and nuclear forces.

The functions $W(\theta)$ and $D_{20}(\theta)$ show a very similar behaviour in the angular region of 20° to 30°. With increasing energy an interference-like pattern develops, shifting the maximum values to smaller angles. In a similar way a structure around $\theta_{cm} \approx 60°$ appears also in all five components of the symmetry functions (eqs. 5 to 7). These maxima in the functions 5 to 7 are correlated with the minima observed in the differential cross section. It is well known that polarization observables stem from interference effects between different partial waves, but the presently observed structures of the symmetry functions only partly reflect the structures of the polarization observables themselves. A priori one might expect a large deviation from the symmetry in the region between $E_d = 32$ to 45 MeV corresponding to excitation energy of $^6\text{Li}$ from 23 to 31 MeV where broad overlapping $T=0$ and $T=1$ resonances are known to exist. In particular three broad $T=0$ levels have been observed in an analysis of $d-\alpha$ elastic scattering.

The threshold deuteron energy of the present reaction (21.5 MeV) corresponds to the negative Q-value - 14.3 MeV and an excitation of $^6\text{Li}$ of 15.8 MeV. This excitation region has been investigated by the inverse reaction $^3\text{He}(^3\text{H,}\alpha)^2\text{H}$ with $E_t$ in the 1 MeV region. The intermediate $^6\text{Li}$ excitation between 20.3 to 24.3 MeV has been investigated by the $t-^3\text{He}$ elastic scattering by Haglund et al., where two broad $T=1$ ($0^-,2^-$) levels are known.

Departures from the symmetry relations reflect the effect of either isospin impurities in the particles of the entrance channel or the impurity introduced by the reaction mechanism and the intermediate state or the breakdown of isobaric analogy between the particles in the exit channel. Naturally also a combination of these reasons is possible. For the present reaction it seems most unlikely that isospin mixing in the entrance channel due to the spin singlet production $\alpha+d\rightarrow\alpha+d^*$ occurs since the contribution of this reaction at our energies is rather small. Furthermore, we remark that the $T=1$ states for the $A=6$ system, which could in principle mix with $T=0$ states do not couple to the $\alpha+d^*$ channel because they are unnatural parity states. Although at 32 and 41 MeV DWBA calculations can produce $W(\theta)$ for the cross sections and $D_{11}(\theta)$ for the vector analyzing power of about the correct magnitude due to the slightly different neutron and proton transfer amplitudes which are sensitive to a difference in $^3\text{He}$ and triton wave function, there is a disagreement in the angular distribution since not even the sign is correctly reproduced in the angular region of the largest values of $W(\theta)$ and $D_{11}(\theta)$. This suggests that the reaction mechanism is not dominated by the one nucleon exchange, but involves the coupling to the various channels.

Indeed it is clear that the coupling to the $5+1$ channels ($5\text{He}+p,5\text{Li}+n$) can be extremely relevant because of the charge asy mmetry effects in the level scheme. In particular the 22.32 and 25.66 MeV $^5\text{Li}$ levels are analogous to the 21.35 and 24.39 MeV $^5\text{He}$ levels, respectively, whereas the 23.66 MeV $^5\text{Li}$ level and the broad 28.5 to 29.5 MeV $^5\text{He}$ level do not have partners. For a detailed analysis a more elaborated calculation based on the resonating group model or an R-matrix analysis should be performed. Preliminary results of
a resonating group calculation have been reported by Bruno et al.\textsuperscript{26} and further calculations are still in progress. Considering only positive parity channels the calculation generates the correct shape of the angular distributions of cross section and vector analyzing power at $E_d=35$ and 45 MeV.

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CHARGE SYMMETRY OF THE NUCLEAR INTERACTION
AND THE N-N SCATTERING PARAMETERS

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ABSTRACT

The precise determination of the N-N scattering parameters is one of the most sensitive tests of the charge symmetry breaking (CSB) of the nuclear forces. Analysis of low-energy scattering experiments and bound nuclear mirror systems give evidence of a departure from charge symmetry. Yet the subject has been controversial, due to the model dependence of the Coulomb corrections, and the significant difference of results with reactions involving two or three hadrons in the final state. Charge asymmetric potentials which include an updated value for the $p-\omega$ mixing, as well as other boson exchanges, are in good agreement with the experimental findings.

INTRODUCTION

The conjecture of the charge symmetry of the nuclear forces was advanced by Heisenberg shortly after the discovery of the neutron. In the absence of electromagnetic effects proton and neutron, the two charge states of the nucleon, should have identical properties. This simple idea was a far-reaching contribution to the understanding of the structure of matter, being the first internal symmetry introduced which acts on the particle identity, independent of space-time. A renewed interest in the study of the charge symmetry breaking originates in the development of a more fundamental description of the strong forces in terms of structureless constituents, which are the carriers of the symmetries within the hadrons, and in a new generation of high precision experiments at intermediate energies, which shed new light into this intricate field. A precise determination of CSB is important to find out whether the departures from charge symmetry can be explained on the basis of electromagnetism alone, or if there are isospin-breaking components in the purely hadronic sector. We know today, that it is not likely that CSB effects (such as the neutron-proton mass difference) would be generated radiatively within electromagnetism or a broader theory of the particle interactions. We are thus led to reexamine the origin of CSB at the hadronic substructure level and to see if the experimental results are consistent with the emerging picture.

Let us summarize the contents of this talk. After discussing the origin of isospin breaking in QCD, we describe the possible evidence of charge asymmetry from low-energy N-N scattering and energy differences in mirror nuclei. The evidence of CSB from new experiments at intermediate energies is not touched at all, since this is treated in detail by other speakers at this conference. We give some attention to the study of the reaction $\pi^{-}d \rightarrow \gamma nn$, which is an ideal tool for the extraction of the low-energy properties of the neutron- neutron interaction. Finally, we review briefly recent theoretical calculations of the charge asymmetry of the nuclear forces to see how well they confront the experimental findings.

Invited talk presented at the Symposium Workshop on Spin and Symmetries, TRIUMF, Vancouver, Canada. June 30-July 2, 1989
CSB IN THE CONTEXT OF QCD

The fundamental interactions among the quark constituents are described in quantum chromodynamics (QCD) by the Lagrangian

\[ \mathcal{L} = - \frac{1}{4} F_{\mu\nu}^a F^{\mu\nu a} + i \bar{\psi} \gamma^\mu D_\mu \psi - \bar{\psi} M \psi, \]  

where \( F_{\mu\nu}^a \) represents the field-strength and \( D_\mu \) the standard covariant derivative. Since exchanging a proton by a neutron corresponds to the replacement of a u-quark by a d-quark, we write the quark field in the form of an isospin doublet in analogy with the n-p doublet

\[ \psi = \begin{pmatrix} u \\ d \end{pmatrix}, \]  

and the generators of rotations in isospin space \( \mathbf{T} = \frac{1}{2} \mathbf{r} \), with

\[ \tau_1 = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \quad \tau_2 = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}, \quad \tau_3 = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}. \]  

To better understand the symmetries of the QCD Lagrangian, we expand (2) into its left (L) and right (R) chirality projections

\[ L = \frac{1}{2} (1 - \gamma_5) \psi, \quad R = \frac{1}{2} (1 + \gamma_5) \psi, \]  

and write \( M \) as a 2 x 2 diagonal matrix with the u- and d- quark masses, \( m_u \) and \( m_d \), as its diagonal elements. The QCD Lagrangian is split into two distinct terms: a scale-invariant term \( \mathcal{L}_\chi \) which only has a field content and is invariant under independent left or right \( SU(2) \) rotations

\[ \mathcal{L}_\chi = - \frac{1}{4} F_{\mu\nu}^a F^{\mu\nu a} + i L \gamma^\mu D_\mu L + i R \gamma^\mu D_\mu R, \]  

and a mass term \( \mathcal{L}_M \) which breaks the \( SU(2)_L \times SU(2)_R \) invariance of \( \mathcal{L}_\chi \)

\[ \mathcal{L}_M = - \frac{1}{2} (m_u + m_d) (\bar{L}R + \bar{R}L) - (m_u - m_d) (\bar{L}T_3 R + \bar{R}T_3 L). \]  

If \( m_u = m_d \), the QCD Lagrangian is invariant under \( SU(2)_{L+R} \sim SU(2)_V \), the isospin group, which corresponds to identical rotations of the L and R fields. The charge symmetry operation corresponds to a \( \pi \) rotation about the y-axis in isospin space. Charge independence is related to arbitrary rotations. Electromagnetism breaks charge symmetry since the charge operator \( Q = T_3 + Y \) (the hypercharge \( Y \) has the value 1/6) is not invariant under isospin rotations

\[ [e^{i\pi T_2}, Q] \neq 0. \]
The quark mass difference has the same dependence on $T_3$ as $Q$, $\delta m \sim (m_u - m_d) T_3$, thus

$$[e^{i\pi T_2}, \delta m] \neq 0,$$

and this is the origin of non-electromagnetic or explicit charge symmetry breaking effects in the N-N interaction. The u-d quark mass difference is actually responsible for the mixing of mesons with different isospin quantum numbers, and the result of such mixing is translated into medium-large distance effects between the nucleons. From a practical point of view, the particle mixing is parametrized in terms of a tadpole which depends linearly on the quark mass difference. Accordingly, any charge asymmetric effect should vanish in the isospin limit $e \rightarrow 0$, $m_u = m_d$. Additional quark mass effects could also be expected from the short range properties of the N-N wave function.

We know that isospin is badly broken at the quark level ($m_d/m_u \approx 1.8$). Why is isospin such a good symmetry? It is generally accepted that the quark and lepton masses are generated by the Higgs mechanism from the different Yukawa couplings in the electroweak sector of the theory. In the standard model, there are no zeroth-order relations among quark masses (isospin is not a natural symmetry of the model), and the masses are introduced as free parameters. Since QCD is a theory with a scale $\Lambda$, the observed conservation of isospin to within a few percent reveals the smallness of the u- and d-quark masses with respect to the hadronic scale $\Lambda \sim 1$ GeV (accidental symmetry), rather than the degeneracy of $m_u$ and $m_d$ (exact SU(2) symmetry). The contribution from non electromagnetic terms are of the order of $(m_u - m_d)/\Lambda$, and this effect is of the order of $0.5\%$. Let us recall that the weak interactions also violate isospin since only the left-handed doublet in Eq. 2 is gauged.

**CSB FROM MIRROR NUCLEI AND LOW-ENERGY N-N SCATTERING**

Many of the new experiments at intermediate energies\(^4\) aim at a direct detection of (CSB) in the neutron-proton (n-p) system by measuring polarization observables, as in n-p elastic scattering, or some other form of asymmetry, as in the inelastic reaction $np \rightarrow d\pi^0$. Experimentally, this method has an advantage over a comparison between the neutron-neutron (n-n) and proton-proton (p-p) systems, in that only a single experiment is required. Most of what is known about charge symmetry, however, comes from low-energy scattering experiments between nucleons.\(^2\) and the measurement of binding energies in mirror nuclei.

Early evidence of charge asymmetry in the nuclear forces came from the precise determination of the binding energies in the $^3H-^3He$ mirror nuclei.\(^5\) Of the 764 keV binding energy difference, only $683 \pm 29$ keV are attributed to direct electromagnetic contributions, leaving $81 \pm 29$ keV totally unexplained.\(^6\) The disagreement suggested a slightly stronger short-range n-n attraction to account for the differences\(^7\), which are observed throughout the periodic table. This discrepancy if known as the Nolen-Schiffer anomaly.

Low-energy scattering experiments between nucleons are among the simplest and most sensitive tests of the isospin symmetry of the nuclear interaction. Since low-energy scattering is dominated by the $^1S_0$ partial wave, the large scattering length corresponding to this almost bound state is particularly sensitive to charge asymmetric effects, which are inherently very small and consequently difficult to observe otherwise. Using the first order perturbation formula to evaluate the difference in the scattering lengths, $\delta a = a_{nn} - a_{pp}$, from a charge asymmetric component $\delta V$ of the nuclear potential

$$\frac{\delta a}{a} = am \int dr \ u_0^2(r) \ \delta V(r),$$

(9)
where \( u_0(\tau) \) is the asymptotic N-N wave function, it follows that

\[
\frac{\delta a}{a} \approx \frac{a}{b} \frac{\delta V}{V}, \tag{10}
\]

with \( b \) the radius of the potential. Since the coefficient \( a/b \) is a factor of order 10, any charge asymmetric effect is largely magnified. Typically, a 1\% change in the potential produces a change of 3 fm in the scattering length. The contribution of non-electromagnetic terms from quark mass differences discussed in the previous section is at the 0.5\% level. We could thus expect a 1.5 fm difference for \( \delta a \).

The actual determination of CSB from low-energy N-N scattering experiments has been a controversial issue. In the absence of direct neutron-neutron collision results,\(^8\) our present knowledge of the n-n interaction comes mainly from final-state interactions (FSI’s) where two or more final nucleons (two of which are neutrons) are present in the final state. The pronounced enhancement of the spectrum of a third particle near the end point of its kinematics is a measure of the very strong scattering of the two neutrons at low energies.\(^9\) The extraction of the n-n scattering parameters from the data depends on the theoretical model utilized. In the presence of three strongly interacting particles in the final state the calculation becomes quite involved and theoretical uncertainties difficult to evaluate. It is basically for this reason that different values of \( a_{nn} \) have often been found. To avoid large errors, the analysis of the data is usually restricted to a small kinematical region where n-n quasifree scattering dominates. The quoted “world average” value\(^10\) for the n-n scattering length \( a_{nn} = 16.6 \pm 0.6 \text{ fm} \) is smaller in absolute value than the Coulomb corrected value of \( a_{pp} = 17.1 \pm 0.4 \text{ fm} \),\(^11\) suggesting a stronger p-p force. The above value for \( a_{nn} \) is a straight average over some 40 measurements, most of them with three strongly interacting particles in the final state and with large uncertainties. There is no reason to give much relevance to this value.

Fig. 1. Experimental values of n-n scattering lengths.
There is a significant difference between the kinematically complete experiments from the neutron-induced deuteron breakup reaction \( nd \to pnn \) and other reactions with three hadrons in the final state such as \( ^3Hd \to ^3He nn \) or \( ^3Ht \to ^4He nn \). The results from \( nd \to pnn \) differ also considerably according to the distinct kinematical region studied: neutron pick-up \( (a_{nn} = -16.7 \pm 0.5 \text{ fm})^{12} \) or proton knock-on \( (a_{nn} = -20.7 \pm 2.0 \text{ fm}) \) processes. The values of \( a_{nn} \) for different experiments are shown in Fig. 1, which is taken from Ref. 8. The results are listed in chronological order from left to right. The last point corresponds to a recent high precision determination of \( a_{nn} \) from the reaction \( \pi^- d \to \gamma nn \) which gives the value \( a_{nn} = -18.5 \pm 0.4 \text{ fm}^{13} \). The difference

\[
|a_{nn}| - |a_{pp}| = 1.4 \pm 0.6 \text{ fm}
\]  

is consistent with the charge asymmetry in the \( ^3H - ^3He \) system and our simple order of magnitude estimate of the non-electromagnetic effects, but almost three standard deviations above (in absolute value) the “world average” value.\(^{14}\) A mechanism based on a specific model for the three-body forces which operates differently for neutron pickup and proton knock-on processes has been suggested\(^{15}\) to account for the difference between the results from \( \pi^- d \to \gamma nn \) and \( nd \to pnn \). Since no conclusive evidence of three-body forces has been found yet, it is difficult to ascertain the validity of the mechanism proposed to solve the discrepancy.

**EXTRACTION OF \( a_{nn} \) FROM \( \pi^- d \to \gamma nn \)**

As pointed out by Watson and Stuart almost four decades ago\(^{16}\), the capture reaction \( \pi^- d \to \gamma nn \) is ideal for the extraction of the n-n scattering parameters since the three particles in the final state are detectable and only the two neutrons interact strongly. Consequently, the determination of \( a_{nn} \) is free from the theoretical uncertainties inherent to other nuclear reactions. The extracted value of \( a_{nn} \) depends mostly on the asymptotic properties of the n-n wave function and is largely independent of the short-range description of the nuclear force. The scattering length is determined at low energies where the normalized spectrum is insensitive to secondary effects. For this reason, the extraction of \( a_{nn} \) with the different models proposed so far\(^{16-24}\) gives essentially the same result within 0.4 fm. The extraction of the effective range parameter \( r_{nn} \) has a greater dependence on the theoretical model since it represents the zero-energy variation of the n-n wave function within the range of the nuclear interaction. We shall not discuss here the measurement of this parameter which has little sensitivity to charge effects. The final-state interactions between the outgoing neutrons in \( \pi^- d \to \gamma nn \) has been discussed in the frame of potential theory\(^{17,21,24}\) with dispersion relations\(^{18,22,23}\) and inverse scattering methods.\(^{24}\) It is also important to compute the rescattering effects,\(^{18,20,21,23}\) where the pion collides elastically with the neutron before being absorbed radiatively by the proton. The effect of FSI is illustrated in Fig. 2, where the solid and dotted lines correspond to a RSC calculation and a separable potential respectively. The dashed-dotted line is the zero range approximation and the dashed line corresponds to a dispersion relation solution (the curves are from Ref. 24).

**Fig. 2. Sensitivity of the normalized photon spectrum to final-state interactions.**
The extraction of $a_{nn}$ results from a high-statistics, high-precision, simultaneous measurement of the photon spectrum from the capture of the pion from rest in the reactions $\pi^- d \rightarrow \gamma nn$ and $\pi^- p \rightarrow \gamma n$ using the SIN spectrometer. The experimental response function (720 keV FWHM) from $\pi^- p \rightarrow \gamma n$ is used to fold the theoretical spectrum (Fig. 3). The theoretical curve corresponds to a Bargmann potential for $a_{nn} = -18.6$ fm and $r_{nn} = 2.83$ fm. Our final result $a_{nn} = -18.5 \pm 0.4$ fm includes the theoretical uncertainty. In a recent analysis, the neutrons were detected in coincidence with the photon to rule out a systematic error in our previous experiment based on the analysis of the photon spectrum. The neutron time-of-flight spectrum (Fig. 3) is far less dependent on theoretical assumptions, since $a_{nn}$ is extracted in this case from the low energy n-n spectrum, where the theoretical uncertainty is negligible. The neutron spectrum gives $a_{nn} = -18.7 \pm 0.6$ fm.

![Fig. 3. Spectra for $\pi^- d \rightarrow \gamma nn$. Left: photon spectrum from $\pi^- capture in H and D. Right: neutron time-of-flight spectrum taken in coincidence with a photon.](image)

**THEORETICAL CALCULATIONS OF CSB**

The large mixing of the $\rho - \omega$ mesons is the dominant component of any CSB calculation. The construction of charge asymmetric potentials from $\rho - \omega$ mixing accounts for 80% of the measured $^3H - ^3He$ binding energy anomaly, and generates a difference of about 0.7 fm for $|a_{nn}| = a_{pp}$. The calculated charge asymmetric values, which relied on a decade old determination of $\langle \rho H \omega \rangle$, have been updated to incorporate an accurate value of the mixing extracted from colliding beams experiments on $e^+e^-$ annihilation into $\pi^+\pi^-$ pairs near the $\rho$ mass. The new value for the mixing matrix element

$$\langle \rho | H | \omega \rangle = 4520 \pm 600 MeV^2,$$

rescales previous estimates of CSB from $\rho - \omega$ mixing by almost a factor 1.2 accounting for most of the asymmetry of the N-N scattering lengths, as well as the binding energy differences in the $A=3$ system. The $\pi - \eta$ particle mixing is small due to the large mass difference of the $\pi$ and $\eta$ mesons. In addition, the $\pi - \eta/\eta'$ mixing generates a contribution of the opposite sign, the total effect being very small.
In the quark model, the $\rho - \omega$ mixing is expressed in terms of quark states as follows

$$\langle \rho | H | \omega \rangle = \frac{1}{2} \left( \langle \bar{u}u | H | \bar{u}u \rangle - \langle \bar{d}d | H | \bar{d}d \rangle \right).$$  \hspace{1cm} (13)

Since it is assumed that SU(2) is not spontaneously broken in the quark massless vacuum, $\langle 0 | \bar{u}u | 0 \rangle = \langle 0 | \bar{d}d | 0 \rangle$, the minus sign in Eq. 12 corresponds to $m_u - m_d < 0$.

The shifts in the scattering lengths are presented in Table 1. Secondary effects from two pion exchanges (TPE)\textsuperscript{30} and $\gamma \pi^0$ boson exchanges\textsuperscript{31} are also included. The effect from radiative corrections at the $\pi NN$ vertex (which induces a charge asymmetric vertex $g_{\pi nn}/g_{\pi pp}$) is very small.\textsuperscript{32} The results of a new calculation of CSB\textsuperscript{33} which includes the effect of the neutron-proton mass difference on the one pion exchange potential (OPE), contributing to charge III CSB potentials\textsuperscript{2}, are also shown in Table 1. I have rescaled the $\rho - \omega$ contribution of this reference with the new value (12). The difference of the $\rho - \omega$ contribution from these two calculations is due to the different couplings of the $\rho$ and $\omega$ mesons that the authors have used. The $\pi - \eta$ mixing contribution from Ref. (33) has been dropped because it was obtained by an inappropriate scaling argument. It is clear from Table 1 that a consistency check should be done on the TPE contribution. The rescaled value of the $\rho - \omega$ potential from Ref. 33 describes well the CSB effects of mirror nuclei for orbital configurations near closed shells.\textsuperscript{34}

Table 1

<table>
<thead>
<tr>
<th>Ref.</th>
<th>$\rho - \omega$</th>
<th>$\pi - \eta$</th>
<th>$\pi - \eta'$</th>
<th>TPE</th>
<th>OPE</th>
<th>$\gamma \pi^0$</th>
<th>TOTAL</th>
</tr>
</thead>
<tbody>
<tr>
<td>28\textsuperscript{a})</td>
<td>0.87</td>
<td>0.07</td>
<td>-0.10</td>
<td>0.30</td>
<td>0.19</td>
<td></td>
<td>1.33</td>
</tr>
<tr>
<td>33\textsuperscript{b})</td>
<td>1.16</td>
<td>0.11</td>
<td>-0.21</td>
<td>0.19</td>
<td></td>
<td></td>
<td>1.25</td>
</tr>
<tr>
<td>Exp.</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td>1.4 ± 0.6</td>
</tr>
</tbody>
</table>

a) Couplings $g_\rho$ and $g_\omega$ extracted from data
b) Couplings $g_\rho$ and $g_\omega$ from Bonn potential.

The CSB in the N-N interaction from short distance explicit quark effects have been examined in the context of the resonating group quark cluster model,\textsuperscript{35–37} and also in a QCD inspired model.\textsuperscript{38} The results are highly model dependent and very sensitive to the assumed values of the nucleon core $R_C$. The difference of scattering lengths for different models and assumptions is shown in Table 2. A significant contribution from short-distance quark effects is not compatible with the new value (12) for the $\rho - \omega$ mixing and the value (11) for the difference in the scattering lengths. It is important, however, to have agreement first on the CSB calculations of bosonic secondary effects, before deciding if explicit quark effects are needed at all to explain the data.
Table 2

<table>
<thead>
<tr>
<th>Ref.</th>
<th>$R_C$</th>
<th>TOTAL</th>
</tr>
</thead>
<tbody>
<tr>
<td>35</td>
<td>0.6</td>
<td>2–3.6</td>
</tr>
<tr>
<td>36</td>
<td>0.8</td>
<td>0.1</td>
</tr>
<tr>
<td></td>
<td>1.0</td>
<td>0.9</td>
</tr>
<tr>
<td>37</td>
<td>0.46</td>
<td>1.5–2.5</td>
</tr>
<tr>
<td></td>
<td>0.8</td>
<td>3.3–4.3</td>
</tr>
<tr>
<td>38</td>
<td>0.57</td>
<td>0.5</td>
</tr>
</tbody>
</table>

CONCLUSIONS

We have discussed the evidence for charge symmetry breaking in the nuclear interaction from bound mirror nuclei and low-energy N-N scattering experiments. The value $|a_{nn} - a_{pp}| = 1.4 \pm 0.6$, obtained from the reaction $\pi^- d \rightarrow \gamma nn$, is consistent with the $^3H-^3He$ binding energy difference and recent CSB rescaled $\rho - \omega$ calculations. With our present knowledge of the proton from low and high momentum transfer reactions, it should be possible to reduce the uncertainty in the Coulomb correction of $a_{pp}$. The origin of the significant discrepancy with the values for $a_{nn}$ obtained from deuteron breakup reactions should be found. New data for $nd \rightarrow pnn$ is needed and a full charge dependent Faddeev calculation for the n-n as well as the n-p final state interactions of this reaction. It is likely that the Los Alamos-Oak Ridge experiment from simultaneous fusion-fission sources would be done in a near future and that data on direct neutron-neutron collisions be obtained for the first time.

To conclude, we could ask what is the meaning of the isospin symmetry. The invariance under isospin rotations in charge space for nucleons is rephrased in present theories of the fundamental interactions as the invariance under identical rotations to left-handed (L) and right-handed (R) quark fields. In the standard model isospin is not a symmetry of the electroweak sector of the theory which only gauges the weak doublet. Isospin is not a natural symmetry of the sector of the strong interactions (QCD), which is basically invariant under independent L and R rotations. Isospin is a good symmetry due to the accidental smallness of the light quark masses, $m_u$ and $m_d$, with respect to the QCD scale $\Lambda$. The fundamental QCD symmetry at the level of the light quarks is $SU(2)_L \times SU(2)_R$ and not $SU(2)$. Isospin is not an exact symmetry, but a very useful one.

ACKNOWLEDGEMENTS

It is a pleasure to acknowledge the LMZ group, particularly C. Joseph and B. Gabioud, for our long and fruitful collaboration in the study of the charge symmetry breaking. I am indebted to S.A. Coon for discussions.
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CHARGE SYMMETRY BREAKING IN THE $^3$H-$^3$He SYSTEMS

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ABSTRACT

We solved a set of modified Faddeev equations including the Coulomb potential in addition to a realistic NN potentials with or without a three-nucleon potential, for various number of channels of the three-body system. After thirty two case studies, we deduced the $^3$He Coulomb energy of $683 \pm 4$ keV for the model independent value. Then the charge symmetry breaking and the charge independence breaking forces, fitted to the $^1S_0$ state NN scattering lengths are included in our Faddeev equation. Again, the model independent estimate is made for these effects, yielding $74 \pm 15$ keV. Other various small effects amount to $9$ keV, totaling $766 \pm 2$ keV over than the experimental $^3$H-$^3$He binding energy of $764$ keV.

INTRODUCTION

The validity of the isospin invariance in the strong interaction has attracted interest both in theory and experiment for a long time. The evidence of the charge independence breaking (CIB) and the charge symmetry breaking (CSB) of the NN interaction is obtained directly from the NN scattering experiments. For the NN system, if there is a CIB force, we have $V_{np} \neq V_{nn}$, but $V_{nn} = V_{pp}$, and if there is a CSB force, we have $V_{nn} \neq V_{pp}$. The difference of the scattering lengths $a_{np}, a_{nn}$ and $a_{pp}$, and the effective ranges $r_{np}, r_{nn}$ and $r_{pp}$ in spin-singlet state from the experiment tells us the existence of the CIB and the CSB force in the NN interaction. Because of the difficulty in the nn scattering experiment, the are very different from a group to a group, but the typical values of the scattering lengths are $2, 3, 4$ $a_{np} = -(23.715 \pm 0.015)$ fm, $a_{nn} = -(18.7 \pm 0.6)$ fm and $a_{pp} = -(17.15 \pm 0.15)$ fm. The big difference of $a_{np}$ and $a_{pp}$ data shows that the $np$ attraction is much stronger than the $pp$($nn$) interaction, thus CIB is big. The small difference of $a_{nn}$ and $a_{pp}$ data shows that the $nn$ interaction is slightly stronger than the $pp$ interaction.

The information about the CSB of NN interaction also comes from the mass difference of the mirror nuclei $\Delta E$. The main contributions to $\Delta E$ are from the Coulomb interaction of the pp pair. As pointed by Negele, the mass difference of $^{41}$Sc-$^{41}$Ca nuclei can not be explained by the Coulomb force and the Coulomb anomaly energy is about $580$ KeV. To give this Coulomb anomaly energy we need the CSB effect causing $a_{nn} - a_{pp} = -0.84$ fm.

In this article, we report our recent exact calculation of the CIB and CSB nuclear force effects on the binding energies of the simplest mirror nuclei $^3$H and $^3$He. The experimental binding energies are $E_{\exp}(^3\text{H}) = 8.482$ MeV and $E_{\exp}(^3\text{He}) = 7.718$ MeV. Thus, the mass difference is $\Delta = 0.764$ MeV. This mass difference has

* Presented by T. Sasakawa
been studied for a long time by different approaches in theory. However, due to the long range difficulty of the Coulomb force, $^3$He was treated mostly by the perturbational theory. All of these results showed that the Coulomb energy $E_c$ is smaller than $\Delta$ and the difference $\Delta - E_c$ is about 0.1 MeV. This difference has been thought to be due to the CSB effect.

Now, during this ten years, the exact calculation of $^3$H with two- and three-body forces based on the Faddeev approach has been completed. Based on this, here we perform exact calculations for $^3$He. By the exact calculation we mean that, we solve a modified 52-channel Faddeev equation including the Coulomb potential, a realistic two-nucleon potential, the Tucson-Melbourne (TM) three-nucleon potential exactly, taking account the total isospin $T=3/2$ components. Because of the CIB and CSB nuclear forces, the total isospin is no longer a good quantum number and it takes the value of $1/2$ and $3/2$. In this way, we can include the CIB and CSB force effects in $^1S_0$ state exactly in our calculation, although $T=3/2$ component is very small as we shall discuss in the below.

Among several sources of CSB nuclear force effects, the first thing that we should count is the purely electromagnetic effects, the most important of which is the Coulomb interaction, present in the pp system and missing in the nn system. There are some ways to accommodate the Coulomb force in the Faddeev equations for charged three-body systems. One of them is due to Sasakawa-Sawada. This modified equation has been used for $^3$He binding energy calculations. More in general, we consider a three-body force $W_{123} = W_{12,3} + W_{23,1} + W_{31,2}$. The modified Faddeev equations, including both the Coulomb and the three-body force takes the form

$$\begin{align*}
[E - H_0 - (V_\alpha + u^{c}_\alpha + u^{c}_{\alpha,\beta} + u^{c}_{\alpha,\gamma})]\Psi_\alpha &= [V_\alpha + (u^{c}_\alpha - u^{c}_{\beta,\alpha})]\Psi_\beta \\
+ [V_\beta + (u^{c}_\beta - u^{c}_{\gamma,\alpha})]\Psi_\gamma + W_{\beta,\gamma,\alpha}(\Psi_\alpha + \Psi_\beta + \Psi_\gamma)
\end{align*}$$

(1) 

\((\alpha, \beta, \gamma) = (1, 2, 3), (2, 3, 1), (3, 1, 2).\)

This equation is designed so that when the distance of the pair of particles becomes large, the Coulomb potential decreases as the inverse square of the distance. We use the Tucson-Melbourne two-pion exchange potential taken from Ref. 12 for the three-body force. The method of continued fractions is applied to solve these equations.

CIB AND CSB EFFECTS

Theoretically, CIB nuclear force is caused by several reasons. In the OBE NN potential, the main source of CIB is the mass difference of charged and neutral mesons.

In the long range part, the most important contribution comes from the pion exchange. The CIB potential is caused by the mass difference of charged and neutral
pions. It is given by
\[
(V_{np} - V_{nn})_{\Delta m_p} = \frac{g_{\rho}^2}{f^2} \frac{e^{-m_{\rho} r}}{4\pi r} \left[ 1 - \frac{\mu_+^2 e^{-\Delta m_{\rho} r}}{\mu_0^2} \right],
\]
where the coupling constant is \( f^2/4\pi \approx 0.079 \), the mass of neutral and charged pions are \( \mu_0 = 134.963(\text{MeV}) \) and \( \mu_+ = 139.567(\text{MeV}) \), respectively. Rho mesons will be important at the middle range of nuclear interaction. The CIB potential due to the mass difference of charged and neutral rho mesons is
\[
(V_{np} - V_{nn})_{\Delta m_{\rho}} = \frac{1}{2} \frac{g_{\rho}^2}{4\pi} \frac{e^{-m_{\rho} r}}{m_{\rho}^2} \left[ \left( e^{-\Delta m_{\rho} r} - 1 \right) \right.
\]
\[
\left. - \frac{K_{\nu}(K_{\nu} + 1)}{2} \left( \frac{m_{\rho}^2}{m_{\rho 0}^2} \left( e^{-\Delta m_{\rho} r} - 1 \right) \right) \right],
\]
where the coupling constants are \( g_{\rho}^2/4\pi \approx 2.4 \) and \( K_{\nu} \approx 3.7 \). The mass of neutral and charged rho mesons are \( m_{\rho 0} = 776.0(\text{MeV}) \) and \( m_{\rho +} = 770.4(\text{MeV}) \), respectively. To give the correct NN scattering length, we also introduce a phenomenological CIB potential. We assume that it takes the Woods-Saxon form, 
\[
V_{\text{phen}}(r) = V_0 / [(1 + e^{a(r - R)/a}] \]
where \( R = 0.5(\text{fm}) \) and \( a = 0.2(\text{fm}) \). \( V_0 \) is the strength of the potential and it is adjusted to fit the NN scattering data. We take \( V_0 \) as 6.5(\text{MeV}) in this work. This phenomenological potential affects \(|a_{np}| \) and \(|a_{nn}| \) about 2 ~ 3 fm.

Several authors have calculated the CIB effect in the triton binding energy (but without the full Coulomb effect).\(^{11,23} \) The average t-matrix \( \tilde{t} = (1/3)t_{np} + (2/3)t_{pp} \) for \(^3\text{He} \) and \( \tilde{t} = (1/3)t_{np} + (2/3)t_{nn} \) for \(^3\text{H} \) are used in Ref. 11 and the average potential \( \tilde{V} = (1/3)V_{np} + (2/3)V_{nn} \) is used in even-parity spin-singlet interaction for \(^3\text{H} \) calculation in Ref. 23. In the both cases, the isospin \( T = 3/2 \) component is omitted. On the other hand, we consider the isospin \( T = 3/2 \) component and include the CIB force in \(^1\text{S}_0 \) state exactly in our calculation.

The theoretical investigations of the CSB other than the Coulomb origin in the NN interaction, often give conflicting results. Therefore, the theoretical situation in the last decades has been much uncertain than the experimental one. The best studied origins are \( \pi^0 - \eta \), \( \rho^0 - \omega \) mixing and the \( n-p \) mass difference in the OPE potential.\(^{1,24,25} \) The CSB potential due to \( \rho^0 - \omega \) mixing is given by the Eq.(4) in Ref. 24. In the case of \(^1\text{S}_0 \) state, it can be reduced to the following form,
\[
\Delta V_{\rho\omega} = \frac{g_{\rho} g_{\omega}}{4\pi} \frac{(\rho^0|H|\omega)}{2m_{\omega}} (\tau_1 + \tau_2) e^{-m_{\tau} r} (1 + \beta(\frac{2}{m_{\omega} r} - 1)\frac{m_{\omega}^2}{M^2}),
\]
where \( \beta = (1/4)(2K_{\rho}K_{\omega} + K_{\rho}K_{\omega} + 1) \) is the interaction constant, \( m_{\omega} = (1/2)(m_{\omega} + m_{\rho}) \) is the average mass of the mixed mesons and \( M \) is the mass of a nucleon. We
approximate $m_\omega \simeq m_\rho$ to get Eq. (2). Other parameters are $m_\rho = 769$ MeV, $m_\omega = 782.6$ MeV, $g_\rho = 3.46$, $g_\omega = 15.58$, $K_\rho = 6.1$ and $K_\omega = 0$. The electromagnetic interaction matrix element is given by the experiment as $\langle \rho^0|H|\omega \rangle = -3850$ MeV$^2$. The expression of (2) is exactly the same as the Eq. (8) in Ref. 25. However, the interaction constant $\beta$ is different [comparing the expression of (9) for $\beta$ in Ref. 25]. In any way, we take $\beta$ as a parameter in our calculations. The CSB potential due to $\pi^0\eta$ mixing in $^1S_0$ state is given by

$$\Delta V^{\pi\eta} = -\frac{g_\pi g_\eta}{4\pi} \frac{\langle \pi^0|H|\eta \rangle}{m_\eta^2 - m_\pi^2} (\tau_1 + \tau_2) \frac{1}{4\pi} \left( m_\pi^2 e^{-m_\pi r} - m_\eta^2 e^{-m_\eta r} \right),$$

where the electromagnetic interaction matrix element is $\langle \pi^0|H|\eta \rangle = -4200$ MeV$^2$. Other parameters are given by $m_\pi = 134.96$ MeV, $m_\eta = 548.8$ MeV, $g_\pi = 13.55$, and $g_\eta = 5.09$. The mass difference of n-p also cause a CSB force in OPEP. For $^1S_0$ state, it is

$$\Delta V^{np} = \frac{g_\pi^2}{4\pi} \delta(\tau_1 + \tau_2) \frac{m_\pi^2}{4M^2} \frac{e^{-m_\pi r}}{r},$$

where $\delta = (M_n - M_p)/(M_n + M_p)$.

The CSB effect was investigated by Brandenburg's group$^{,26}$ and Gibson et al.$^{27}$ They estimated the effect of CSB nuclear force by a model independent form factor method, which yields a smaller Coulomb effect of 638 keV than the direct calculation. On the contrary, as in the case of CIB, we also include the CSB force exactly in our calculation.

Before showing our results, we briefly discuss the three-body channel with the total isospin $T=3/2$ component. To solve the Faddeev equations, the Faddeev amplitudes are projected onto a complete set of channel wave functions which represent the orbital angular momentum and spin-isospin variables. We express the channel basis states as $|\alpha\rangle = \{(LS)J, (ls)J_0 M_0; (It)M_T\}$. We divided them into four classes according to the total angular momentum $J$ of a pair particles. They are 6 channel ($^1S_0$ and $^3S_1 - ^3D_1$), 28 channel ($J \leq 2$), 38 channel ($J \leq 3$) and 52 channel ($J \leq 4$). The isospin functions are shown in Table 1. For the NN interactions, the following ones were used: the Reid soft-core (RSC),$^{26}$ de-Tourreil-Rouben-Sprung (TRS),$^{29}$ Paris (PARIS),$^{30}$ Argonne (AV)$^{31}$ and Bonn (BONN)$^{32}$ potentials. We should remark that all of these realistic NN interactions are charge independent and the NN scattering lengths are fitted either to $a_{pp}$ or $a_{np}$ as seen in the second column of Table 2. The difference in the assumed scattering lengths of these realistic potentials have been based on the thoughts that the difference of the experimental value may be caused by the difference in the way of experiments or due to the experimental error, and need not to be taken into account seriously. With above CIB and CSB potentials, the scattering lengths become as shown in Table 2. The charge symmetry breaking shows that this traditional thoughts must be changed.
NUMERICAL RESULTS

Now we discuss our calculated results. First, let us consider the Coulomb effect. We solve the modified Faddeev equations (1) and obtain the binding energies of the trinucleon systems for various realistic NN interactions and the Tucson-Melbourne (TM) three-nucleon force. Specifically, we calculated 32 cases: RSC6, 28; [RSC+TM]6, 28; 2NP6, 28, 38, 52 for 2NP=AV, PARIS, TRS and BONN; [2NP+TM]6, 28, 38, 52 for 2NP=AV, PARIS, TRS. Our results are listed in Table 3. We see that if only the two-body interaction is considered, both $^3$H and $^3$He are underbound about 1.0 MeV. Taking the cutoff mass $\Lambda = 700$ MeV for $\pi NN$ form factor in TM, we obtain the binding energies which is very close to the experimental one for three-nucleon systems. The Coulomb effect $E_c$ increases as the binding energies of the systems increase. We find that there is a remarkable linear relationship between the binding energies of $^3$H and $^3$He,

$$E(^3\text{He}) = E(^3\text{H}) \times 0.9617 - 0.3579 \pm 0.0062 \text{ MeV}$$

Using this line, and putting $E_{\text{exp}}(^3\text{H})$, we obtain a model independent value for $E(^3\text{He})$: $E_{\text{m.i.}}(^3\text{He}) = 7.799 \pm 0.006$ MeV and the model independent Coulomb energy: $E_{c,\text{m.i.}} = 0.683 \pm 0.006$ MeV. The Coulomb energy anomaly of $^3$H and $^3$He $(E_{\text{exp}}(^3\text{H}) - E_{\text{exp}}(^3\text{He})) - E_{c,\text{m.i.}}$ is $0.081 \pm 0.006$ MeV and this value is what we have now to obtain theoretically. It has been suggested that it mainly comes from the CSB force.

We have solved the Coulomb modified Faddeev equation adding the CIB forces. Table 4 shows the calculated results of binding energies of $^3$H and $^3$He with CIB forces. Comparing with Table 3, we draw a conclusion that the CIB effect causes the same mass shift for both $^3$H and $^3$He. Thus CIB effect has almost no contribution to the mass difference of $^3$H and $^3$He. In RSC, TRS and PARIS potentials, the CIB nuclear force effects make increase the absolute value of the scattering length of $a_{np}$, thus CIB effects let increase the binding energy about 0.105 MeV and 0.115 MeV (with three-body force). In AV potential, the value of $a_{np}$ are kept, and as a result $a_{nn}$ and $a_{pp}$ are decreased, resulting the decrease of the binding energy of about 0.200 MeV and 0.220 MeV (with three-body force). In Table 5, we calculated the CIB effects on the mass difference by the perturbational way. The long range CIB force caused by pions is important, and contributes about 60% to the CIB effect.

Although the CSB of nuclear force is small, it contributes to the mass difference of $^3$H and $^3$He. From the experiment, the scattering length difference $\Delta a = |a_{nn} - a_{pp}|$ is $\Delta a_{\text{exp.}} = 1.55 \pm 0.75$ fm. We adjust the parameter $\beta$ in Eq. (2) to let $\Delta a = 1.5$ fm for various realistic NN interactions. In Table 2, we show the scattering lengths that are obtained after incorporating the CIB and CSB potentials.

The binding energies of $^3$H and $^3$He are calculated for various NN potentials and TM three-body force further adding the CSB potential. The results are listed in Table 6. Also the perturbational estimate of the contributions are given in Table 5.
As seen in Table 5, $\rho^0-\omega$ mixing amounts almost to all the CSB effects. We deduce a model independent CSB effect from 32 case studies. As in the case of Coulomb force, a very good linear relationship is obtained as before.

$$E(^3\text{He}) = E(^3\text{H}) \times 0.9507 - 0.3387 \pm 0.0092 \text{ MeV}$$

From this relationship we obtain the binding energy for $^3\text{He}$ as $E_{\text{m.i.}}^{\text{CSB}}(^3\text{He}) = 7.725 \pm 0.009$ MeV. Comparing with $E_{\text{m.i.}}(^3\text{He})$, we find that the CSB effect is $\delta E_{\text{m.i.}}(\text{CSB}) = 0.074 \pm 0.015$ MeV. It should be compared with the the Coulomb energy anomaly of $0.081 \pm 0.006$ MeV. We see that the CSB correction is very close to the Coulomb energy anomaly.

There are some other small effects such as (1) the finite size effect due to the nucleon electromagnetic form factor, (2) the n-p mass difference in kinetic energy, (3) the vacuum polarization potential acting on the pp pair, (4) the magnetic interactions and (5) the momentum dependent electromagnetic interaction due to the relativistic corrections. We make an estimation using the wave function generated from TRS-52(TM700), PARIS-52(TM700) and AV-52(TM700) and averaging the results. For the finite size of proton, we consider the charge form factor in momentum space $G_E^p(q^2) = (\Lambda^2/(q^2 + \Lambda^2))^2$ with $\Lambda = 840$ MeV. The effect of the neutron charge form factor is very small and it can be neglected. The magnetic moment form factor can be given by an approximate scaling relation $G_M^p/\mu_p = G_M^n/|\mu_n| = G_E^p$.

Finally, we discuss the percentage of the total isospin $T=3/2$ component. Taking AV-52(TM700) as an example, we obtained zero and $1.0 \times 10^{-3}$ for $^3\text{H}$ and $^3\text{He}$, respectively, with only the Coulomb forces. Including CIB and CSB forces, we have $4.0 \times 10^{-4}$ and $4.0 \times 10^{-3}$ for $^3\text{H}$ and $^3\text{He}$, respectively.

**CONCLUSION**

Our results are summarized in Table 7. It shows that other effects that we did not take into account should be small.

We thank Dr. M. Maruyama for his useful discussions. The computer calculation for this work has been financially supported in part by the Grant-in-Aid of the Ministry of Education, Science and Culture of Japan, Research Center for Nuclear Physics, Osaka University, and the Cyclotron Radioisotope Center, Tohoku University.
REFERENCE


12. T. Sasakawa and S. Ishikawa, Few-body system 1, 3(1986); 1, 143(1986).


TABLE CAPTIONS

TABLE 1. The isospin state of trinucleon systems, where the upper sign is for $^3$H and the lower one is for $^3$He.

TABLE 2. The scattering lengths of NN interaction is calculated for various two-nucleon potentials. $a_{NN}$ is the scattering length obtained without CIB and CSB nuclear forces. $a_{nn}$, $a_{pp}$ and $a_{np}$ are the scattering lengths obtained with CIB and CSB nuclear force.

TABLE 3. The $E(3H)$ and $E(3He)$ are the binding energies of $^3$H and $^3$He, respectively. $E_c$ is the Coulomb effect on the mass difference in trinucleon systems. The number after the potentials are the number of three-body channels. The number inside the bracket is the results with the three-body force.

TABLE 4. The $E_{CIB}(3H)$ and $E_{CIB}(3He)$ are the binding energies of $^3$H and $^3$He, respectively, for various realistic NN interactions and three-body channels with CIB forces. $\Delta E(CIB)$ is the mass difference in trinucleon systems. The number inside the bracket is the results with the three-body force.

TABLE 5. The perturbation calculation of CIB forces and CSB forces effects on the binding energy of $^3$H for AV-52(TM700).

TABLE 6. The $E_{CSB}(3H)$ and $E_{CSB}(3He)$ are the binding energies of $^3$H and $^3$He, respectively, for various realistic NN interactions and three-body channels with CIB and CSB forces. $\Delta E(CSB)$ is the mass difference of trinucleon systems. The number inside the bracket is the results with the three-body force.

TABLE 7. The contributions of charge asymmetric effects to the $^3$H and $^3$He binding energy difference.
### TABLE 1

| State | I | I_z | Pair   | \( |I_{I_z}; t m_t; M_T(12, 3)\) |
|-------|---|-----|--------|----------------------------------|
| No. 1 | 1 | 0   | np     | \( |10; \frac{1}{2} \pm \frac{1}{2}; \pm \frac{1}{2}(12, 3)\) |
| No. 2 | 1 | ±1  | pp,nn  | \( |1 \pm 1; \frac{1}{2} \mp \frac{1}{2}; \pm \frac{1}{2}(12, 3)\) |
| No. 3 | 0 | 0   | np     | \( |00; \frac{1}{2} \pm \frac{1}{2}; \pm \frac{1}{2}(12, 3)\) |

### TABLE 2 The Unit for Scattering Length : fm

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<th>(a_{np})</th>
<th>(a_{nn})</th>
<th>(a_{pp})</th>
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### TABLE 3 The Unit for Energy : MeV

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<th>(E(3\text{He}))</th>
<th>(E_c)</th>
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<td>RSC-28(TM700)</td>
<td>7.233 (8.064)</td>
<td>6.599 (7.403)</td>
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<td>TRS-52(TM700)</td>
<td>7.543 (8.420)</td>
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<td>0.648 (0.680)</td>
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<td>PARIS-52(TM700)</td>
<td>7.630 (8.273)</td>
<td>6.979 (7.602)</td>
<td>0.651 (0.671)</td>
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<tr>
<td>AV-52(TM700)</td>
<td>7.673 (8.367)</td>
<td>7.026 (7.700)</td>
<td>0.647 (0.667)</td>
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<td>BONN-52</td>
<td>8.317</td>
<td>7.629</td>
<td>0.688</td>
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### TABLE 4 The Unit for Energy : MeV

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<th>(E_{\text{CIB}}(3\text{H}))</th>
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<td>PARIS-52(TM700)</td>
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<td>BONN-52</td>
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TABLE 5 The Unit for Energy : keV

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<td>ρ mass difference</td>
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<td>phenomenon</td>
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TABLE 6 The Unit for Energy : MeV

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<th>$E_{CSB} (^3\text{He})$</th>
<th>ΔE(CSB)</th>
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<td>7.084 (7.714)</td>
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<td>AV-52(TM700)</td>
<td>7.529 (8.213)</td>
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<td>0.703 (0.730)</td>
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<td>BONN-52</td>
<td>8.133</td>
<td>7.381</td>
<td>0.752</td>
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TABLE 7 The Unit for Energy : keV

<table>
<thead>
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<th>Charge asymmetry effects</th>
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<td>Static Coulomb</td>
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</tr>
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<td>CIB and CSB forces</td>
<td>74±15</td>
</tr>
<tr>
<td>Proton finite size effect</td>
<td>-33±3</td>
</tr>
<tr>
<td>K. E. due to n-p mass diff.</td>
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</tr>
<tr>
<td>Vacuum polarization</td>
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<tr>
<td>Magnetic interaction</td>
<td>6±2</td>
</tr>
<tr>
<td></td>
<td>pp 20±1</td>
</tr>
<tr>
<td></td>
<td>nn 14±1</td>
</tr>
<tr>
<td>Momentum dep. e. m.</td>
<td>11±3</td>
</tr>
<tr>
<td><strong>Total (theory)</strong></td>
<td>766±30</td>
</tr>
<tr>
<td><strong>Experiment</strong></td>
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</table>
CHARGE SYMMETRY BREAKING - AN OVERVIEW

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ABSTRACT

Diverse aspects of charge symmetry breaking (CSB) are reviewed. I argue that the mass difference between up and down quarks is responsible for all observed non-electromagnetic CSB.

INTRODUCTION

The purpose of this talk is to provide a common theme for the different aspects of charge symmetry breaking (CSB) that can be studied. My presentation is drawn from a lengthy review in preparation for Physics Reports.¹

Consider the quark mass difference (QMD) hypothesis: the difference between the masses of the up and down quarks plus electromagnetic interactions amongst quarks causes all CSB.

The long-ranged electromagnetic interaction is not very sensitive to whether nuclear matter comes in the form of hadrons or quarks. Thus the currently testable part of the hypothesis involves the quark mass differences, hence the name QMD. There is much evidence to support this hypothesis.

1. The up-down quark mass difference \((m_d - m_u)\) is the origin of a vast number of hadronic mass differences. If two hadrons are related by switching a down to an up quark, the replacement \(u \rightarrow d\) leads to an increase of the hadronic mass.¹² Moreover, within theories of confined quarks, the value of \(m_d - m_u\) that yields mass differences in agreement with experiment also reproduces the value of the \(\rho\omega\) mixing matrix element.³

2. The existence of \(\rho\omega\) mixing allows the \(\omega\) a two pion decay mode. Many \(e^+e^- \rightarrow \pi\pi\) experiments clearly find this mode. A recent experiment has especially good statistics.⁴

3. The CSB effect discovered in a recent TRIUMF 477 MeV measurement of elastic neutron-proton (n-p) scattering⁵ seems to be dominated by a term in the one pion exchange potential proportional to \(M_n - M_p\).⁶ (The quark mass difference is the origin of the nucleon mass difference, so this is a success of the QMD.)

In addition, \(\rho - \omega\) mixing influences nucleon-nucleon scattering when nucleons exchange a mixed \(\rho\omega\) meson. This is a small effect for the published TRIUMF n-p experiment, but could be observable at other energies or scattering angles.⁶

4. \(\rho - \omega\) mixing causes the interaction between two neutrons to be more attractive than the one for two protons.⁷ The potential \(\Delta V = V_\rho^{\omega n} - V_\rho^{\omega p}\) is a medium-ranged force. Thus, it is not very sensitive to short distance uncertainties.⁸ The magnitude of the effect on the scattering length difference is in agreement with current observations.¹⁸⁻¹⁰

5. The interaction \(\Delta V\) is large enough to account for the binding energy difference between the three-body nuclei.⁸

6. Moreover, the same \(\Delta V\) explains¹¹ most of the Nolen-Schiffer¹² anomaly.
The net result is that the QMD hypothesis seems to be valid. It would be very interesting to confirm the influence of $\rho$-$\omega$ mixing in elastic $n$-$p$ scattering experiments.$^{13}$

**THE MASS DIFFERENCE BETWEEN UP AND DOWN QUARKS**

For me, the single most compelling evidence for believing that the mass of the down quark $m_d$ is greater than the mass of the up quark ($m_u$) is that the neutron ($ddu$) is more massive than the proton ($uud$). This mass difference overcomes the influence of electromagnetic interactions which tend to make the positively charged proton more massive than the neutron.

A word of caution about quark masses is in order here. The relation between quark masses used in hadronic descriptions and the fundamental current quark masses that appear in the quantum chromodynamical (QCD) Lagrangian is tricky.$^{14}$ Indeed, it is often argued that one can only determine ratios of quark masses. My purpose here is more phenomenological. It is to seek masses that can consistently be used in describing a wide variety of hadronic and nuclear processes. Indeed, it turns out that the role of the quark masses in understanding different processes is model independent, even though specific values may vary widely.

I want to illustrate the main ideas with the examples most relevant for the present CSB discussion. One starts with a quark model Hamiltonian

$$H = V_{\text{conf}} + \sum_{i=1}^{3} (p_i^2 + m_i^2)^{1/2} + \sum_{i>j} (\alpha Q_i Q_j + \bar{\lambda}_i \cdot \bar{\lambda}_j \alpha_s) S_{ij} ,$$

where

$$S_{ij} = -\frac{\pi}{2} \, \delta(\vec{r}) \, \frac{16}{3} \, \frac{\bar{s}_i \cdot \bar{s}_j}{m_i m_j} + ...$$

For simplicity I assume harmonic oscillator wave functions with oscillator parameter $b = 0.62$ fm and $\alpha_s$ chosen to fit $2/3$ of the $\Delta$-nucleon mass difference. This allows the nucleon to have space available for the necessary pion cloud.$^{15,1}$ The average quark mass value is chosen to be 330 MeV. Setting $m_d - m_u = 2.80$ MeV allows a reproduction of the experimental value $M_n - M_p = 1.29$ MeV.

It is interesting to look at the separate contributions. The difference in expectation values of the sum of mass plus kinetic energy for down and up quarks is 2.11 MeV. This is less than 2.80 MeV because the average kinetic energy of a quark decreases as its mass increases. The difference in the gluon exchange energy is -0.244 MeV. This repulsive (positive) energy term decreases with increasing quark mass. This term tends to make the proton more massive. The well-known electromagnetic effect has the same tendency. Its contribution is -0.573 MeV. The sum of the three terms (2.11-0.244-0.573) is 1.29.

Taken by itself, this example seems ludicrous. I have used one parameter $m_d - m_u$ to "explain" one datum for $M_n - M_p$. However, choosing $m_d - m_u$ in a given model allows one to describe a full range of baryon and mesonic mass differences. A few references are listed$^{16}$. 


RHO-OMEGA MIXING

The most relevant example of meson mixing is that which occurs between the $\rho^0$ and $\omega$ mesons. If charge symmetry holds, the $\rho^0$ is a member of an isotriplet ($I=1$) of $\rho$ mesons and the $\omega$ an isospin singlet. The violation of charge symmetry implies $[H, P_{cs}] \neq 0$. Then recalling that for neutral systems the charge symmetry operator $P_{cs}$ acts as a charge parity [17] operator: $P_{cs} | I, I_{30} \rangle = (-1)^I | I, I_{3} = 0 \rangle$, we find

$$\langle \rho^0 | H | \omega \rangle = \frac{1}{2} \langle \rho^0 | [H, P_{cs}] | \omega \rangle ,$$

(3)

Thus the $\rho^0$ and $\omega$ physical states are mixtures of bare $\rho^0$ and $\omega$ states if CSB occurs.

The $\rho-\omega$ mixing is observed in the reaction $e^+e^- \rightarrow \pi^+\pi^-$. The experimental signature is based on the existence of two interfering amplitudes. The dominant term is $\pi^+\pi^-$ emission from a real $\rho$ meson. The CSB enters via the conversion of an $\omega$ meson into a $\rho$ meson. These two amplitudes are coherent so $\rho-\omega$ mixing can be observed as an interference effect. Indeed, one observes a sharp “knee” in the cross section at the energy corresponding to the $\omega$ meson mass, Ref. 4. The relevant mixing parameter is the $\rho$ self-energy term

$$\mu_{\rho\omega}^2 = 2m_\rho \langle \rho^0 | H | \omega \rangle .$$

(4)

The recent Novosibirsk experiment finds

$$\mu_{\rho\omega}^2 = -4.5 \pm 0.06 \times 10^{-3} \, \text{GeV}^2 .$$

(5)

Using eqs. (4 and 5) one obtains

$$\langle \rho^0 | H | \omega \rangle = -2.9 \pm 0.4 \, \text{MeV} .$$

(6)

How does the value of $-2.9$ MeV compare with expectations based on electromagnetic processes and quark mass differences? The annihilation process $\rho^0 \rightarrow \gamma \rightarrow \omega$ contributes \cite{18} $0.43 \pm 0.04$ MeV, which is obtained using the known $e^+e^-$ decay widths of the $\rho$ and $\omega$ mesons. Thus one wants to understand the remainder of $-3.3 \pm 0.4$ MeV.

One can use the quark model of Eqns. 1 and 2 to understand this value of $-3.3$ MeV. That Hamiltonian can be applied to the $\rho^0$ and $\omega$ mesons. The mesons are treated as quark antiquark pairs, which in the absence of CSB have good isospin. Then one writes the unperturbed basis states as

$$|\rho^0 \rangle = (|u\bar{u} \rangle - |d\bar{d} \rangle)/\sqrt{2} ,$$

$$|\omega \rangle = (|u\bar{u} \rangle + |d\bar{d} \rangle)/\sqrt{2} .$$

(7)

The use of first-order perturbation theory leads to the result

$$\langle \rho^0 | H | \omega \rangle = -\Delta m \left[ \frac{\partial K}{\partial m} + \frac{40}{9} \frac{\alpha_s \pi}{m^3} \langle \psi | \delta^{(3)}(\vec{r}) | \psi \rangle \right] + \frac{\alpha}{6} \langle \psi | \frac{1}{r} | \psi \rangle .$$

(8)
with $K$ as the quark mass plus kinetic energy operator and $\Delta m = m_d - m_u$.

The different quark masses induce changes in the kinetic energy operator that are supplemented by electromagnetic and gluonic interactions of $m_d - m_u$. Note that the negative sign of the mixing matrix element is reproduced by a positive value of $m_d - m_u$, provided the interaction terms do not overcome the change in the kinetic energy. This is so in all the models employed.

A simple computation leads to a $\rho^0 - \omega$ mixing matrix element of $-3.2$ MeV, in very good agreement with the experimental value of $-3.3$ MeV. Thus there is a consistent explanation of both the $np$ mass difference and $\rho^0 - \omega$ mixing.

The values presented here are illustrative and not meant to be an official tabulation. This is because of the model dependence. For example, Godfrey and Isgur (GI) find that using $\Delta m = 8.8$ MeV is necessary to reproduce the experimental value of the $\rho - \omega$ mixing matrix element. The difference between the GI result and the present value is discussed in Ref 1. However the specific value is not relevant. The main point is consistency. Once a model of confinement is chosen, a single value of $m_d - m_u$ allows one to explain mass difference and meson mixing data.

$p - \omega$ MIXING AND $^1S_0$ NUCLEON-NUCLEON SCATTERING

Next consider the influence of meson mixing on the nucleon-nucleon force. A virtual meson on its way from one nucleon to another can undergo isospin mixing. The largest of these exchanges is that due to $\rho^0 - \omega$ mixing. This is because of the small $\rho^0 - \omega$ mass difference. The exchange of mixed $\pi \eta, \pi \eta'$ mesons is smaller, if appropriately chosen strong interaction coupling constants are used.

Consider the $nn$ and $pp$ $^1S_0$ states. Coon et al. have derived a convenient and accurate form for the $p - \omega$ exchange contribution to the difference between the $nn$ and $pp$ potentials:

\[
\Delta V^{p\omega} = \frac{g_{p\omega}}{4\pi} \left| \frac{H_{em} \omega}{2m_v} \right| e^{-m_v r} \left[ 1 + \beta \left( \frac{2}{m_v r} - 1 \right) \frac{m_v^2}{M^2} \right].
\]

The form makes it clear that this potential is not sensitive to the $p$ mass, which is still imprecisely known; the radial dependence is mainly an exponential; and, isolates the relativistic corrections of order $1/M^2$. The constant $\beta$ in eq. (9) is

\[
\beta = \frac{1}{2} \left[ k_s k_v + \frac{1}{2} (k_s + k_v) \right],
\]

where $k_s$ and $k_v$ are the isoscalar and isovector tensor coupling constants. Using the published coupling constants yields $\beta = 1.22$.

The exponential feature of eq. (9) is essential. The potential is actually of medium range, essentially that of a $\sigma$-exchange term ($M_\sigma = 550$ MeV).

Coon and Barrett compute $a_{nn}$ and $a_{pp}$ with the Reid soft core (RSC) or de Toureil Rouben Sprung (dTRS) potentials supplying the charge-independent strong interaction. The results (with $\beta = 1.22$) are

\[
|a_{nn}| - |a_{pp}| = 0.9 \text{ fm (RSC)} = 1.35 \text{ fm (dTRS)}.
\]
These values may now be compared with experimental values. This represents considerable progress.

$^1S_0$ SCATTERING LENGTHS - EXPERIMENTAL

The measured value of the proton-proton scattering length has long been precisely known. However, it is necessary to remove the effects of the Coulomb interaction to make comparisons with the neutron-neutron system. We argue that the proton-proton scattering length $a_{pp}$ has the value $a_{pp} = -17.3 \pm 0.3$ fm. The uncertainty of 0.3 fm is due to the model dependence encountered in removing the effects of Coulomb scattering.

This is a much smaller uncertainty than obtained by the non-local phase equivalent potentials of Sauer. We claim that these potentials are ruled out. The essential point is the modern notion that wave functions are related to physical channels such as nucleon-nucleon, delta-delta, six-quark bag etc. Thus one can speak of the baryonic charge density of the physical two baryon wave function. We find that extreme variations of this charge density yield an uncertainty of at most 0.3 fm. Furthermore, the non-local potentials causing larger uncertainties lead to very unusual $^3$He charge densities in disagreement with data.

Next turn to the neutron-neutron scattering length. The measurements of the photon spectrum from the reaction $D(n, \gamma)n$ gave

$$a_{nn} = -18.50 \pm 0.42 \text{ fm}$$

Recently, a kinematically overdetermined measurement of a neutron spectrum in coincidence with a photon gave $a_{nn} = -18.7 \pm 0.6$ fm.

The reaction $D(n, nn)p$ has been studied in many kinematic situations. Typically the data are analyzed with the Faddeev equations using simple separable $S$ wave $NN$ forces. There are uncertainties due to off-energy-shell effects and also to the possible influence of three-body forces. The different kinematic regions provide different between extracted values of $a_{nn}$. It has long been known that the nucleon-nucleon final state interaction is sensitive to variations in the short range force. Thus the reaction $D(\pi^- , \gamma)n$ provides the best value for $a_{nn}$. We quote $a_{nn} = -18.5 \pm 0.3$ fm.

Combining the $nn$ and $pp$ information yields $|a_{nn}| - |a_{pp}| = 1.2 \pm 0.4 \text{ fm}$, which is remarkably close to the earlier predictions of Coon and Co.

THE $^3$He-$^3$H BINDING ENERGY DIFFERENCE

The binding energy difference is given by $B(^3\text{H}) - B(^3\text{He}) = 0.764 \text{ MeV}$. The neutron rich $^3$H is more deeply bound than $^3$He which contains two protons. Indeed, the electromagnetic repulsion between the two protons is responsible for most of the difference. But how large are the Coulomb and other electromagnetic effects? This question was studied by Brandenburg et al. and most recently by Coon and barrett (CB). The key point is to use measured electromagnetic form factors of the two nuclei. CB find that electromagnetic effects account for $693 \pm 19 \pm 5$ keV. The first uncertainty is due to experimental errors in determining the form factor, the second is a rough guess at the model dependence of the extraction procedure. Thus $71 \pm 19 \pm 5$ keV remain to be explained by CSB effects.
Coon and Barrett make a model independent evaluation of the influence of $\Delta V$ and obtain $89 \pm 14$ keV. This agrees with the experiment.

**NOLEN-SCHIFFER ANOMALY IN MIRROR NUCLEI**

Mirror nuclei are members of an isospin doublet. As a first approximation, one may think that the two nuclei are related by changing a proton in one nucleus into a neutron of the same wave function. More formally the two states are related by the charge symmetry rotation.

If CS holds the two states are degenerate. Thus information about the CSB nuclear force can be inferred by comparing binding energy differences of mirror nuclei with calculated electromagnetic effects. It has long been known$^{32,33}$ that the electromagnetic effects do not entirely account for the experimentally measured binding energy differences.

For nuclei heavier than $A=3$ the binding energy difference should mainly be determined by the electromagnetic interaction of the odd nucleon. In a systematic study over a wide range of nuclei Nolen and Schiffer$^{33}$ considered the direct electromagnetic effects such as the Coulomb interaction, electromagnetic spin-orbit interaction and the change in kinetic energy due to the neutron-proton mass difference.

The dominant term of Nolen and Schiffer (NS) is due to the Coulomb interaction of the last nucleon with the core. NS also included a variety of correction terms and found that the calculated electromagnetic energy differences were too low by about 7%. This is similar to the situation in the three-body nuclei, in which electromagnetic effects account for about 680 of the necessary 760 keV.

The Nolen-Schiffer paper inspired much theoretical work. More detailed nuclear structure efforts aimed to examine effects such as isospin mixing, correlations and core polarization were launched. We reviewed$^1$ the attempts to explain the anomaly using only computed electromagnetic effects. Such terms fall short of accounting for the experimental values. However, the often-quoted shortfall-value of 7% could be decreased to about 5% if different effective interactions and all of the effects of core polarization are carefully computed.

The success of the $\rho^0-\omega$ mixing term in accounting for the nucleon-nucleon CSB scattering length difference and also the $A=3$ binding energy difference is noted above. Blunden and Iqbal$^{34}$ took the next step and computed the binding energy difference of the mirror nuclei with $A = 11, 13, 15, 17, 27, 29, 31, 33, 39$ and $41$. These workers included nucleon-nucleon CSB potentials derived from $\rho^0-\omega$ and $\pi^0-\eta$ mixing, and the influence of neutron-proton mass difference in OPEP and TPEP. We note that the effects of form factors and the exchange term which reduce the size of the effect have been included by Blunden and Iqbal. As shown in Table 1, the effect of $\rho^0-\omega$ mixing dominates the nuclear CSB. CSB nucleon-nucleon forces explain about 85% of the discrepancy between the Hartree-Fock theory (with the Skyrme II interaction of Ref. [35] and experiment. The $\rho^0-\omega$ contribution increases with increasing mass number, which is seen experimentally.

We regard the agreement as very good. Effects such as the non-monopole core polarization$^{36}$, quark and relativistic effects (see below), the uncertainties in the effective interaction and possible CSB three-nucleon forces can be invoked to describe the small remaining discrepancies. Thus we may conclude that the NS anomaly is well explained. A theoretically derived CSB meson exchange potential, determined by quark mass differences, constrained by $e^+e^-$ annihilation and nucleon-nucleon
scattering data accounts for the binding energy anomaly in the mirror nuclei ranging from \( A = 3 \) to \( A = 41 \).

Table I.\(^a\) Contributions to the unexplained binding energy difference in keV. The column SkII are the values required after known electromagnetic and nuclear structure effects are removed (from Sato's calculation with the Skyrme-II force). The \( \rho^0 - \omega \) contribution has been rescaled to incorporate the current value of the mixing matrix element.

<table>
<thead>
<tr>
<th>A</th>
<th>orbital</th>
<th>( \rho^0 - \omega )</th>
<th>total</th>
<th>SKII</th>
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<tr>
<td>15</td>
<td>( p^{-1}_{3/2} )</td>
<td>161</td>
<td>189</td>
<td>190</td>
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<tr>
<td></td>
<td>( p_{1/2}^{-1} )</td>
<td>200</td>
<td>256</td>
<td>290</td>
</tr>
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<td>17</td>
<td>( d_{5/2}^{-1} )</td>
<td>116</td>
<td>129</td>
<td>190</td>
</tr>
<tr>
<td></td>
<td>( l_s_{1/2}^{-1} )</td>
<td>192</td>
<td>228</td>
<td>210</td>
</tr>
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<td>( d_{3/2}^{-1} )</td>
<td>169</td>
<td>223</td>
<td>270</td>
</tr>
<tr>
<td>39</td>
<td>( l_s_{1/2}^{-1} )</td>
<td>256</td>
<td>303</td>
<td>270</td>
</tr>
<tr>
<td></td>
<td>( d_{3/2}^{-1} )</td>
<td>248</td>
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<td>430</td>
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<tr>
<td>41</td>
<td>( f_{7/2}^{-1} )</td>
<td>154</td>
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<td></td>
<td>( 1p_{3/2}^{-1} )</td>
<td>228</td>
<td>265</td>
<td>340</td>
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<td></td>
<td>( 1p_{1/2}^{-1} )</td>
<td>249</td>
<td>303</td>
<td>330</td>
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\(^a\) from Blunden and Iqbal\(^{34}\)

**SUMMARY**

The quark mass difference leads to experimentally observed hadronic mass differences, \( \rho^0 - \omega \) mixing and a CSB nucleon-nucleon potential derivable from \( \rho^0 - \omega \) exchange. This potential accounts for the observed CSB in the nucleon-nucleon, three-body and heavy nuclear systems. About 85% of the NS anomaly (missing Coulomb energy) can be accounted for. The quark mass difference hypothesis is consistent with all of the observed CSB. An independent verification of the influence of the \( \rho - \omega \) mixing term in n-p elastic scattering would be of great interest.
REFERENCES

13. The QMD hypothesis does not require the presence of a robust $\rho-\omega$ mixing effect in nucleon-nucleon interactions. My current belief (G.A. Miller, Phys. Rev. C39, 1563 (1989) is that at low momentum transfers QCD leads to meson-nucleon dynamics. In that case, the $\rho-\omega$ term must influence $nn$ scattering.
16. See ref. 2 above. A recent caveat has been obtained by T. Goldman, K. Maltman, and G.J. Stephenson, LANL preprint, 1989. They find an effect of the electromagnetic vertex correction (EVC) to the quark-gluon interaction. If this (EVD) is large, the conventional lore would be revised. But at the moment the size of the EVC is not very well determined.
19. S. Godfrey and N. Isgur, Phys. Rev. D34, 899 (1986). Actually the text contains an oversimplification. Godfrey and Isgur point out that the mass difference data are very sensitive to the average quark mass, $m$. Thus both $m$ and $m_d - m_u$ are constrained by mass differences.
26. See ref 17.
   See also Y. Wu, S. Ishikawa, and T. Sasakiwa, 1989 Tohoku University preprint.
CONCLUSIONS AND RECOMMENDATIONS OF THE
SYMPOSIUM/WORKSHOP ON SPIN AND SYMMETRIES
TRIUMF, June 30–July 2, 1989

PARITY VIOLATION

1. How well does $A_z$ in $p-p$ scattering need to be measured at 15 MeV to be significant? — Accurately enough to distinguish between Kisslinger's and more conventional meson-exchange calculations but, in any case, not worse than $2 \times 10^{-8}$. This should be feasible, given currently available polarized ion source technology and state-of-the-art treatment of systematic errors. A very accurate low-energy measurement is desirable to check the result from PSI, because of the difficulty of the measurement, and/or to check the energy dependence of $A_z$.

There were questions about the energy dependence of $A_z$ in Kisslinger's calculations and whether there could be difficulties with unphysical short-range non-local effects (as found by Sauer, for example, where these introduced large charge-dependent effects, but phase shifts were left unaltered). To check the calculations, a measurement at low energy was suggested (the TRIUMF measurement of $A_z$ at 222 MeV determines a different combination of $h_p$ and $h_\omega$ than do those at low energy), but it was concluded there must be easier and more direct ways to check the calculations than a parity measurement. Kisslinger has not found sensitivity in few-body scattering or electromagnetic interactions.

2. Could $n-p$ measurements be made?

- The parity-violating neutron spin rotation in parahydrogen is to be measured at ILL. This is a very important experiment because, taken together with the 222 MeV TRIUMF measurement, it completes the minimum set of experiments required to determine the six weak meson-nucleon coupling constants.
- The value of $A_z$ for fast neutrons interacting with $^4$He and H is very interesting to measure, but is thought to be too difficult at present because of problems of intensity. But note that $^4$He is not an “elementary” particle. Also, for $^4$He the parity-allowed analyzing power, which is the prime cause of systematic error, is large. $^4$He would give the same constraints as $^{21}$Ne (the result which is difficult to interpret — see below).
- $A_z (n-p)$ could be extracted, in principle, from the $p-d$ and $p-p A_z$ measurements at 800 MeV made at LANL. The accuracy of these measurements is not great enough ($10^{-7}$, limited by statistics) for the result to be significant, however. A measurement with the new optically pumped source at LANL could improve the precision of the result significantly. The intensity of neutron beams is also increasing, so that a direct $n-p$ measurement may become feasible in a few years.

3. The smallness of $f_r$ is not understood — It was suggested that a lattice gauge calculation be made to explain the small value. Is there an analogy between the $\Delta I = 1/2$ rule and the apparent suppression of isospin 1 weak coupling constants compared with isospin 0? Input from particle theorists may be needed.
4. The 5.1 GeV measurement of $A_z$ used a water target. Is the extraction of $A_z(p-p)$ a problem? — It's thought not to be. The measurement needs to be repeated, if only to confirm the result because of the great difficulty of the measurement. There has been no concrete criticism of the measurement itself. The result must be taken at face value.

5. It was noted that there are few new initiatives to measure parity violation. It was thought that experimentalists were being too pessimistic in viewing the time and effort required, but that the effort is certainly very worth while. Sources of systematic error in the measurements are now well explored and detailed in a number of publications for future experimentalists to read. The TRIUMF experiment is the ONLY NEW MEASUREMENT of $A_z$ and will be very important, provided the stated accuracy is achieved ($\pm 2 \times 10^{-8}$).

6. Can good calculations be made of $A_z(p-d)$ below breakup threshold? — Yes, although these are very lengthy calculations. There is a well-established three-body formalism and non-separable interactions can now be used. Calculations of $A_z(p-^4\text{He})$ are not on the same footing although it should be possible to calculate phases below breakup threshold.

7. $np \rightarrow d\gamma$ is interesting as it gives an extra constraint on weak couplings. There are ideas to improve the $^{18}\text{F}$ result by a factor of three. Also, a measurement is possible on $^{18}\text{O}$ — many E1 transitions have been measured over the last couple of years, so the nuclear structure is thought to be well in hand. $^{18}\text{O}$ sets the same constraints as $^{19}\text{F}$. It is crucial for shell model calculations to fit the low-lying E1 transitions before reliable parity-violating effects can be calculated with confidence. There is general agreement that for $^{21}\text{Ne}$ the nuclear structure calculations are sufficiently unreliable at present that the result based on the measurement of circular polarization should not be included in the current data set. Improved calculations can be expected when the $4\hbar\omega$ excitations are included.

8. Electron scattering — $^{12}\text{C}$ is a choice candidate for a parity experiment to get rid of nuclear effects. Experiments at CEBAF on H and He are complementary to measurements of $\sin^2 \theta_w$ at LEP or SLC. The question was raised whether uncertainties on $G^n_E$ would ultimately limit our ability to put limits on extensions to the standard model (e.g. extra gauge bosons).

**CP and T VIOLATION**

Although CP symmetry violation has been known for many years, there exists only one experimental number, $\epsilon$, from kaon decay with a tentative second $\epsilon'$, also from the kaon system, expected shortly (there appears to be at present a discrepancy between the CERN and FNAL results). Any nonzero number from any other system would be VERY IMPORTANT in helping to determine the origin of CP violation and constraining the fertile imagination of theorists. Limits on the magnitude of CP violation are less useful, especially in nuclear systems where there are uncertainties in interpretation. There is no direct evidence of time-reversal violation but it is strongly implied by CPT invariance together with CP violation. It is difficult to construct models that have $T$-invariance violation without parity conservation violation, and if such models are constructed it is expected that manifestations of $T$-invariance violation will be suppressed in nuclei, in part because the interaction is short ranged.
1. Types of Tests

The tests of CP or T nonconservation fall into three main classes — hadronic, electromagnetic, and weak. The most accurate hadronic tests have used detailed balance; a recent new class of tests consists of neutron transmission through polarized targets. For low-energy neutrons, nuclear systems offer the possibility of enhancements by factors of $10^3$ to $10^5$ when carried out at or near $p$-wave resonances. Thus these systems offer a good place to look. The main drawback is the uncertainties due to nuclear structure. This will make limits obtained by this method more difficult to interpret. However, it is expected that the ratio of matrix elements calculated for parity violation and for time reversal violation should be more reliable. Furthermore, parity violation was first seen in nuclear systems where there was a nuclear enhancement and was then studied in more theoretically accessible systems. Something similar might easily happen for time reversal violation as well.

The first type of electromagnetic tests makes use of gamma-ray correlations between competing multipoles. The most accurate ones have been interpreted to set limits on a $T$-odd $P$-conserving interaction strength of $< 10^{-7}$ and on a $T$-odd $P$-nonconserving strength of $< 10^{-5}$. A second type of test is the search for the static electric dipole moments of the neutron and atomic systems. For the neutron dipole moment the limit is currently $11 \times 10^{-28}$ e·cm at 95% C.L. which provides a limit $\theta_{\text{QCD}}$ of $5 \times 10^{-10}$. The lack of an electric dipole moment of $^{199}$Hg ($d_e \leq 2 \times 10^{26}$ e·cm) and of the neutron both set limits on a $T$-odd $P$-nonconserving interaction strength of $< 10^{-11}$!

The weak interaction tests fall into two classes, those involving $K_0\overline{K}_0$ (or $B\overline{B}$) system and those involving other systems where the polarization of the particles in the initial and/or final state are measured. CP violation has so far only been seen in kaon decay. The standard model predicts similar effects in the $B$ system. However, their observation will require a $B$ factory and hence one does not expect such experiments to be done for at least eight years. The most accurate tests to date have been carried out in $\Lambda^\pi$ decay of polarized neutrons and $^{19}$Ne. Other tests, though difficult, may be done in the next few years, e.g. $K_L^0 \rightarrow \pi^0 e^+ e^-$, $p\overline{p} \rightarrow \Lambda\overline{\Lambda}$, $p\overline{p} \rightarrow \Xi^- \Xi^+ \rightarrow \Lambda \pi^0 \pi^0$ and $\eta \rightarrow \mu^+ \mu^-$. The first of these, namely $K_L^0 \rightarrow \pi^0 e^+ e^-$, looks for asymmetry in the Dalitz plot while the remaining reactions look at polarization asymmetries. The reaction, $\eta \rightarrow \mu^+ \mu^-$, is especially interesting in view of the large $\eta$ yield observed at the SATURNE accelerator in the reaction $pd \rightarrow \overline{3}H\eta$ ($\eta$-factory?).

Finally, it was mentioned that one should perhaps try to look elsewhere for $T$-reversal violation such as in macroscopic systems. For example, gravity may exhibit $T$-reversal violation.

2. Conclusion

Any measurement of nonzero $T$ or CP violation will be very important and will have a profound influence on our understanding of its origin. It should be kept in mind that in the standard model the parameters giving rise to $T$ or CP violation are not fixed by the model (i.e. one has a model and not a theory).

**CHARGE SYMMETRY BREAKING**

The classic test of charge symmetry (CS) involves the comparison of the strong $pp$ and $nn$ scattering lengths ($a_{pp}$ and $a_{nn}$). Experimentally one can obtain the Coulomb-modified value of $a_{pp}$ relatively easily but unfolding the effect of the Coulomb potential is inherently model dependent. The commonly accepted value is 17.2 fm with a theoretical error of perhaps
±0.5 fm. On the question of this error the discussion was lively with some feeling 0.5 fm was an upper limit and others that it was optimistic in view of the possible nonlocality of the short-distance \( N-N \) force. For \( a_{nn} \), the experimental problem is harder with the best values of \( 18.4 \pm 0.4 \) fm coming from the \( \pi^{-}d \rightarrow nn\gamma \) reaction. There was some concern that the values obtained from extensive studies of \( n-d \) break-up are systematically lower by some 2 fm. There was a strong feeling that a modern state-of-the-art calculation, including all \( N-N \) partial waves, might resolve the discrepancy, and this was felt to be a high priority. Conventional meson-exchange models of \( |a_{nn}| - |a_{pp}| \) typically give around 1.3 fm which is in agreement with the experimental difference. However, the error on the experimental value is rather large \( (|a_{nn}| - |a_{pp}|)_{\text{Expt}} = 1.2 \pm (0.5^2 + 0.4^2)^{1/2} \), and it is not clear how to improve it since it is dominated by theoretical uncertainties.

Another famous low-energy test is the \( ^3\text{H}-^3\text{He} \) mass difference. The direct Coulomb correction is under control with modern three-body estimates in close agreement with model-independent estimates (as they should be). Conventional meson-exchange corrections reproduce the remaining \( 80\pm24 \) keV discrepancy coming mainly from \( \rho-\omega \) mixing. The need for more theoretical work, notably on electromagnetic coupling to virtual \( \Delta \)'s and three-body effects, came through in the discussion. Thus although the situation looks good it is not yet a closed problem.

Turning to higher-energy reactions: The forward-backward asymmetry in the \( np \rightarrow d\pi^0 \) reaction has not yet been studied in an experiment dedicated to that aim. However, two experiments are planned for the near future, one at LAMPF and one at TRIUMF. One mechanism for charge symmetry breaking (CSB) which is of special interest in this case is \( \eta-\pi^0 \) mixing. (In this regard it would be particularly useful to study the complementary reaction \( np \rightarrow d\eta \).) Since the effect of a major part of the \( \eta-\pi^0 \) mixing contribution is to yield an effective coupling constant \( f_{\pi^0_{nn}} \) different from \( f_{\pi^0_{pp}} \), one needs as much help as possible to separate this from the difference caused by the \( u-d \) mass difference in the nucleon. While a number of theoretical mechanisms have not yet been estimated, like the nucleon mass difference in \( \rho \)-exchange and electromagnetic effects involving the \( \Delta \), the theoretical analysis is nevertheless sophisticated. There is an urgent need for data for the \( np \rightarrow d\pi^0 \) reaction, preferably involving as broad a range of angles as possible and at least two energies.

There was some discussion of tests of charge symmetry and charge independence in heavier systems, such as \( d-^3\text{He} \) and \( d-^3\text{H} \). The general tone was cautious, with worries over the interpretation in view of Coulomb corrections – both direct and implicit in wave function differences.

Largely because of the beautiful new data from IUCF the longest discussion involved the test of class-IV CSB in \( n-p \) elastic scattering. With the only earlier data being a \( 2\sigma \) observation at TRIUMF (at 477 MeV), this \( 4\sigma \) measurement doubles the world data. This accurate result is not before its time, because the theoretical situation has matured over the last 3 or 4 years. For example, it is clear that the earlier TRIUMF measurement is insensitive to \( \rho-\omega \) mixing but rather is dominated by the effect of the nucleon mass difference \( (\delta m_N) \) on pion exchange. This is, of course, unambiguous except for the range of the \( \pi NN \) form factor and a little model dependence in the short-distance \( N-N \) distorted waves. However, a high accuracy measurement at either 477 MeV or 350 MeV (as planned at TRIUMF) would be very useful in isolating the \( \rho \)-meson and pion contributions to \( \delta m_N \) and hence indirectly in extracting the effect of \( \rho-\omega \) mixing from data at lower energies.

It has become clear that the DWBA is the best theoretical method to study \( np \) CSB – although in the region where \( \Delta \)-production is important the effect of channel coupling
will need to be treated explicitly. Since the CSB potential is usually calculated in a boson-exchange model, the momentum space DWBA calculation involving the Bonn potential is clearly the most consistent. It also happens to give the best agreement with the new IUCF data at 183 MeV. At this energy the Bonn potential (with $g^2_{\omega NN}/4\pi = 20$) implies that $\rho-\omega$ mixing gives about half of the CSB effect ($A_n - A - p$). While the data analysis is not final and the result may move a little before publication, the agreement does seem significant.

Further discussions dealt with the question of whether or not a $\rho-\omega$ contribution is present is significant. Suppose one considers conventional, time-ordered, valence quark and gluon models (supplemented by pion exchange) of the $N-N$ system, such as resonating group or coupled cluster calculations. Then the quark mass difference, which in a $q-\bar{q}$ system generates $\rho-\omega$ mixing, cannot give a class-IV force. Only by introducing $q-\bar{q}$ pairs explicitly in the time-ordered calculations can such a quark model hope to give an effect analogous to $\rho-\omega$ mixing. Since these models have been relatively successful in describing the short-distance $N-N$ interaction (conventionally described by $\omega$-exchange) without such admixtures, unambiguous observation of $\rho-\omega$ mixing would be important.

The conclusion of this discussion was that there is a clear need for at least TWO MEASUREMENTS of $\Delta A$ at a level of precision at least $\pm 5 \times 10^{-4}$. One of these should be in a region sensitive to $\rho-\omega$ mixing and the other in a region insensitive to it (to pin down the other contributions). A third measurement at high energies where $\Delta$ production is important would open up new physics possibilities and stimulate more theoretical work.
PROGRAM
SYMPOSIUM/WORKSHOP ON SPIN AND SYMMETRIES
TRIUMF, June 30 – July 2, 1989

Friday, June 30, 1989

08:00    Registration at TRIUMF
TRIUMF Auditorium, Chairperson: R.E. Mischke
Scientific Secretaries: P. Blunden, J. Birchall

08:25 – 08:30    E.W. Vogt, TRIUMF – Welcome!

08:30 – 09:15    “A measurement of the Flavor Conserving Weak Hadronic Interaction
in \( \bar{p}p \) Scattering at 222 MeV”
S.A. Page, University of Manitoba

09:15 – 10:00    “Parity Violation in the Nucleon-Nucleon Interaction”
W.C. Haxton, University of Washington

10:00 – 10:30    Coffee Break

10:30 – 11:00    “Status of Parity Violation in \( \bar{p}N \) Scattering at Higher Energies”
M. Simonius, ETH

11:00 – 11:30    “A Measurement of Parity Violation in \( \bar{p}p \) Scattering at 14.5 MeV”
P.D. Eversheim, Universit"at Bonn

11:30 – 12:00    “A Measurement of Parity Violation in Electron Scattering from Carbon”
G.W. Dodson, MIT

12:00 – 12:45    “Parity Violation Measurements at CEBAF”
R. Carlini, CEBAF

12:45 – 14:00    Lunch

14:00 – 14:30    “Pion Exchange Contribution to the Weak Longitudinal Analyzing Power
in \( \bar{p}p \) Scattering”
R.R. Silbar, LANL

14:30 – 15:00    “Relativistic and Inelastic Effects for the Proton-Proton Parity
Violating Interaction”
D.E. Driscoll, University of Washington

15:00 – 15:30    “Weak \( \bar{p}p \) Analyzing Power and Short Range Interactions”
L.S. Kisslinger, Carnegie-Mellon University

15:30 – 16:00    Coffee Break

16:00 – 17:30    Discussion Session
W. Haeberli, University of Wisconsin

19:00    Salmon Barbecue in the TRIUMF Courtyard

Saturday, July 1, 1989

TRIUMF Auditorium, Chairperson: P.C. Gugelot
Scientific Secretaries: B. Jennings, G. Roy

08:30 – 09:15    “CP Violation and Decay Asymmetries”
J. Ng, TRIUMF

09:15 – 10:00    “Hyperon-Antihyperon Decay Asymmetries and CP Violation”
D. Hertzog, University of Illinois

10:00 – 10:30    Coffee Break

10:30 – 11:15    “Nuclear Enhancement of Time Reversal Non-Invariance”
E.D. Davis, Max Planck Institut, Heidelberg
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| 11:15 - 12:00 | "Tests of Parity and Time Reversal Invariance in Neutron Scattering"  
C.R. Gould, North Carolina State University |
| 12:00 - 12:45 | "Tests of Symmetries Involving Spin at KEK"  
A. Masaike, Kyoto University |
| 12:45 - 14:00 | Lunch                                                                                     |
| 14:00 - 14:30 | TRIUMF Auditorium, Chairperson: J.D. Bowman  
Scientific Secretaries: G. Roy, B. Jennings |
| 14:30 - 15:00 | "On the Electric Dipole Moment of the Neutron"  
B.R. Heckel, University of Washington |
| 15:00 - 15:30 | "The Neutron Electric Dipole Moment in Various Models of CP Violation"  
B.H.J. McKellar, University of Melbourne |
| 15:30 - 16:00 | Coffee Break                                                                             |
| 16:00 - 17:30 | Discussion Session  
E.M. Henley, University of Washington |

**Sunday, July 2, 1989**

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| 08:30 - 09:15 | "Charge Symmetry breaking in n-p Elastic Scattering at 183 MeV"  
W. Jacobs, IUCF |
| 09:15 - 10:00 | "Charge Symmetry Breaking in n-p Scattering"  
L.G. Greeniaus, TRIUMF |
| 10:00 - 10:30 | Coffee Break                                                                             |
| 10:30 - 11:00 | "Calculations of Charge Symmetry breaking n-p Elastic Scattering"  
A.G. Williams, University of Washington |
| 11:00 - 11:45 | "Charge Symmetry Breaking in the n-p System"  
J.A. Niskanen, TRIUMF/University of Helsinki |
| 11:45 - 12:15 | "Forward-Backward Asymmetry in the Reaction n-p → d-π^0"  
J.C. Hiebert, Texas A&M University |
| 12:15 - 12:45 | "Charge Symmetry Breaking in Few-Nucleon Systems"  
W. Grüebl, ETH |
| 12:45 - 14:00 | Lunch                                                                                     |

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| 14:00 - 14:30 | TRIUMF Auditorium, Chairperson: W.T.H. van Oers  
Scientific Secretaries: N.E. Davison, S. Bass |
| 14:30 - 15:00 | "Charge Symmetry of the Nuclear Interaction and the N-N Scattering Parameters"  
G.F. de Téramond, Universidad de Costa Rica |
| 15:00 - 15:45 | "Charge Symmetry Breaking in the ^3H-^3He System"  
T. Sasakawa, Tohoku University |
| 15:45 - 16:15 | Coffee Break                                                                             |
| 16:15 - 17:30 | Discussion Session  
A.W. Thomas, University of Adelaide |
| 17:30 | End of the Symposium/Workshop |
LIST OF PARTICIPANTS

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