ENERGY BALANCE AND TEMPERATURE IN A CO₂ LASER PRODUCED PLASMA

by

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Abstract

The absorption of intense $\leq 10^{14}$ W/cm² CO₂ laser radiation by an underdense plasma is investigated experimentally. The plasma is produced by the CO₂ laser beam focussed onto a stabilized laminar gas jet emanating from a Laval nozzle. The electron temperature of the laser produced plasma is measured by soft x-ray diagnostics which yield a 300 eV thermal and a 2000 eV suprathermal temperature. Time resolved studies are made of the radial expansion of the plasma to determine the absorbed energy and temperature using a modified blast wave analysis that involves the identification of the Chapman-Jouguet detonation point in the expansion. The absorbed energy is also determined from Ulbricht sphere photometry. Information on the plasma dimensions and electron density is obtained by time resolved ruby laser interferometry.

The electron temperature and absorbed energy measurements are used to corroborate one another and to provide a self-consistent picture of the laser plasma coupling. Although vacuum laser intensities are sufficiently high where saturation of collisional (inverse bremsstrahlung), absorption is expected to occur, the experimental evidence indicates that linear inverse bremsstrahlung accounts for the observed thermal electron temperature and measured level of absorption. The energy balance shows that there is a sufficient level of absorption to accommodate the measured thermal electron temperature as well as < 15% of the electrons at a 2 keV suprathermal temperature.

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Chapter I: Introduction

Shortly after the invention of the ruby laser in 1960 by Maiman, experiments began to produce plasmas by focussing laser radiation onto materials. An extensive review of these numerous studies can be found in Hughes¹. These plasmas were comparatively much hotter and denser than those typically found in electrical discharge devices. The possibility of achieving sufficiently high temperatures and densities to obtain controlled thermonuclear fusion has been the impetus for much of the research in plasma physics since the 1950's. Accordingly, the extrapolation of laser produced plasmas to conditions where thermonuclear fusion can occur remains to this day a topic of intensive research efforts throughout the world.

The main idea behind the laser fusion program is the isentropic implosion of a deuterium-tritium mixture. The increase in density produced by such an implosion can reduce the energy required to produce a given thermonuclear output energy by a factor of as much as 10^4 below that needed for heating an uncompressed plasma².

Most of the early implosion experiments in the 1970's used simple thin shell glass microballoons filled with D-T gas at high pressure which acted like "exploding pushers". These experiments were characterized by short plasma scale lengths and short (<1 ns) laser pulses of several kJ energy. Although these experiments yielded large amounts of neutrons and high plasma temperatures, thermonuclear ignition cannot be achieved with any feasible laser system based on an exploding pusher scheme. Since

1980³ the emphasis has turned towards ablative compression studies in which the D-T gas is adiabatically compressed. These experiments are done with targets that have a comparatively thick plastic outer shell which is ablated by a multಿನanosecond laser pulse resulting in an extensive tenuous underdense corona. (The electron density of the corona plasma is said to be underdense when the laser light frequency is greater than the plasma frequency and can propagate through the plasma. Laser light is reflected at the critical density (n_{cr}) and plasma densities greater than this are said to be overdense.) The interaction between the laser and this underdense coronal plasma is important in these ablative pusher experiments where phenomena such as stimulated Brillouin (SBS) and Raman scattering (SRS) can have serious detrimental consequences on the efficiency of laser energy coupling to the target. In SBS and SRS, the incident laser light interacts with density perturbations in the plasma resulting in a scattered laser beam and subsequent enhancement of these perturbations. These phenomena can cause a substantial reflection of the incident laser energy.

The effective absorption of laser light is perhaps the most important issue in all laser fusion schemes. The objective is to heat the body of the electron distribution avoiding the generation of suprathermal electrons which can penetrate the cold dense plasma above the critical density layer causing unwanted preheating of the target. Several absorption mechanisms are known to occur in the underdense corona. The most important mechanism is collisional inverse bremsstrahlung absorption which

arises from the interaction of the electrons in the Coulomb field of the ions. Electrons oscillate under the influence of the laser electric field and subsequently transfer their laser field induced oscillatory motion to thermal kinetic energy. Although the validity of the theory for inverse bremsstrahlung absorption been tested and verified for low laser intensities of $\leq 10^{10}$ W/cm², at higher intensities where the oscillatory velocity of the electrons v_{os} , becomes comparable to the thermal velocity v_{th} , this theory is expected to fail as electron-ion collisions become increasingly less effective for energy transfer of the oscillatory motion to thermal energy. To date, most laser plasma experiments are concerned with the absorption of laser radiation by solid targets. In these cases, the presence of a critical density layer complicates the overall picture so that conclusions concerning the role of inverse bremsstrahlung absorption are somewhat insubstantial.

This thesis experimentally examines the absorption of CO₂ laser light by a laser produced underdense plasma. In order to determine the absorbed laser energy, the plasma is diagnosed by soft x-ray measurements, Ulbricht sphere photometry, interferometry and time resolved photography. The main question is how the absorbed laser energy is correlated with the measured electron temperature. This enables a comparison between the expected absorption due to inverse bremsstrahlung and measured values for an underdense plasma. The energy balance study allows the investigation of the distribution of the absorbed energy among the various processes which take place in the laser-plasma interaction, e.g. hot electron production, radiation losses,

etc.

The organization of this thesis is as follows. The theory of absorption of laser radiation by a plasma is discussed semi-quantitatively in the following chapter to provide some physical insight to the processes as well as to introduce some of the complications and effects that are expected to arise for the experimental conditions encountered. A simple electron heating model is presented along with estimates of the expected electron temperature. These estimates are then supplemented by the results of several computer simulations that were performed using available hydrodynamic computer codes.

Since this thesis is mainly an experimental work describing the interaction of a CO₂ laser beam and the plasma it produces, brief descriptions of the apparatus used follow in Chapter III. To ensure production of a strictly underdense plasma, the CO₂ laser beam is directed and focussed upon a variable density gas jet target. The device for producing the laminar gas target and its characterization for use as a target is described in this chapter. In addition to the hardware of the laser and its target, the diagnostic instrumentation developed and implemented in the course of the experiments is also presented here.

The experimental results for the electron temperature and energy absorption are contained in Chapter IV. Error estimates are considered and the results are compared with the data obtained from other laboratories for comparable experimental conditions.

A unifying self-consistent picture of the absorption process of the laser-plasma interaction is developed in Chapter

V. The results of various calculations are tabulated and the energy inventory is presented. The nature of the suprathermal temperature measured by soft x-ray diagnostics receives consideration and speculation regarding its possible origin.

The results are summarized and concluding statements concerning the nature of the absorption of laser radiation for CO₂ laser intensities $\leq 10^{14}$ W/cm² in underdense plasma are made in Chapter VI. The most probable mechanism for the generation of the suprathermal electrons is proffered. Since no investigation can ever produce all the answers, a few suggestions for further work are included as well.

The thesis and the new significant results it contains can be summarized as follows. The thesis describes the absorption and energy balance of a unique laser-underdense plasma interaction based on x-ray diagnostics, interferometry, streak photography and Ulbricht sphere measurements. A new method of analysis of the streak records of the radial plasma expansion is introduced and implemented to provide an independent estimate of the plasma temperature, pressure and absorbed laser energy. The Ulbricht sphere and x-ray diagnostics measurements, along with interferometric electron density information, yield the absorbed energy and electron temperature as well, enabling a comparison of the results for inverse bremsstrahlung absorption and to make an energy inventory for the interaction. All of the results can be explained on the basis of inverse bremsstrahlung absorption, although saturation of this absorption mechanism is predicted by theory for the experimental conditions encountered.

Chapter II: Absorption of Laser Radiation

The study of the energy balance requires a detailed account of the energy absorption mechanisms and of the energy loss processes. Several different absorption processes may in whole or in part, account for the absorption observed in laser-plasma experiments. It is therefore appropriate to presently review some of these various absorption mechanisms.

The purpose of this chapter is to provide some insight into the relevant physics by which a laser beam is absorbed by a plasma. The discussion will be limited to semi-quantitative reviews of the physical processes of absorption mechanisms since a thorough theoretical treatment of this subject is beyond the scope of this thesis. The review of the physical processes involved in the collisional absorption of laser light serves as a prelude to a discussion of some of the anomalous mechanisms that can occur. In addition, some numerical estimates of the expected experimental results will be presented along with the results of some computer simulations that were performed.

II-1 Inverse bremsstrahlung absorption

Inverse bremsstrahlung absorption is considered to be the most important mechanism for the absorption of laser light by a plasma. It is the process by which the electron population increases its thermal energy through the absorption of photons in the Coulomb field of the ions. Other absorption processes to

be discussed later create an energetic component of hot electrons which usually constitutes a minority of the electron population.

The inverse bremsstrahlung process arises from the oscillation of the electrons in the electric field of the incident laser light. This directed energy of electron motion is randomized and hence converted into heat energy by electron collisions with ions. The equation of motion of the electron under these circumstances is approximately

$$m(d\vec{v}/dt) + m\nu_{ei} \vec{v} = -e\vec{E},$$

here the collisional damping term $m\nu_{ei} \vec{v}$ contains the ion collision frequency ν_{ei} . There is an induced current $\vec{J}$, due to the electron motion along the direction of the electric field vector

$$\vec{J} = n e \vec{v},$$

and $\vec{J}$ is in turn governed by Maxwell's equations. The simultaneous solution of Maxwell's equations along with the above equations, with the assumption of plane wave oscillatory dependences of the form

$$\begin{pmatrix} \vec{E}(\vec{r}, t) \\ \vec{B}(\vec{r}, t) \\ \vec{v}(\vec{r}, t) \end{pmatrix} = \begin{pmatrix} \vec{E} \\ \vec{B} \\ \vec{v} \end{pmatrix} \cdot \exp(i(\vec{k} \cdot \vec{r} - \omega t)) ,$$

produces the dispersion relation for transverse ($\vec{k} \cdot \vec{E} = 0$) electromagnetic waves:

$$\begin{aligned}
 k^2 c^2 / \omega_L^2 &= 1 - \omega_p^2 / \omega_L (\omega_L + i \nu_{ei}), \\
 &= 1 - \omega_p^2 / \omega_L^2 + (i \nu_{ei} / \omega_L) (\omega_p^2 / \omega_L^2),
 \end{aligned}$$

where the expansion is done assuming $\nu_{ei} / \omega_L \ll 1$, ω_p is the plasma frequency $\omega_p = 4\pi n_e e^2 / m_e$. Since the energy of a wave is proportional to its amplitude squared, the spatial damping rate of wave energy, k_{ib} is then twice the imaginary part of k as obtained from the above dispersion relation:

$$k_{ib} \approx \frac{\nu_{ei} \omega_p^2 / c \omega_L^2}{(1 - \omega_p^2 / \omega_L^2)^{1/2}}.$$

The scaling of the electron-ion collision frequency ν_{ei} is obtained by the following reasoning. The collision frequency is $\nu_{ei} = n_i \sigma v_{th}$ for an electron of thermal velocity v_{th} colliding with background ions of number density n_i . The collision cross section is $\sigma = \pi b^2$ where the distance of closest approach b is given by the balance of potential and kinetic energy:

$$Ze^2 / b = 0.5 m_e v_{th}^2, \quad \text{where } v_{th} = (kT_e / m_e)^{1/2}$$

The collision frequency scaling is obtained by substituting the above expression for b and is

$$\nu_{ei} \propto n_i (2Ze^2 / m_e v_{th}^2) v_{th} \propto n_e Z / T_e^{1.5},$$

so that the inverse bremsstrahlung absorption coefficient is

$$k_{ib} \propto \left\{ \frac{Z n_e^2}{T_e^{3/2}} \right\} \left(1 - \frac{n_e}{n_{cr}} \right)^{-1/2},$$

where $n_{cr} = 1 \times 10^{19} \text{ cm}^{-3}$ for CO_2 laser light wavelength $\lambda = 10.6 \text{ } \mu\text{m}$. From the above relation, it is apparent that inverse bremsstrahlung absorption is greatest for low temperature, high density, high Z plasmas.

A much more rigorous calculation of the inverse bremsstrahlung coefficient involves the use of kinetic theory which takes into account the actual distribution function of the electrons and actual position of the ions under the influence of the laser electric field⁴. In addition, quantum mechanical corrections to the classical results are incorporated into a Guant factor. Numerically, to sufficient accuracy, the absorption coefficient obtained is the Johnston and Dawson value⁵:

$$k_{ib} \propto 9.74 \times 10^{-35} \frac{Z n_e^2}{T_e^{3/2} (1 - n_e/n_{cr})^{-1/2}} \text{ cm}^{-1} \quad (T_e \rightarrow \text{eV}, n_e \rightarrow \text{cm}^{-3})$$

Billman and Stallcop⁶ have considered other low temperature mechanisms such as electron-neutral atom collisional absorption, photoionization, etc. Their absorption coefficient is higher than that of the Johnston and Dawson value by as much as a factor of 5 for electron temperatures of the order of a few eV. However, for $T_e > 20 \text{ eV}$ the difference between the two absorption coefficients is negligible.

II-2 Saturation of Collisional Absorption

At very high laser intensities the inverse bremsstrahlung coefficient as calculated by Johnston and Dawson is expected to saturate with increasing intensity so that the absorption becomes inefficient. This is because the calculation of the absorption coefficient assumes the electron distribution to remain Maxwellian whereas actually, the electron-ion Coulomb collision frequency is modified by the electron's coherent motion in the laser electric field. Schlessinger and Wright⁷ review and compare the different approaches to quantify the effect of intense fields on the inverse bremsstrahlung absorption rate. The rigorous theoretical treatments produce a correction term to the absorption rate that is in agreement with the heuristic treatment of Faehl and Roderick⁸ to within 10%. Instead of having a typical mean square velocity v_{th}^2 , the electron now has an effective average velocity

$$\langle v^2 \rangle = v_{th}^2 + \langle v_{os}^2 \rangle$$

where v_{os} is the oscillatory velocity of the electron

$$v_{os} = eE(t)/m\omega_0, \quad E(t) = E_0 \cos(\omega_0 t + \Delta).$$

The absorption coefficient is proportional to the electron ion collision frequency which in turn has the velocity dependence

$$\nu_{ei} \propto \langle v^2 \rangle^{-1.5}$$

This suggests that the inverse bremsstrahlung coefficient must be reduced by the factor:

$$k_{ib} \rightarrow k_{ib} / (1 + \langle v_{os}^2 \rangle / 3v_{th}^2)^{-3/2}.$$

This simple theory agrees with other calculations to within 10% up to at least $v_{os}/v_{th} = 10$. For the experimental conditions in this thesis $v_{os}/v_{th} \leq 3$.

A second non-linear modification to inverse bremsstrahlung was discussed by Langdon⁹. It arises from the fact that the electron distribution can be far from Maxwellian when the laser intensity is high. For a Maxwellian distribution, the heating rate for an electron in the electric field of the laser is

$$0.5 m_e v_{os}^2 \nu_{ei}.$$

However, the rate at which this electron can share this energy with other electrons to form a Maxwellian distribution is proportional to

$$\nu_{ee} v_{th}^2,$$

where ν_{ee} is the electron-electron collision frequency $\nu_{ee} = \nu_{ei}/Z$. Therefore, the heating rate can exceed the rate at which electrons can form a Maxwellian distribution, establishing the condition for the formation of a non-Maxwellian distribution:

$$Z v_{os}^2 / v_{th}^2 > 1.$$

It is the low velocity electrons, $v_e \ll v_{th}$, that are the most strongly perturbed by collisions. When the above condition is satisfied, the number of these low energy electrons is reduced. This has the effect that the amount of collisional absorption decreases below the value calculated for a Maxwellian distribution.

The reduction as calculated by Langdon takes the form of

$$k'_{ib} \rightarrow k_{ib} \exp(-v^2/\omega_{th}^2),$$

where v_{ω} is the electron velocity for which $v_{ei} = \omega_L$, since the inverse bremsstrahlung absorption rate is proportional to the number of electrons in the distribution which have $v_{ei} = \omega_L$. Generally, this reduction in the absorption rate is greater than that calculated by the heuristic model and is the same at $v_{os}/v_{th} \approx 7$.

II-3 Resonance Absorption

Another important possible mechanism of absorption of laser light that is considered in this thesis is that of resonance absorption. This is also known as collisionless absorption since the collision frequency does not enter into the absorption rate for this mechanism. For the case of light incident on a plasma density gradient at angle to the gradient direction Snell's law and the dispersion relation

$$\omega_L^2 = \omega_p^2 + c^2 k_L^2 ,$$

can be combined to show that the turning point for the light wave occurs at the density $n_e = n_{cr} \cos^2 \theta$. For p polarized light, the electric field $\vec{E}$, is in the plane formed by $\vec{k}$ and ∇n . At the turning point, the electric field is in the direction ∇n . This component of the electric field can tunnel into the critical density layer becoming very large and resonantly excite electron plasma waves. The polarization dependence of resonance absorption is apparent from Maxwell's equations for electromagnetic waves in a plasma:

$$\nabla \cdot (\epsilon \vec{E}) = 0 = (\vec{E} \cdot \nabla \epsilon + \epsilon \nabla \cdot \vec{E}) ,$$

$$\text{where } \epsilon \simeq 1 - \omega_p^2 / \omega^2 = 1 - n_e / n_{cr} .$$

The electron density perturbation δn_e of the electron plasma wave follows:

$$\begin{aligned} \nabla \cdot \vec{E} &= -4\pi e \delta n_e \\ &= -\vec{E} \cdot \nabla \epsilon / \epsilon \propto \vec{E} \cdot \nabla n_e / \epsilon . \end{aligned}$$

Therefore, if $\vec{E} \cdot \nabla n = 0$, then $\delta n_e = 0$ and there are no driven electron waves. But if there is a component of the electric field along the density gradient, $\vec{E} \cdot \nabla n \neq 0$, then $\delta n_e \rightarrow \infty$ since $\epsilon \rightarrow 0$ as the electron density approaches critical density n_{cr} . For a linear density gradient the power absorbed from the laser field may be shown to be¹⁰

$$P_{abs} = \omega_0 L E_d^2 / 8\pi ,$$

where L is the density gradient scale length and E_d is the component of the electric field that drives the waves and has been calculated by Kruer¹¹. The fraction of incident light that can be absorbed by the resonance process is limited to 50%, and values of this absorption as high as this value have been observed for CO₂ laser-solid target experiments.

A second important consequence of resonance absorption is the generation of hot-energetic electrons. The laser energy coupled into the plasma oscillations can dissipate through Landau damping appearing then as kinetic energy of heated electrons. Two dimensional computer simulations have shown that the resonantly driven wave field can grow to sufficient intensity that the electrons are trapped and accelerated in one wave period. The hot electrons are found to be characterized by a suprathermal Maxwellian distribution whose temperature is approximately¹²

$$T_H = 14(I\lambda^2)^{0.33} T_c^{0.33} \quad (\text{keV}) ,$$

where I is the laser intensity in units of 10^{16} W/cm², λ is the laser wavelength in μm and T_c is the thermal temperature. For the experimental conditions encountered in this thesis, the hot electron temperature that may be expected is 6 - 10 keV.

II-4 Ion-Acoustic Turbulence

Another likely mechanism for absorption is ion acoustic turbulence, a phenomenon which has been invoked to account for the high levels of absorption in experiments which cannot be explained in terms of inverse bremsstrahlung or resonance absorption. The idea is that if by some means a continuous noise spectrum of ion acoustic waves is excited, then this resulting turbulence can scatter electrons effectively, reducing the heat flow, thus creating a higher temperature than can be expected from collisional (inverse bremsstrahlung) absorption alone. Also the electron ion collision frequency is increased by the level of turbulence and hence there is a high amount of absorption of laser radiation.

There are several models of ion wave turbulence (see e.g. Monchicourt¹³). One common cause is thought to be the following. The energy flux $\vec{q} = \kappa \nabla T$, where κ is the thermal conductivity, produces a distortion of the electron distribution function $f(\vec{r}, \vec{v})$. Its maximum is shifted by

$$\vec{v}'_e = -e\vec{E}/m_e \nu_{ei} = \alpha/\sigma (e\nabla T/m_e \nu_{ei}) ,$$

$$\text{since } \vec{j} = \sigma \vec{E} + \alpha \nabla T = 0 ,$$

where σ is the electrical conductivity and α is the Seebeck coefficient. There is a return current of cold electrons $n_e \vec{v}'_e$ which compensates the hot electrons of $f(\vec{r}, \vec{v})$ which carry the heat flux. It was shown by Gray and Kilkenny¹⁴ is that the heat flux is carried by electrons with $v_e = 2.3$ to 3 times v_{th} .

The ion acoustic instability appears when the electron drift velocity exceeds the ion acoustic speed

$$v_e \geq (ZT_e / m_i)^{0.5}.$$

When the above condition is fulfilled, a continuous turbulent spectrum of ion acoustic modes is generated whose effect is an increase in the momentum transfer between ions and electrons since the electrons can now scatter from ion charge bunches rather than individual ions. In the presence of ion turbulence the energy damping rate is¹⁵

$$\nu_{ei}^* = \omega_L / 2 \sum_{\vec{k}} (\delta n_i(\vec{k}) / n_{cr})^2 \text{Im}(1/\epsilon(\vec{k}, \omega_L)) (\vec{k} \cdot \vec{E})^2,$$

where ϵ is the plasma dielectric function and $\langle \text{Im}(1/\epsilon(\vec{k}, \omega_L)) \rangle = 0.3$ for $n_e = 0.25n_{cr}$ for fluctuations with $k\lambda_D = 0.5$, λ_D being the Debye wavelength. The term $(\vec{k} \cdot \vec{E})^2$ is usually averaged over all $\vec{k}$ and taken to be 0.5. The fluctuation spectrum for a gas jet experiment similar to the one in this thesis has been measured by a Thomson scattering technique at the University of Alberta¹⁶ where $(\sum_{\vec{k}} \delta n_i(\vec{k}) / n_{cr})^2 = 0.04$. Should the same level of turbulence be present in the present experiment, this would result in $\nu_{ei}^* / \nu_{ei} = 1.20$ producing a 20% increase in the absorption level. Because of a lower temperature and higher electron density, the Alberta experiments demonstrate a 4X enhancement in the absorption over that which can be attributed to inverse bremsstrahlung alone.

Estabrook¹⁷ has shown, in one dimensional computer

simulations, that ion acoustic turbulence can produce a suprathermal temperature given approximately by

$$T'_e = T_e (1.5) / (1 - n_e / n_{cr}) .$$

For the electron densities of $0.25n_{cr}$ as encountered in the experiments of this thesis, the hot temperature due to ion acoustic turbulence can be expected to remain below one keV.

II-5 Thermal Conductivity

The measured electron temperature in this thesis is used to determine the energy content of the plasma. The electron temperature that is ultimately attained by the plasma depends on the thermal conductivity since for short lived laser plasmas, electron thermal conduction is the dominant energy transport mechanism. In this section, the thermal conductivity of a laser produced plasma is examined in view of some of the various phenomena that can effect it.

Since the electron thermal speed is always much larger than the ion speed, it is the electron thermal conductivity that is dominant. In the classical theory of electron thermal conduction, the heat flux is assumed to be given by Fourier's law, $\vec{q} = -\kappa \nabla T_e$. The most commonly used value for the thermal conductivity κ , is that of Spitzer and Harm¹⁸ where a perturbation approach was used to solve the Boltzmann equation with a Fokker-Planck collision term. The perturbation parameter used is λ_e / L_T , the ratio of the electron mean free path λ_e , to

the temperature gradient scale length $L_T = |T_e / \nabla T_e|$. When $\lambda_e \ll L_T$ and the partial time derivatives $\delta/\delta t \ll \nu_{ei}$ the heat flux is

$$|q| \propto (\lambda_e / L_T) q_{T \max},$$

which is equivalent to Fourier's law with the thermal conductivity

$$\kappa = 20(2/\pi)^{1.5} k_B (\epsilon \delta) \left[\frac{(kT_e)^{5/2}}{(m_e)^{1/2} e^4 Z \ln \Lambda} \right],$$

where ϵ, δ , are the Z dependent Spitzer factors for a non-Lorentzian gas with no net current flow. The quantity $q_{\max}$ is the flux of energy across a plane in a classical gas which is

$$\begin{aligned} q_{\max} &= 2k_B T_e n_e / 4 (8k_B T_e / \pi m_e)^{0.5} \\ &= 0.8 (n_e v_{th}) k_B T_e. \end{aligned}$$

Several effects can occur in a laser produced plasma that invalidate the theory of Spitzer and Harm. First, the perturbation approach to solving the Fokker-Planck equation fails for sufficiently large values of λ_e / L_T . Gray and Kilkenny¹⁴ examined Spitzer's¹⁸ calculation with the criterion that the perturbed electron distribution function become negative with v_e in the range $2.3 v_{th} < v_e < 3 v_{th}$. A negative distribution function is physically unreal and is a consequence

of the omission of higher order terms in the solution of the Fokker-Planck equation. The criterion determined by Gray and Kilkenny sets the limit:

$$q/q_{\max} = 2.4 \lambda_e / L_T < 0.036 ,$$

(for $Z = 1$). For the plasmas encountered in the experiments of this thesis, the typical radial temperature gradient scale length is

$$L_T = |T_e / \nabla T_e| = 300 \text{ eV} / (250 \text{ eV} / .05 \text{ cm}) = 0.06 \text{ cm} .$$

The temperature gradient considered is that in the radial direction perpendicular to the laser axis. In the axial direction, the plasma is heated uniformly by the laser and the gradients there are small as measured in a similar experiment by Wyndham et al.¹⁹ The electron mean free path is

$$\begin{aligned} \lambda_e &= \tau_{ee} v_{th} \\ &= 1.44 \times 10^{13} \frac{T_e^{3/2}}{n_e Z \ln \Lambda} \text{ cm}, \quad (T_e \text{ in eV}, n_e \text{ in cm}^{-3}) \end{aligned}$$

where Z is the charge state and $\ln \Lambda$ is the Coulomb logarithm ≈ 10 . For an electron temperature of 300 eV and $n_e = 2.5 \times 10^{18} \text{ cm}^{-3}$, $Z = 7$, $\lambda_e = 74 \mu\text{m}$. This makes $\lambda_e / L_T = 0.12$ so that the criterion of Gray and Kilkenny is not satisfied.

Bell et al.²⁰ extended the work of Spitzer and Harm by numerical simulations of the Fokker-Planck equation which examine the heat flow in temperature scale lengths a few times the electron mean free path. For the case of $\lambda_e / L_T = 0.12$ their

result is that the thermal conductivity is reduced from that of the classical value by a factor of 3.5. The experiments of Wyndham et al. suggest that for $\lambda_e/L_T = 0.12$ the reduction of the conductivity is a factor of 2.

In addition to the inadequacy of the Spitzer-Harm conductivity for situations of large thermal gradients, two other effects can further decrease the thermal conductivity. Since the Coulomb collision rate is proportional to v^{-3} , when $v_{os}/v_{th} > 1$ the collisional transport can be considerably modified. Secondly, if $T_e \gg T_i$ as is usually the case in laser produced plasmas, ion acoustic turbulence can occur, reducing the electron mean free path and the thermal conduction. Manheimer²¹ calculates the thermal conductivity in the presence of ion acoustic turbulence. Gray and Kilkenny present this result in the form of a ratio between the anomalous and Spitzer conductivity as

$$\frac{\kappa_{an}}{\kappa_{Spitzer}} = \left[\frac{1}{\gamma^2 n_e \lambda_D^3} \right] \cdot \left(\frac{n}{\Delta n} \right) \cdot \left[\frac{1}{320} \frac{\Lambda}{\pi \delta \epsilon} \right]$$

Assuming a $\Delta n/n = 0.02$ as measured by Al-Shiraida et al.¹⁶ and a $T_e = 300$ eV, $n_e = 2.5 \times 10^{18} \text{ cm}^{-3}$ this ratio is 0.04.

In summary, there are several effects all of which decrease the thermal conductivity from the classical Spitzer value. This has serious consequences in computer code simulations since the thermal conductivity is proportional to $T_e^{5/2}$, large temperatures can predict a physically unrealistic heat flux. In practice, the thermal flux used in simulations is an interpolation between

Fourier's law and the free streaming limit

$$\vec{q} = \left[\frac{1}{-\kappa \vec{\nabla} T} + \frac{1}{\phi \vec{q}_{\max}} \right]^{-1}$$

where ϕ is a factor varied to fit the observed electron temperatures with the results from simulations. Many experiments indicate $\phi = 0.03$, but values as high as 0.15 have been obtained¹⁹.

II-6 Electron Heating Model

An oversimplified model for the heating of electrons under the action of laser radiation is presented in this section. The purpose is to arrive at an estimate of possible electron temperatures for typical experimental conditions. The calculations shown below were initially performed for the case of laser interaction with a preformed z-pinch plasma column. However, with some interpretation the results can be applied to the interaction of intense laser radiation with a low pressure gas jet target as well.

In this model, a focussed laser beam impinges upon a uniform density plasma that is initially at 50 eV. This removes complicated initial effects of plasma formation from a neutral gas target, and also validates the use of the Johnston and Dawson inverse bremsstrahlung coefficient to calculate the laser absorption. For mathematical expediency, the energy losses due to volume expansion of the plasma are neglected, an assumption

partly justifiable on the basis that there is little expansion in the duration of a nanosecond laser pulse. The energy losses considered in this model are thermal conduction, radiation and energy transfer to the ions. The thermal conduction is that as calculated by Spitzer and initially no saturation effects are included in the absorption coefficient. The use of the Spitzer conductivity would tend to underestimate the electron temperature while the exclusion of bremsstrahlung saturation results in an overestimate. However, it is the general trend in such models to neglect these effects e.g. Handke et al.²² Burnett and Offenberger²³.

The model determines the change in T_e with time by means of the energy transport equation

$$\frac{d}{dt} \left\{ \frac{3}{2} n_e k_B T_e \right\} = I/\ell (1 - \exp(-k_{ib}\ell)) - \vec{\nabla} \cdot \vec{q} \\ - (P_{ff} + P_{fb}) - \frac{d}{dt} \left\{ \frac{3}{2} n_i k_B T_i \right\} ,$$

where k_{ib} is the inverse bremsstrahlung absorption coefficient, ℓ is the absorption length and P_{ff} and P_{fb} are the continuum radiative powers discussed and evaluated in Appendix B. Energy transfer to the ions is calculated by Spitzer's²⁴ equation:

$$\frac{dT_i}{dt} = (T_e - T_i) / t_{eq} ,$$

where t_{eq} is the equipartition time between electrons of mass m_e and ions of mass m_i :

$$t_{eq} = \frac{3m_e m_i k^{3/2}}{8(2\pi)^{1/2} n_e Z^2 e^4 \ln \Lambda} \cdot (T_e/m_e + T_i/m_i)^{3/2}.$$

Thermal conduction losses are given by the divergence of Fourier's law (for $Z=2$)

$$\nabla \cdot \vec{q} = \nabla \cdot \left(3.6 \frac{\nabla T}{T} 0.8 n_e k T_e v_{th} \right).$$

Axial gradients in temperature can be considered negligible in comparison to radial gradients (transverse to the beam axis). Therefore, the temperature gradients are estimated from either the geometry of the focal region or the size of the laser produced plasma. The temperature profile assumed has the form shown in Figure II-1 where the distance d is the diameter of the

Figure II-1 Form of the assumed radial temperature profile.

focal region or the plasma radius < 1 mm. The background temperature T_0 is set to 50 eV and T_e is the temperature

attained at time t into the laser pulse. With the temperature profile shown, the divergence of the heat flux may be approximated to first order as:

$$\underline{\nabla} \cdot \underline{\vec{q}} \propto \underline{\nabla} (T_e^{2.5} \cdot \underline{\nabla} T_e) \approx T_e^{2.5} \cdot ((T_e - T_0) / \frac{1}{2} d^2) \quad .$$

Further approximations include linearizing the exponential term $(1 - \exp(-k_{ib})) \approx k_{ib}$ which is justified on the basis that the scale lengths of the interaction are $\approx 1\text{mm}$ so that $k_{ib} \ll 1$. For simplicity, the laser pulse is approximated by a linear ramp pulse $I = I_0 \alpha t$, where I_0 is the maximum intensity in the pulse. The electron temperature is expected to follow the laser pulse intensity and so calculations are carried out to the time $t = 1/\alpha$, specifically to 1.3 ns.

The equations for electron heating are solved numerically using a second order Runge-Kutta method for various incident laser intensities. Thermal conduction begins to have a significant effect only at intensities greater than 10^{11} W/cm^2 . Due to the relatively long electron-ion equipartition times of $> 40 \text{ ns}$, there is little energy transfer to the ions and their temperature is seen to increase only a few eV. The corresponding reduction in electron temperature due to energy transfer to the ions is thus only a few percent. Similarly, the inclusion of radiation loss terms was found to produce a negligible effect on the resultant electron temperature.

The results of the calculation are shown in Figure II-2. The uppermost curve is calculated assuming no thermal conduction and other losses and is obtained by a straightforward

Figure II-2 Results from the numerical solution of the energy transport equation.

integration of the energy transport equation, yielding:

$$T_e = \left(\frac{5}{2} \left(\frac{2}{3} T_0^{2.5} + 1.18 \times 10^{-6} I_0 \right) \right)^{0.4}, \text{ at } t = 1.3 \text{ ns}$$

with T_e , T_0 in eV and I_0 in W/cm^2 . The calculation for Spitzer conductivity is shown as the center curve through the cross hatched region. The upper curve represents the results obtained with the thermal conductivity lowered from the Spitzer value by a factor of two as is suggested by the experimental results of Wyndham et al.¹⁹ The lower curve was calculated by allowing the inverse bremsstrahlung absorption coefficient k_{ib} , to saturate at the higher intensities according to the model of Faehl and Krueer⁸

$$k'_{ib} \rightarrow \frac{k_{ib}}{(1 + 7.02 \times 10^{12} I_0 / T_e)^{3/2}},$$

with I_0 in W/cm^2 and T_e in eV. Although the reduction in the collisional absorption is as high as a factor of 3 at intensities approaching 10^{14} W/cm^2 , the resulting electron temperature is not drastically reduced by the saturation effects because most of the heating occurs at early times at low intensities, when the absorption is much higher due to a lower electron temperature. Similar calculations were repeated for the case of a fully ionized nitrogen gas jet target which produced predicted electron temperatures in the range of 400 - 600 eV for laser intensities $10^{13} - 10^{14} \text{ W/cm}^2$. The increase in temperature in a nitrogen target compared to the helium is due to the linear dependence on Z in the absorption coefficient, whereas the thermal conductivity decreases with increasing Z .

The calculations presented here are crude in the sense that no time dependent ionization or hydrodynamic effects are considered. For the case of a preformed 50 eV target plasma the ionization issue can be ignored. However, for a laser produced plasma this may not necessarily be so. The hydrodynamics of the phenomenon may also influence the state of ionization which in turn determines the level of absorption and electron temperature and vice versa. Therefore, to model a laser produced plasma with considerable more finesse requires the simultaneous solution of a system of coupled hydrodynamic differential equations which

describe the conservation of mass, energy and momentum together with a system of equations describing the time dependent state of ionization. The complexity of the problem requires the use of large computer codes. For comparison with the predicted electron temperatures calculated in this section, the results of several computer code simulations are presented in the following section.

II-7 Results from Hydrodynamic Computer Codes

The implementation and interpretation of computer codes to laser-plasma experiments is a thesis topic in itself. Only a brief mention of the results of the several computer codes used will be presented here to obtain estimates of the expected electron temperature for the experimental conditions encountered in this work.

The one dimensional Lagrangian code MEDUSA has been available since 1974²⁵ and has been in use at UBC since 1980²⁶. A complete description of the code and its method of solution of the hydrodynamic equations is given in the above references. Of relevance to this thesis is that the absorption of laser radiation is by inverse bremsstrahlung and the thermal conductivity is the classical value of Spitzer. Unique to the UBC version of MEDUSA are the numerical changes effected by J. Kwan²⁷ to the inverse bremsstrahlung absorption coefficient and thermal conductivity. These changes brought the inverse bremsstrahlung coefficient into agreement with that of Johnston and Dawson and the thermal conductivity was corrected for

varying values of target Z .

Since MEDUSA is a one dimensional code, all the hydrodynamic phenomena can only take place in one dimension along the laser axis. Obviously, if radial plasma expansion is an important factor in the actual experiments, the output of the code will overestimate T_e and is then subject to interpretation. That plasma expansion cannot be neglected is evidenced by the fact that if the actual gas jet molecular density of 4×10^{-5} g/cm³ is used as an initial target density, the code almost instantaneously produces full ionization and an overcritical density target so that there is no transmission of the laser beam. Since the intention is to study the interaction of a laser beam with an underdense plasma, the initial target density was lowered by a factor of 10 to 3 to get resulting electron densities of 0.1 and 0.45 n_{cr} respectively.

Generally, the electron temperatures obtained from the MEDUSA runs are substantially over 1 keV as Figure II-3 shows. These high electron temperatures are a result of the fact that for an underdense plasma, there are negligibly small thermal gradients axially along the laser beam axis. Therefore, the thermal conduction losses are quite small resulting in high electron temperatures. It was observed that in order to reduce the resulting electron temperatures by a factor of 2, it was necessary to artificially increase the thermal conductivity by a factor of 100, thus indicating the small role of thermal conduction in a one dimensional model. The analytic solution for no thermal conduction from the previous section is also plotted in Figure II-3 for the case of 0.2 n_{cr} electron density.

Figure II-3 Temperature scaling with incident laser energy from MEDUSA computer runs.

Obviously, to model the interaction of a laser beam with an underdense plasma a two dimensional code should be used. N. H. Burnett of NRC has simulated the interaction using a 2-D Eulerian code for a nitrogen target with density of 9×10^{-6} g/cm³ and a laser intensity of 10^{13} W/cm². The result of interest is that at the peak of the pulse the electron temperature is 540 eV for a target of $0.25 n_{cr}$ initial electron density.

Another simulation was performed by E. Hiob using the well documented 2-D Eulerian code CASTOR²⁸. The initial conditions for the CASTOR simulation was a target density of 3.7×10^{-5} g/cm³ and a 6 Joule incident laser pulse of triangular temporal shape with a rise time of 2 ns. The beam focal radius was 100 μ m. The initial electron temperature was 1 eV and a time dependent atomic physics package²⁹ was implemented. The electron

temperatures along the laser axis at the time of the peak of the pulse were in the range of 500-700 eV, the axial density being $0.26 n_{cr}$. If the neutral gas density were to be instantaneously fully ionized, the mean electron density would be $1.1 n_{cr}$ however, the electron density remains underdense throughout the simulation. Initially, the level of ionization is low, the average Z being 4-5 at times up to 0.5 ns, after which plasma expansion is sufficient to keep the axial electron density substantially below the critical density even after near full ionization is achieved.

Therefore, the 2-D codes are in approximate agreement with each other in predicting an electron temperature of the order of 500 - 700 eV. The numerical solution of the energy transport equation with radial heat conduction also predicts electron temperatures in this range for laser intensities $> 10^{13}$ W/cm².

A shortcoming of all these predictions for the electron temperature is that they must assume some initial electron temperature otherwise, the inverse bremsstrahlung coefficient becomes infinite. The assumption of an initial electron temperature makes an imperceptible difference to the maximum electron temperature, but the initial ionization that arises from this initial temperature does create a preformed electron density profile which does not resemble the complicated multicomponent structure actually observed in the interferograms (Chapter IV). Secondly, the computer simulations do not self-consistently model the path of the laser through the ever-changing density contours. All of the incident laser light is sent through the plasma with no refraction taken into account. A

significant amount of the incident laser beam can in principle be refracted by a tenuous precursor upstream plasma reducing the laser heating of the denser core of the plasma.

Chapter III: The Apparatus and its Application

The main objective of this thesis is to obtain a complete energy balance of the laser-plasma interaction by experimental methods. The energy balance is used to explain the electron temperature obtained by an independent measurement. The means by which these measurements are made are indicated in this chapter.

Brief descriptions of the apparatus used in the experiments are presented in this chapter. The technical details are left to a minimum and instead the underlying physical principles and their relation to the design of the apparatus will be emphasized.

In simple terms, the experimental apparatus consists of a laser, a target and the associated diagnostics for the investigation of the laser produced plasma. The overall depiction of the experimental geometry is shown in Figure III-1. Each of the components of the experiment will be treated in detail subsequently in this chapter.

Some preliminary work involved a z-pinch device described by Houtman³⁰ and characterized by Hilko³¹. Attention quickly turned to a variable initial density gas jet target first described by Burnett³² and later utilized extensively by Offenberger et al.³³ The gas jet target will be described in section III-1 along with its characterization for use as a suitable target. A few pertinent results from the theory of supersonic nozzle flow are also mentioned since they determine the jet operating conditions.

The description of the CO₂ laser system will be

Figure III-1 Overall depiction of the experimental geometry.

appropriately succinct as many of the technical details of the components are contained in a UBC laboratory report³⁴ and a concurrent UBC thesis³⁵. The operation of the laser system in conjunction with the other components of the experiments is discussed in section III-2.

The laser produced plasma is investigated using several diagnostics. Apparatus was designed and developed by the author for the measurement of electron temperature by soft x-ray diagnostics, energy absorption by an Ulbricht integrating sphere, and plasma expansion by an image converter streak camera. These devices and their application in the experiments are treated sections III-3, 4 and 5.

III-1 Laminar Gas Jet Target

The variable density gas jet target was selected as the optimum target for the energy balance investigation. It provides a fairly uniform high-temperature, high-density, underdense plasma which is conducive towards absorption studies.

In contrast, the majority of laser-plasma interaction experiments in laboratories throughout the world are performed on various solid targets, largely in the interest of ascertaining the hydrodynamic efficiency i.e., the ratio of accelerated target energy to absorbed laser energy, (see e.g., Nishimura et al.³⁶) In such targets, the laser-produced-plasma density decays exponentially along the laser axis into the surrounding vacuum. Since different absorption and parametric processes occur in different density regions, it is desirable instead, to have a target of a single uniform density to simplify the laser-plasma investigation. To this end, gas targets have several advantages over solid targets. Firstly, at laser intensities of 10^{13} W/cm², the oscillatory electron energy $0.5m v_{os}^2$, (where v_{os} is the quiver velocity of the electron in the laser field), exceeds the ionization energy of most low Z gases so that 100% ionization can be anticipated in times less than the pulse duration. Thus the plasma density is roughly equal to the density of the neutral gas in which it is formed so that the plasma density can be controlled by varying the gas pressure. Gaseous targets consisting of large volume gas cells into which laser light is focussed were initially used by Yablonovitch³⁷ to elucidate resonant absorption phenomena in a

uniformly overdense plasma. Burnett³⁸ studied stimulated Brillouin backscatter in He gas maintained at pressures such that at full ionization, the electron density was $0.1 - 0.5 n_{cr}$. A major disadvantage of this method, however, is that in the critical density case, the laser intensity can never exceed that required to achieve ionization since any further increase in the intensity will cause the ionization front to move upstream with no intensity increase at the plasma front. Similarly, in the underdense target case, Burnett³⁸ found a complete absence of ionization downstream of the focal point and "explosive" refraction around the focal volume due to the cone-shaped ionization region ahead of the focus, terminating the interaction. To overcome the difficulties of plasma generated upstream towards the incident laser beam, the stabilized supersonic gas jet is used by Burnett³², Offenberger et al.³³ In this target, stable two-dimensional flow is obtained by discharging high pressure nitrogen gas through a convergent-divergent (Laval) nozzle into ambient He gas at a pressure of a few Torr. The 1 mm thick 1 cm wide jet of variable density simulates plane solid targets. Helium gas is used to stabilize the flow thus ensuring a low density plasma compared to that formed from nitrogen in the event of upstream breakdown. The details of the particular Laval nozzle are shown in Figure III-2. The particular design is due to N. H. Burnett of NRC.

For such a nozzle, the critical design feature is the radius of curvature in the diverging part and this is related to the desired Mach number as shown in Shapiro³⁹. The Laval nozzle consists of a pair of "jaws" set in a Lucite cone. This sits

Figure III-2 Details of the Laval nozzle.

atop a short brass pipe connected to a solenoid valve. A high pressure nitrogen reservoir of 0.5 liter volume provides the nitrogen gas. To facilitate precise filling of the required ambient He gas and reservoir pressures, an automatic gas handling system was incorporated Figure III-3 where electric solenoid valves are operated in a preset sequence by the electronics. The electronics senses the output from capacitive pressure transducers and allows the target chamber and nitrogen reservoir to be automatically refilled to the desired pressures whenever required.

A manual push button is used to activate the solenoid valve atop the nitrogen reservoir and establishes laminar flow. A piezo probe located near the nozzle senses the change in gas pressure and is used as a timing signature. The output pulse from the piezo probe enters a triggered variable delay pulse generator. The pulse from this delay unit is used to activate other time delay units for the operation of ruby or CO₂ lasers

Figure III-3 Electrical and gas systems of the target chamber.

as circumstances dictate.

The operating conditions of the jet are determined by the geometry of the jet nozzle and a few results from the gas dynamics of nozzle flow (see e.g. Landau and Lifshitz⁴⁰). In the convergent part of the nozzle the maximum amount of mass flux at the constriction is given by

$$Q_{\max} = \rho v S_{\min},$$

where $S_{\min}$ is the cross sectional area at the constriction. From the theory for streamline flow, the maximum flux occurs when the gas velocity is equal to the local velocity of sound such that the critical density and velocity respectively are:

$$\rho_* = \rho_0 \left(\frac{2}{\gamma + 1} \right)^{\gamma/(\gamma-1)}, \quad (1)$$

and

$$v_* = c_0 (2/(\gamma + 1))^{0.5}. \quad (2)$$

where γ is the ratio of specific heats. The flow becomes supersonic in the divergent part of the nozzle since $Q = \rho v S = \text{constant}$, throughout the nozzle. At the widest part, conservation of mass flux is written as:

$$S_{\min} \rho_* v_* = S_{\text{th}} \rho_{\text{th}} v_{\text{th}}, \quad (3)$$

where

$$v_{\text{th}} = c_0 \left\{ \frac{2}{\gamma-1} \right\} \left(1 - (\rho/\rho_0)^{(\gamma-1)} \right); \quad (4)$$

a general result for streamline flow. The subscripts $*$, o , and th apply to conditions at the constriction, the reservoir and the exit plane of the nozzle.

The substitution of equations (4), (2), and (1) into (3) produces the relation between the cross-sectional areas $S_{\min}$ and S_{th} and the reservoir and exiting gas densities:

$$\frac{S_{\min}}{S_{\text{th}}} \left[\frac{2}{\gamma + 1} \right]^{(\gamma+1)/2(\gamma-1)} \left[\frac{2}{\gamma-1} \right]^{-\frac{1}{2}}$$

$$= \rho_{th} / \rho_0 \left[1 - (\rho_{th} / \rho_0)^{\gamma-1} \right]^{\frac{1}{2}} .$$

For nitrogen, $\gamma = 1.4$, so that to a good approximation:

$$\rho_{th} / \rho_0 = S_{min} / S_{th} (0.259) = (P_{th} / P_0)^{1/\gamma} . \quad (5)$$

The required reservoir pressure to reach a certain exit pressure for a given area ratio S_{min} / S_{th} , is determined by the above equation. When the ambient He gas pressure is less than P_{th} then the jet pressure falls to the lower He pressure as the jet exits in an expansion fan that originates at the jaw edges. This phenomenon, called "underexpansion", manifests itself visually in the "candle flame" appearance of the jet boundaries as made visible by shadowgraphy in Figure III-4. This allows a

Figure III- 4 Ruby laser shadowgraphs of the gas jet target.

(a) $P_{th} = P_{He}$ (b) $P_{th} > P_{He}$

test of the steady state nozzle flow equation (5) as well as a verification of correct pressure settings to obtain an optimum target with parallel boundaries and subsequent uniform density.

Calculations based on the method of hodograph characteristics (Shapiro,³⁹ chapter 15) show that if the background pressure is lower than the exit pressure by 30% then the jet width will be roughly double its width at the exit at a distance 5 mm from the exit and thus be clearly visible. The background He pressures necessary to obtain parallel target boundaries for a given reservoir pressure for a jet with $S_{\min}/S_{th} = 0.0846$ (jaw separation = 0.0111cm) are shown in Figure III-5 where it is seen that there is good agreement between the experimental

Figure III- 5 Stabilized jet operating pressures determined from ruby laser shadowgraphy.

values and those predicted by equation (5) (solid line). Further shadowgraph series were taken for various $S_{\min}/S_{th}$ ratios by B. Hilko and agreement with theory was within 20%. More recently, Houtman and Meyer⁴¹ have utilized a multipass schlieren

technique to enhance the image of the density gradients which is necessary for low He pressures. Their results on the gas jet target as operated in the experiments verified that the jet target consisted of parallel laminar boundaries as expected.

Ruby laser Rayleigh scattering off the gas jet was also used to corroborate the above observations of the gas jet pressure. The intensity W of scattered laser light of incident power P is $W = \sigma_R n_0 P dz d\Omega$ where σ_R is the Rayleigh scattering cross section, n_0 the density of molecules, dz the length of gas the light is scattered from and $d\Omega$ the solid angle subtended by the collection optics. The Rayleigh scattering system shown in Figure III-6a is calibrated by filling the target chamber with a known pressure of N_2 gas and observing the scattered ruby laser signal level as detected by the collection optics and photomultiplier tube. The ruby laser scattered intensities from the gas jet target are then compared with the calibration data taking into account that the nitrogen emanating from the Laval nozzle is at 69 degrees Kelvin whereas the calibration data is for nitrogen at room temperature. The set-up for Rayleigh scattering and the results are shown in Figure III-6b. Once again, the experimental value of 10 ± 4 Torr for the gas jet is in good agreement with the exit pressure predicted by equation (5) for a nitrogen reservoir pressure of 2280 Torr.

Thus the molecular density of the gas jet can be determined by the reservoir pressure and jaw separation with considerable confidence.

(b)

(a)

Figure III-6 Rayleigh scattering on the gas jet.

(a) experimental set-up

(b) results, calibration and target shots

The CO₂ Laser System

Laser-plasma interaction phenomena consist of two basic types, fluid motion and microscopic interactions that occur in two different time scales. Hydrodynamic effects, such as heat and mass flow occur in time scales typically much longer than those in which the microscopic physics, absorption, parametric instabilities, etc., takes place. In order to investigate the laser plasma interaction in a uniform plasma it is desirable to minimize the complicating hydrodynamic phenomena. Therefore a short single laser pulse of a few ns duration is a useful tool for laser plasma studies.

The production of a high intensity, short-duration laser pulse is achieved in our system by using a fast electronically

controlled optical gate to "slice" a short pulse from a longer laser pulse and subsequently amplify it using other amplifier units. The system is designed with several criteria: efficient use of the gain medium, avoidance of exceeding damage thresholds for optical components, avoidance of parasitic oscillations, control of the backscatter beam and control of the beam quality. The first and second criteria are satisfied by the optical system which allows the laser beam to pass through each amplifier cavity twice with appropriate beam expansion and collimation upon each pass. Parasitic oscillations, i.e. self-lasing of the amplifiers, is suppressed by tilting the optical flats used as windows for the amplifiers and by inserting saturable absorber gas cells in the beam path. Saturable absorbers, e.g. SF_6 gas, suppress small amplitude oscillations but become bleached at large intensities allowing relatively unaffected transmission of the higher intensities. Spatial filters, each consisting of a pinhole aperture at the focal planes along the beam path, precede each amplifier stage. These pinhole apertures ensure a smooth radial intensity profile of the laser beam.

The CO_2 laser system incorporates the technology of transversely excited atmospheric (TEA) systems known since 1970⁴². The system closely resembles the 60 J NRC system described by Tan et al.⁴³ in that it consists of several stages: pulse selection, preamplification and main amplifiers. A schematic representation of the laser system is shown in Figure III-7 where the optical beam path through the various components is depicted.

Figure III-7 Schematic diagram of the optical lay-out of the CO₂ laser system.

The pulse switched out to the amplifier modules originates at the hybrid laser section. The hybrid laser⁴⁴ consists of a small TEA section with a low pressure CW section in the same laser cavity. The low pressure section is operated continuously so that there is a significant level of radiation of a single longitudinal mode at the time when the high gain pressure section is pulsed. This mode then grows first and depletes the available gain in the TEA section before other modes have built up appreciably. Thus the output from the hybrid laser is a high-power single longitudinal mode pulse superimposed on a low-power CW output.

In general, a CO₂ laser can be made to emit at least 400 individual vibration-rotation lines in the 8.7 - 11.8 μm wavelength region. Without any attempt to select the emitted wavelength, oscillation may occur on either one of several

lines, of which the P(18), P(20) and P(22) lines have the highest gain. For this reason, a temperature controlled Fabry-Perot etalon is used as the output coupler at the end of the CW section. By careful control of the temperature of the etalon and its alignment, the P(20) line ($10.58 \mu\text{m}$) is maintained as the emitted wavelength during cw operation. The P(20) wavelength has the optimum gain characteristics for subsequent amplification in the TEA amplifiers. Nonetheless, the stability of the emitted wavelength has a tendency to be somewhat transient so that the output of the CW section must be continuously monitored by a spectrum analyzer and the temperature and alignment of the etalon adjusted accordingly.

The polarization of the output beam of the hybrid laser is set by the intracavity Brewster windows and by the germanium polarizer immediately preceding a Pockels cell. Prior to the time of peak output of the hybrid laser gain-switched pulse, a co-axial cable spark gap is discharged. This provides a 2 ns FWHM high voltage pulse which is applied to the Pockels cell that consists of a slab of GaAs mounted in a 50 transmission line arrangement. The half wave voltage of 13 kV across the GaAs crystal causes the plane of polarization of the laser beam to be rotated by 90 degrees. As a result, the Germanium polarizer following the GaAs crystal becomes partially reflecting and switches out a pulse whose duration is equal to that of the voltage pulse applied to the Pockels cell.

The switched out 2 ns pulse is then directed into the Lumonics K103 preamplifier which has been discharged some few hundred nanoseconds prior to the time of entry of the pulse. The

switched out pulse of 100 μ J energy double passes through the active gain medium of the K103 amplifier so that its energy is augmented to about 100 mJ.

Some of the 100 ns duration gain switched pulse from the hybrid laser passes through the Pockels cell polarizer arrangement, and so an SF₆ saturable absorber cell is placed after the K103 amplifier to eliminate the low level "leak through" radiation. The 100 mJ beam from the K103 passes through this SF₆ cell and into a second spatial filter placed inside a multiple section TEA amplifier unit designated as the 3-stage amplifier in Figure III-6. The strategic location of the spatial filters in the amplifier chain also provides necessary overdense breakdown plasma isolators for any retro-reflected pulse which becomes amplified in traversing backwards down the laser chain.

The rear mirror of the three stage amplifier is contained in a gas cell which is used to prevent self-lasing of this amplifier and to provide preferential wavelength selection for the P(20) line. This gas cell is filled with 1 - 2 Torr SF₆, 20 Torr ethanol, 100 Torr Freon 502 and 640 Torr of He. The laser beam is double passed through both the three stage and Lumonix 600 modules. The exit beam radii of the 3-stage and Lumonix 600 amplifiers are 2 cm and 5 cm respectively. The input intensity to these units is sufficiently high so that on the return path they are operated in the saturated regime where a linear growth in energy density with distance traversed is obtained. All the available energy density per unit length is thus extracted. From the geometry and experimental gains measured by Houtman and Walsh^{3,4} the available extractable energy from the three stage

amplifier operated at 1 atmosphere pressure is 3 - 5 Joules, while that from the Lumonix 600 is about 15 Joules. Thus in theory, a total available energy of some 20 Joules can be extracted.

In actual practice, the typical observed energies are about 10 Joules. The relatively low yield can be attributed to the inherent problem of TEA lasers, which is arcing in the discharge. Normally, a glow discharge at atmospheric pressure is unstable and rapidly develops into an arc. However, the time to develop and arc is finite and hence arc formation can be inhibited by excitation of the gas mix with pulses whose duration is less than the arc formation time. Also effective and necessary is uniform preionization of the gas volume by UV radiation from an auxilliary low energy discharge prior to the main discharge as first demonstrated by Seguin and Tulip⁴⁵. Although these features are incorporated into the laser modules, arcing remains a sporadic phenomenon. To remedy this situation, a small controlled concentration of tripropylamine vapour is introduced into the continuous flowing gas mix. This impurity has low ionization potential and a quantum efficiency for photoionization of nearly unity. Seguin et al.⁴⁶ found that a trace amount of tripropylamine is sufficient to raise the photoelectron density by up to four orders of magnitude. Addition of tripropylamine into our laser system is found to be an effective arc suppressant at the cost of some absorption of the CO₂ laser beam by the vapour.

X-ray Diagnostics

Soft x-ray diagnostics are used extensively for electron temperature determination of laser produced plasmas since other methods such as Thomson scattering and spectroscopy are comparatively more cumbersome and unsuitable for short-lived small scale high temperature plasmas.

The method used to measure the electron temperature of the plasma in the experiments is the foil absorber technique pioneered by Jahoda et al.⁴⁷ on a theta pinch plasma. The electron temperature is inferred from measurements of the transmitted x-ray intensity through various metal foils. The technique relies upon the fact that a hot plasma of 100 eV electron temperature emits primarily continuum radiation. The peak of the continuum spectrum is at the wavelength

$$\lambda_p = 6200/T_e, \text{ \AA}, (T_e \text{ in eV})$$

and is in the soft x-ray range. Since the wavelength dependence of the continuum emissivity varies as $\exp(-hc/\lambda kT_e)$, it is a sensitive function of the temperature in the soft x-ray wavelength range.

Absorption measurements yield information about the spectrum of incident x-rays and hence the electron temperature by the shape of the plot of the transmitted intensity versus absorber thickness. The absorption coefficients of materials as a function of wavelength are very well known and documented (see e.g. Henke et al.⁴⁸) and have the general form of

$$\mu(\lambda) = C\lambda^n,$$

where n is of the order of 3 and C is a constant proportional to Z^a , a being determined empirically. The transmission of x-rays emanating from a plasma emitting bremsstrahlung and recombination radiation through a foil thickness x is expressed as

$$\begin{aligned} T(x) &= \int I(\lambda) \exp(-\mu(\lambda)x) d\lambda \\ &= \int A \exp(-hc/\lambda kT_e) \exp(-\mu(\lambda)x) d\lambda, \end{aligned}$$

where A is (see e.g. Donaldson⁴⁹):

$$A = \left(\frac{7.58 \times 10^{-18} n_e^2 V}{T_e^{1/2}} \right) \sum_p f_p Z_p^2 + \sum_p f_p \left[\frac{\chi_H}{T_e} \right] \sum_n \left(\frac{\chi_{p,n}}{\chi_H} \right)^2 \left(\frac{\xi_n}{n} \right) \exp \left(-\frac{\chi_{p,n}}{T_e} \right).$$

Here n is the electron density in cm^3 , V is the volume of the emission source, Z is the charge, n_i density of ion species p , f_p the fractional abundance ($=n_p/n_e$), n_e the density of ion species p , χ_H is the ionization potential (in eV) of hydrogen, $\chi_{p,n}$ is the ionization potential from the n th shell of ion species p and ξ_n is the number of states available in the n th shell for a recombining electron.

The form of the integrand $I(\lambda) \exp(-\mu(\lambda)x)$ is shown in Figure III-8 (Elton⁵⁰) for various thicknesses of Be and Al for bremsstrahlung emission from a plasma of $T = 500$ eV. The x-ray transmission (in arbitrary units), in Figure III-8 is shown as a

function of x-ray wavelength with foil thickness as a parameter.

Figure III-8 X-ray transmission through aluminum and beryllium foils of thicknesses indicated in mg/cm^2 .

Two features immediately become apparent upon examination of Figure III-8. Firstly, the overall transmitted intensity drops many orders of magnitude in comparison to the increase of absorber thickness. This allows the determination of the electron temperature by straightforward comparison of the fall of transmitted x-ray intensity as a function of absorber thickness. Secondly, for comparatively thin foils, most of the transmitted x-ray intensity lies in the k-edge wavelength of the absorber (e.g. for Al the k edge is $7.9 \text{ \AA}$). Thus using suitably thin foil absorbers of various differing Z allows a determination of the x-ray spectrum directly. Increasing the foil thicknesses diminishes the transmitted intensity but also shifts the peak of the transmitted spectrum towards shorter wavelengths. Thus, when a wide variety of foil absorbers are used, the wavelength response of the x-ray detector must be considered as well.

Of all commonly available x-ray detectors, the

photomultiplier-scintillator system has the most sensitivity per x-ray photon absorbed. Plastic scintillators (NE 110) were used to detect the transmitted x-rays and their visible light output in turn is detected by a photomultiplier coupled to the scintillator via a suitable light guide. Implicit in this technique is the assumption that the visible light output from the scintillators is proportional to the x-ray energy absorbed. Indeed, the response of plastic scintillators have been verified by several investigators^{51, 52} to be linear in the x-ray range of 0.6 - 15 keV. The transmittance properties of the foils together with the absorption properties of the scintillator determine the response R of the x-ray detection system

$$R(\lambda) = \exp(-\mu_a(\lambda)x_a) \cdot (1 - \exp(-\mu_{sc}(\lambda)x_{sc})) ,$$

where the subscripts "a" and "sc" refer to absorber and scintillator respectively. The significance of this result is that the x-ray detector has a finite wavelength range of sensitivity. This can be put into practical use by judicious selection of scintillator thickness x_{sc} and absorber thickness x_a such that $R(\lambda)$ is a maximum in a narrow wavelength range, (Bernstein⁵³). For laser produced plasmas which are commonly characterized by a two component temperature, the scintillator must be thick enough to absorb the x-ray photons being transmitted through the foils, yet thin enough so that the hard x-ray components from the suprathermal electrons do not interfere significantly when the softer part of the x-ray spectrum is being sampled. In the experiments described in this thesis, the thickness of the scintillators used was 0.6 cm so

that x-ray photons of energies less than 15 keV are efficiently absorbed.

Various x-ray detectors were constructed and utilized in the course of the experiments. The important design features common to all the arrangements used successfully in the experiments are clearly illustrated in the detector used for measurements on the z-pinch plasma column^{5,4} shown in Figure III-9. The z-pinch plasma served as a useful test-bed for the

Figure III-9 X-ray diagnostics set-up for the z-pinch.
 (a) overall set-up
 (b) detector head detail

development of x-ray diagnostics of high sensitivity. Since the z-pinch plasma and the gas jet plasma are both not particularly copious sources of x-rays, a large detector area = 5 cm^2 is exposed to the x-ray source for optimum sensitivity. Since delicate foils of thicknesses as low as $5 \text{ }\mu\text{m}$ are used, provision must be made for evacuation of air trapped behind the foil and the scintillator to prevent foil rupture upon evacuation of the

target chamber. The foils must be directly placed against the scintillator so that they are supported against the blast from the collapsing pinch or laser produced plasma. A 0.5 mm vacuum channel drilled into the scintillator allows escape of any air entrapped between foil and scintillator. The scintillator is cemented with index matching epoxy onto a Lucite light guide 60 cm in length which in turn is optically coupled with Dow Silicone 200 fluid onto the face of an RCA photomultiplier tube. The measured light attenuation from the light guide was 50 %. To reduce electrical noise from the plasma and laser triggering Krytron discharge units from interfering with the photomultiplier signals, the photomultiplier is encased in a doubly shielded enclosure. The outermost part of the electrical shield is an extension of the screened room housing the experiment's electronic instrumentation. The signal ground and the screened room extension are electrically isolated and commonly grounded at the screened room. The electrical noise level is thus reduced from several volts to a tolerable level of 10 mV.

Fluorescence of the detector from irradiation by plasma bremsstrahlung was kept to a minimum by the exclusive use of low Z materials in the construction of the x-ray detector. For example, the fluorescence conversion efficiency for brass compared to aluminum is twenty-five⁵⁴ times larger.

The x-ray detection apparatus and the principle of this temperature measuring technique was confirmed by determining the electron temperature of the z-pinch plasma. The results obtained are shown in Figure III-10 where the relative transmitted

Fig III-10 X-ray transmission data for the z-pinch plasma.

intensities through various aluminum foils are plotted against the theoretical relative transmissions for various electron temperatures. It is seen in Figure III-9 that the results suggest a T_e of 40 eV which is in agreement with the temperatures previously obtained by Thomson scattering, spectroscopy and other means⁵⁶.

The design of the x-ray detector for the laser incorporates the same design features as for the z-pinch detector but involves a few more considerations mainly due to the problems that arise from suprathermal electrons. The energetic electrons which are attributed to various parametric processes generated by the laser plasma interaction, can have mean free paths exceeding the dimensions of the vacuum chamber, enabling them to strike the surfaces producing characteristic K_α radiation. These electrons must also not be allowed to reach the x-ray detectors where they could penetrate the absorber foils. Both of these effects unless controlled, can produce a additive

background level to the observed signals. When determining the electron temperature by taking signal ratios, the presence of a background level would lead to an erroneously high figure for the electron temperature (see e.g. Lachambre et al.⁵⁷ and Martineau et al.⁵⁸).

The x-ray detector system used for the electron temperature measurement of the laser produced plasma from the gas jet target is shown in Figure III-11. To facilitate an electron temperature

Figure III-11 X-ray diagnostics set-up for the gas jet target.

measurement on a single shot basis, there are now four x-ray channels each with a scintillator, 1 cm diameter fiber optics light guide and an RCA 8575 photomultiplier tube. The scintillator-foil assemblies are contained in a black Lucite 8 cm diameter disk which is also enclosed in a tubular Lucite

enclosure that serves as shielding of the detectors from the electrons. The tubular housing also contains a pair of permanent bar magnets behind an aperture. These magnets provide a field of 700 gauss over a distance of 5 cm so that in the geometry used in the experiments, electrons of energies less than 10 MeV are swept away from the detectors.

Linearity of the photomultipliers was established using a co-axial spark gap light source whose duration was limited to 20 ns by the cable length. The saturation signal level of the photomultipliers was determined by the observation of the peak signal output voltage as a function of the light intensity varied by the inclusion of neutral density filters as shown in Figure III-12. All four signal outputs from the photomultipliers

Figure III-12 Photomultiplier Characterization

(a) set-up

(b) p.m. response scaling with light intensity

are delayed with respect to one another using cable delays, terminated by 50 Ohms and then displayed on a Tektronix 7904

oscilloscope with 7A12 and 7A24 dual trace plug-in amplifiers operated in the adding mode. The time response of the system was checked by injecting a 680 ps pulse from an N₂ TEA laser⁵⁹ of UV wavelength 337.1 nm onto the scintillator surface. The observed penetration depth of the beam is about 2 mm into the scintillator causing it to fluoresce. The resulting signals as measured on the same oscilloscope as the x-ray signals shows that the FWHM of the response is 10 ns. Since the laser pulse is only 2 ns there will be considerable temporal smoothing of the detected x-rays. In addition, there is spatial smoothing as well since the only spatial discrimination is that provided by the aperture and the solid angle subtended by the x-ray detectors. Both temporal and spatial smoothing properties of x-ray diagnostics for short laser pulse plasmas have been considered by Puell et al.⁶⁰ who concluded that a temperature equal to the maximum temperature is obtained by the ratio method if the transmission of the foils used is not greatly different from one another. A similar conclusion was reached for the spatial integration assuming reasonable density gradients for the plasma and considering that bremsstrahlung power is proportional to the electron density squared.

However, the foil absorbers used in the experiments in this thesis have a range of different transmission characteristics so that the temporal and spatial smoothing effects as they affect the experimental results will be discussed at some length in Chapter V.

Ulbricht Integrating Sphere

As part of the energy balance, the scattered fraction of laser light that does not contribute to plasma heating must be known. Optical energy balance measurements are the most direct means of gaining this information. In laser produced plasma experiments, a large fraction of the non-absorbed laser light may be scattered in a diffuse manner. Hence, a reflectometer to be used in such experiments should have a 4π collection angle. An Ulbricht spherical photometer⁶¹ is ideal for such a measurement and has been first used by Godwin et al.⁶² for a laser produced plasma. It basically consists of a hollow sphere with its interior surface fashioned to be a diffuse reflector. The attractive feature of this device is that the measured intensity, which is proportional to the reflected laser light, is independent of the angular distribution of the radiation.

The photometric properties of the Ulbricht sphere are immediately apparent from the general laws of radiation, namely that the intensity of illumination of a surface is inversely proportional to the square of its distance from the source and that the luminous intensity of an element of a diffusing surface is directly proportional to the cosine of the angle between the normal to the surface and the line of emission⁶³.

In the schematic diagram of Figure III-13, the flux per unit solid angle falling normally onto the surface at point x' from the surface element dS at x with normal luminance b is $b dS \cos\theta/d^2$.

The area element at x' is at an angle θ to the line of

Figure III-13 Schematic of an Ulbricht sphere.

sight xx' so that the illumination at x' , I due to the element dS , is $I = b \, dS \cos\theta \cos\theta / d^2 = b \, dS / 4R^2$ since $d = 2R \cos\theta$ by geometry of the sphere. As x is an arbitrary point on the sphere, then it follows that any element of the surface is illuminated by all others equally. Therefore, the illumination of the surface by reflected light only is equal at all points regardless of any asymmetry of the light emitted by the source.

If I_d is the average illumination of the sphere by direct light and I_m is the illumination by diffuse reflected light then the average total illumination is

$$\begin{aligned}
 I &= I_d + I_m \\
 &= I_d + mI_d + m^2 I_d + m^3 I_d + \dots,
 \end{aligned}$$

where m is the reflection factor of the sphere's surface. Since the number of reflections is infinite and $m < 1$, then $I = I_d m / (1 - m)$.

A more thorough treatment⁶⁴ includes consideration of light lost through the necessary ports of the sphere. The main result here is that the throughput of the sphere, i.e. the ratio of

exiting flux to that entering the sphere, is

$$\tau = \frac{f_e m}{1 - m(\sum_i f_i)} ,$$

where the factor f_e is the ratio of the exit port area to the area of the sphere surface, and $\sum_i f_i$ is the ratio of the summation of all the port areas to that of the sphere surface.

The use of Ulbricht spheres for CO₂ laser produced plasmas has been comparatively restricted to a few experiments in the published literature: Garban-Labaune et al.⁶⁵, Kristal⁶⁶, Nishimura et al.⁶⁷. For CO₂ laser produced plasmas the reflected light is infrared so that the interior surface of the sphere must be sand-blasted and gold coated to provide a diffusely reflecting surface. By its very nature, the Ulbricht sphere is a time integrating device. Kristal⁶⁶ has measured the decay pulse shape and found it to be of the form $I(t) = I_0 \exp(-t/\tau)$ where τ is the decay constant approximated by T/α , T being the representative transit time of the sphere, α is the loss per bounce. The time integrating feature of the Ulbricht sphere can in principle lead to an overestimate of the reflected laser energy since the energy radiated in bremsstrahlung can contribute to that attributed to reflection alone. However, as shown in Appendix B, the power radiation losses from the observed plasma are 3 orders of magnitude smaller than the laser power input and the reflected (refracted and or scattered) infrared radiation is of the same order as the input power. The Ulbricht sphere used in the experiments is depicted in Figure III-14. It is comprised of two 18 cm diameter aluminum hemispheres the interior of which was sandblasted and uniformly

Figure III-14 Ulbricht sphere photometric set-up.

gold-coated by vacuum evaporation. The incident laser beam enters and exits the sphere through two 2 cm ports. The exit port is necessary to avoid high intensity transmitted laser radiation from damaging the gold surface. The incident laser power and the transmitted energy are measured for each shot by the arrangement shown in Figure III-15. A beam splitter in front of the focussing lens sends a small fraction of the incident laser beam, as well as the retro-reflected light collected by the focussing lens, onto a photon drag detector through suitable IR attenuators. The photon drag detector for monitoring the incident laser energy is calibrated periodically relative to the Apollo energy meter which is used for the transmitted energy measurement. Calibration is achieved by direct comparison of the photon drag detector signal relative to the Apollo energy meter when the latter is placed in front of the target chamber. The actual details of the IR incident and backscatter radiation are

Figure III-15 Arrangement for the infrared energy inventory.

in the thesis of Bernard³⁵. The reflected laser energy from the plasma is measured by a Gen-Tec ED - 200 energy meter placed directly above the gas jet as shown in Figure III-14. A NaCl salt flat positioned in front of the Gen-Tec meter served as a filter preventing direct heat transfer from the plasma to the energy meter. To check that visible radiation from the plasma makes no contribution to the measurements, the filtering salt flat was replaced by a Plexiglas disk. No detectable signal was observed with the plasma present and with the inclusion of the Plexiglas disk.

Calibrating the Ulbricht sphere essentially involves firing the CO₂ laser beam into the evacuated sphere with the rear port covered by a matching cap, the inside concave surface of which was also sandblasted and gold coated. Comparison of the Gen-Tec signal to the incident beam photon drag signal under these

circumstances establishes the calibration of the Ulbricht sphere. The ratio of incident laser energy to that measured by the Gen-Tec meter was determined to be 74 ± 2 .

Isotropy of the sphere was checked by tilting the back port cap at various degrees of incidence to the incident laser beam and observing if there was any change in the calibration of the Gen-Tec signals. No backscatter signals were observed during the calibration ($\leq 0.01I_0$) to within the accuracy of the measurements. The integrating sphere as used appeared to be entirely symmetric in its radiometric properties.

Image Converter Streak Camera

Although the time duration of the laser-plasma interaction is short-lived (< 4 ns), the plasma persists for a long time afterward. Observations of the light output from the plasma by a Hamamatsu vacuum photodiode indicates that the visible emission has a $1/e$ decay time of about 60 ns. It was therefore desirable to obtain a time resolved record of the growth of the plasma. In particular, the amount of absorbed laser energy can in principle be determined by comparison of the radial growth of the plasma with theoretical blast wave space-time curves. (see e.g. Leonard and Mayer⁶⁸). For this purpose a streak camera unit was constructed which utilized an RCA C73435 image converter tube with a co-axial cable discharge circuit which provides the gate and deflection plate voltages. A schematic diagram of the streak camera unit is shown in Figure III-16. A similar circuit of this design was first used by Cavenor and Meyer⁶⁹ in the study of

al.⁷¹). The mode-locked output train of pulses are separated in time by about $2\ell/c$, where ℓ is the cavity length and c the speed of light. The streak speed of the camera is established once a streak record of the mode-locked train of pulses is obtained. For the calibration, the streak camera spark gap circuit was triggered by a Krytron circuit activated by a variable time delay pulse unit which itself was activated by the time delay unit used to operate the ruby laser.

A bright light source of time duration > 100 ns was necessary for backlighting the plasma for time resolution of the radial growth of its shadow. Since Q switched laser pulses are limited to 60 ns at best and have as much temporal jitter, a triggered high pressure spark gap light source was constructed. The spark gap light source was powered by a low impedance 1 Ohm transmission line as shown in Figure III-17. The background light source was measured to have a decay time of 500 ns. Both the light source and the streak camera are triggered by a 2 ns 13 kV pulse from the hybrid laser Pockels cell driver system via suitable cable delays to compensate for the time of arrival of the CO₂ laser pulse and laser triggering.

Studies of the plasma expansion on faster time scales were performed using a Hamamatsu M1763 temporal disperser unit. The high gain of this unit as provided by the inclusion of a micro-channel plate in the streak tube, allowed sub-millimeter spatial resolution. Plasma expansion was recorded up to 5 ns after the onset of the plasma, which is the upper limit of the time base of this unit. A schlieren system was also set up and used in these studies. The schlieren system essentially consisted of a

Figure III-17 Arrangement for time synchronization of the spark-gap light source for shadow streak photography.

submillimeter foil speck placed in the focal plane of the imaging lens as shown in Figure III-18. The turning mirror for

Figure III-18 Set-up for time-resolved schlieren photography.

the ruby laser beam was mounted on a gimbal mount and allowed a vertical displacement of the image on the streak camera slits so that the radial expansion of the plasma at different points along the focal region could be studied.

Chapter IV: Results from the Experiments

In order to examine how the laser energy is coupled to the plasma, several of the physical parameters of the plasma in addition to the energy absorption must be measured. The quantities of primary importance are the electron temperature, the density and size of the plasma. This chapter reports the results of the experiments using the apparatus and techniques described in Chapter III. The discussion of the results here will be restricted to the possibilities of systematic error and the validity of the results. The plausibility of the experimental results can be ascertained to some degree by comparison with the available published data from other laser laboratories. A more detailed discussion of the significance of the experimental results are left to Chapter V.

IV-1: Electron Temperature Measurements

This section describes the electron temperature measurements obtained by the x-ray absorption technique described in Chapter III. Initial observations are reported that were made at comparatively low laser intensities impinging on a low density target such that no hot electron generation was observed. The interaction is thus simplified by the absence of the complicating feature of hot electron generating parametric instabilities. A scaling of electron temperature with incident laser energy is obtained for the interaction of CO₂ laser light

with an underdense plasma.

The experimental conditions were then changed so that production of hot electrons is observed since this thesis work is part of an overall program initiated to study hot electron generating instabilities. The corresponding changes in the x-ray emission character under these circumstances are noted accordingly. The presence of hot electrons introduce several problems that are considered and evaluated in terms of their consequences on the measurements. X-ray absorption data obtained with a variety of x-ray absorbers provide x-ray spectral information which is presented and summarized.

IV-1a: Initial Observations at Low Intensities

As mentioned in Chapter II, the electron temperature of a laser produced plasma is expected to depend on the incident laser intensity. As reviewed by Donaldson et al.⁷², the dependence of electron temperature on the incident laser flux intensity can be obtained from relatively simple hydrodynamic analytical models. These result in scaling laws of the form $T_e \propto I^n$ where I is the laser flux, T_e the temperature, and n varies from $4/9$ to $2/3$ depending upon the assumptions employed. The models have been used with some success to explain the observed electron temperature scaling obtained in planar solid target experiments. The gas jet simulates a solid planar target in geometry, but with the important difference that there is no critical density layer and no consequent shock heating of the target material behind it. Thus, it is of interest to compare

the scaling with incident laser flux for the underdense gas jet target.

For the scaling study of electron temperature, a 2 channel version of the x-ray detector described in Chapter III was used, the electron temperature simply being determined by the ratio of the x-ray intensities transmitted through two suitably different absorber foils. The gas jet target was operated at low pressure conditions such that if full ionization were to occur as expected, the maximum electron density would be $5 \times 10^{18} \text{ cm}^{-3}$. Since the range of accuracy of temperature determination from any pair of foils is limited as shown in Figure IV-1, various aluminum foil pairs in the range of thickness from 6 - 18 μm were selected for different temperature ranges. The results are

Figure IV-1 Ratio of transmitted x-ray intensities through 9 and 18 μm foil vs. T_e

summarized in the plot of Figure IV-2 where the solid line is the fit of $T = 41E^{0.60}$, E being the incident laser energy. The intensity scaling of 0.6 can be compared with that from other laboratories. Direct comparison should be approached with some

Figure IV-2 Electron temperature scaling with incident laser energy. Laser intensities $\leq 10^{13}$ W/cm².

caution since different experiments use different conditions such as focal spot size and time duration of pulse. Nonetheless, it is of interest to mention a few of these results. There are two parameters to be compared here, electron temperature and the scaling with laser intensity. In terms of pulse duration and intensity the results obtained here can be compared to those of Pepin et al.⁷³ where a 4J, 1.8 ns CO₂ laser pulse was focussed onto a polyethylene target. The reported temperatures ranged from 150-300 eV and the intensity scaling is 0.3 - 0.4. Yamanaka⁷⁴ summarizes the results of the Ecole Polytechnique and Limeil laboratories as well as that of Osaka, all of which show a scaling of $I^{0.25}$ and a temperature range from 100 eV at 10^9 W/cm² to 500 eV at 10^{14} W/cm². More recently, Nishimura et al.⁶⁷ have reported a 0.6 power scaling for the black body temperature of CO₂ laser irradiated gold targets.

The observed temperatures from the gas jet are therefore

within the same range obtained for similar intensities on solid targets. In addition to the above mentioned data, Enright et al.⁷⁵ reports a temperature T_e of 250 eV from a CO₂ laser produced polyethylene plasma at intensities of 10^{14} W/cm². Similarly, Violet and Petruzzi⁷⁶ find $T_e = 300 - 500$ eV for 10 - 40 Joule CO₂ laser pulses of 1 - 3 ns duration on planar polyethylene targets. Recent results in the Soviet literature⁷⁷ indicate a thermal temperature of 220 eV obtained under similar conditions as those above.

It should be mentioned that the majority of x-ray emission in solid targets comes from the conductively heated region behind the critical density layer. The computer simulations of Marsh and Thompson⁷⁸ for laser irradiation of a carbon slab reveal that the supercritical conductively heated region is characterized by a Planckian spectral distribution corresponding to temperatures of < 200 eV. Preceding the critical density layer is the tenuous optically thin corona plasma which is laser heated to several keV producing a bremsstrahlung spectrum in the photon energy range of 2 keV and beyond.

Only a few results of electron temperature measurements on gaseous targets irradiated by CO₂ lasers have been published. Giles and Offenberger⁷⁹ measure 150 eV in an oxygen gas jet target irradiated by a 40 ns pulse of irradiance 6×10^{12} W/cm². Baldis and Walsh⁸⁰ measure a Brillouin backscatter red shift corresponding to a 200 eV plasma of electron density $0.1 - 0.2 n_{cr}$. In contrast, Yablonovitch³⁷ claims an electron temperature of 1 keV for a nitrogen plasma using only 0.1 J of input energy. Yablonovitch's results are particularly interesting in that a

resonant nature in the x-ray production is observed. As the gas jet target pressure is increased to a density where a critical electron density is expected for full ionization, a sudden onset of x-ray and electron production is observed (Kolodner and Yablonovitch⁸¹) and is attributed to the resonance absorption mechanism.

IV-1b: Change in X-ray Emission with Target Conditions

Similar behaviour as reported by Kolodner and Yablonovitch was observed with the gas jet target with an increase of target density and a subsequent improvement in the focal spot radius. The focal spot radius was measured by the multiple slit Fraunhofer diffraction burnmark technique as described by O'Neil et al.⁸² and Bernard³⁵. The focal spot was reduced from 250 μm used in the x-ray absorption measurements described earlier to 50 μm thus increasing the focal intensity to $\leq 10^{14}$ W/cm².

A critical density in principle can be reached for full ionization if the pressure of the nitrogen emanating from the nozzle is 4 Torr or greater. Indeed, when the jet pressure was increased to 4 Torr, the x-ray signal levels through a 9 μm foil increased by a factor of 45 from that observed from the 2 Torr jet target. At this higher initial target density, the x-ray signals had a tendency to fluctuate considerably from shot to shot resulting in a large scatter of the data. There also appeared to be no apparent scaling of the electron temperature with incident laser energy. The x-ray signals observed through 300 μm of Al with a 0.3 N.D. filter for various jet target

pressures are shown in Figure IV-3. Energetic electron emission

Figure IV-3 Hard x-ray intensity as a function of target pressure

was observed simultaneously with an electron spectrometer by McIntosh⁸³. These results showed a similarly striking increase as the x-rays with target pressure. The subsequent investigations were all performed on a 5 Torr jet target as the upper limit since higher target pressures resulted in a breakdown He plasma without the presence of N₂. For the 5 Torr jet target, the two channel x-ray detector measurements consistently implied a constant temperature of 2 keV irrespective of the incident laser energy. Some representative data is shown in Figure IV-4 which is the temperature obtained by means of a 100 μ m and 300 μ m aluminum foil pair.

The sudden increase in electron production and x-ray intensity with target pressure along with the observed high electron temperature suggests the onset of parametric

Figure IV-4 Electron temperature versus laser energy.
Laser intensities $\leq 10^{14}$ W/cm²

instabilities causing the production of energetic electrons. It has been shown by other investigators, Martineau et al.⁵⁸, Lachambre and Neufeld⁵⁷, Pepin and Baldi⁸⁴, that in the presence of energetic electrons the two foil ratio method of electron temperature determination can be greatly erroneous as different electron distributions produce different spectra of x-ray radiation according to their energy distributions. This fact is dramatically demonstrated by the results of Dyer et al.⁸⁵ who obtained 100 eV to 2600 eV from the same laser plasma using various pairs of foils.

IV-1c: The Effect of Hot Electrons

Of primary concern is the effect of electrons bombarding materials in the target chamber. Eidmann and Sigel⁸⁶ demonstrated this effect as being important since proper shielding of their x-ray detectors from the radiation emitted by

the target chamber walls reduced the signal yield by a factor of 3 -10. In the physics of bremsstrahlung emission, there are two relevant mean free paths: the mean free path for binary collisions giving bremsstrahlung and the mean free path for scattering through 90 degrees by multiple small angle Coulomb collisions. This second mean free path is typically larger than the first and determines whether an electron can be a runaway electron (Spitzer²⁴). The value of this mean free path can be calculated from Spitzer's equation

$$\lambda = \frac{E^2}{2\pi n_e e^4} \frac{1}{\ln \Lambda}$$

$$\approx 7.68 \times 10^{11} E^2 / n_e, \text{ cm, (E in eV, } n_e \text{ in cm}^{-3}\text{)}$$

for electrons of density n_e , energy E and Coulomb logarithm $\ln \Lambda$.

For electrons of energy 2 keV, the mean free path is over a cm for a typical density of $2.5 \times 10^{18} \text{ cm}^{-3}$. Typical laser produced plasmas are much smaller than this so that electrons of energy of a few keV produce no bremsstrahlung in their passage through the plasma. Photons with energies in this range are not produced in the plasma. The remaining origins are then either electron collisions with the N_2 gas surrounding the plasma or more likely, with materials in the target vacuum chamber. It is well known that the interaction of fast electrons with the target material is the same as that of fast electrons stopping in non-ionized material as far as bremsstrahlung and K_α emission are concerned for radiation of photons of high energy i.e., $h\lambda \gg kT_e$ (McCall⁸⁷). Thus electrons bombarding neutral nitrogen yield the same radiation as if the nitrogen were ionized. Accordingly, to maintain this physical equivalence of radiation

throughout the target chamber, all metal surfaces in the vicinity of the focal volume were covered by a layer of Plexiglas.

Energetic electrons have been observed to leave the focal volume.^{88, 89} These electrons are characterized by an energy of 100 keV and were estimated to constitute 1% of the total number of electrons in the focal volume. The radiation contribution of these electrons impinging upon materials in the target chamber was considered as follows. The x-ray production from a target with atomic number Z is given by⁹⁰

$$\text{x-ray power/beam power} = 7 \times 10^4 Z E ,$$

where E is the electron energy in MeV. With the assumption that the 100 keV electrons are emitted in a pulse duration of 2 ns, the x-ray power from these electrons bombarding the Plexiglas in the chamber is estimated to be $10^2 - 10^3$ Watts. In comparison, the total integrated bremsstrahlung (Jahoda⁴⁷) is

$$\int_{\lambda} E d\lambda = 1.5 \times 10^{-32} T_e^{0.5} g N_e N_i \text{ Watts, } (T_e \text{ in eV}),$$

which for T_e of the order of several hundred eV and total electron number of 10^{16} yields an estimate of just a few watts, an insignificant power compared to that produced by electron bombardment of Plexiglas. The Gaunt factor g above, is ≈ 1 .

The shortest wavelength that can be emitted by an electron of energy E is the Duane-Hunt limit:

$$\lambda_{\min} = hc/E ,$$

which is 0.01 nm for 100 keV electrons. As this wavelength is considerably shorter than that emitted by thermal electrons in the plasma, the shape of the spectrum of x-rays emitted by the energetic electrons becomes an issue. Several results from the theory of production of x-rays from x-ray tubes become useful (Green and Cosslett⁹¹). The number of photons N emitted as continuum radiation in the wavelength range $\lambda_{\min}$ to λ_b is given by the integration of Kramers' formula:

$$N_c = m \cdot 1.8 \times 10^{-14} \frac{Z}{\lambda_{\min}} \ln \left[\frac{\lambda_b}{\lambda_{\min}} \right] + \left[\frac{\lambda_{\min}}{\lambda_b} \right] - 1$$

where m is the number of bombarding electrons per second. Calculations reveal 10^8 photons/sec with the peak of the continuum spectrum at $2\lambda_{\min}$. Green and Cosslett obtain a relation for the number of photons/sec N_k of K_α line radiation which has the form:

$$N_k = \omega_k mA \times \frac{E}{E_k} \ln \left\{ \frac{E}{E_k} \right\} - \left(\frac{E}{E_k} - 1 \right) ,$$

where ω_k is the fluorescence efficiency and A is a constant. Both ω_k and A are target dependent and are tabulated in Green and Cosslett. For a carbon target the calculations show that 10^{18} photons/sec are produced at the carbon K_α line of 4.4 nm. Therefore, the hard continuum radiation produced by the bombardment of energetic electrons is not a significant

contribution to the detectable radiation. The soft x-ray radiation of the K_{α} lines of low Z material need not be considered since the transmission of the foils used falls off sharply at 1.0 nm. For example, the fraction of 4.4 nm K radiation transmitted through a 5 μ m Al foil is $\exp(-(\mu(\lambda)x)) = 10^{-19}$.

The effect of high Z material in the vicinity of the Laval nozzle has been observed in the experiments. A factor of 2 decrease in the x-ray signal intensity was seen to occur when the stainless steel jaws of the nozzle were covered with Plexiglas.

IV-1d: X-ray Measurements at High Laser Intensity

An extensive series of observations of the 5 Torr N_2 jet x-ray emission were made using aluminum foil filters with the four channel detector. The effectiveness of the electron screening by the magnets was checked by observing the relative x-ray signals before and after rotating the magnet assembly by 90 degrees. Upon removal of the magnets from the detector the photomultiplier tube signals were observed to saturate, indicating that the inclusion of the magnets is necessary. Light leaks through the foils were checked by the replacement of the magnet assembly by a 6 cm thick clear Plexiglas cylinder which effectively attenuates electrons and x-rays but is transparent to the visible light. The dynamic range of the photomultipliers was extended by the inclusion of gelatin neutral density filters at the ends of the optical fibers leading to the

photomultipliers. Calibration of the detectors relative to one another was performed whenever the neutral density (N.D.) attenuation was changed. This was done by having the same thickness of foil on each channel and observing the x-ray pulses from the plasma. Because of the fluctuation in x-ray intensity from shot to shot, many x-ray signals were recorded and analyzed for a particular laser energy. All the aluminum foil data consistently yielded an electron temperature of 2 ± 0.4 keV irrespective of incident laser energy. The observations were extended to include the transmissions through copper foils as well. The data are summarized in Figure IV-5, where the error

Figure IV-5 X-ray transmission data from aluminum (o) and copper (Δ) foils.

bars denote the standard deviation of the mean of ten or more shots. The copper foil data are plotted on the same graph by

calculating the cut-off thickness of the copper foil and placing each copper datum on the abscissa corresponding to an equivalent thickness of aluminum foil required to achieve the same cut-off energy. The cut-off energy is defined as that photon energy at which the attenuation of transmitted x-rays is $1/e$ i.e., $\mu(E) \cdot x = 1$, where x is the foil thickness and $\mu(E)$ is the photon energy dependent mass absorption coefficient of the foil material.

The observed temperature of 2 keV is much too high to be the thermal temperature of the bulk of the electrons in comparison with other CO₂ laser-plasma experiments using similar irradiance. The results from other investigators^{74, 72, 92} indicate that a lower temperature is obtained for foils with cut-off energies in the range 2 keV and lower. Thus, a series of measurements were taken using thin beryllium foils from 15 to 90 μm thickness. The data are displayed in Figure IV-6 where it is seen that the temperature now is about 300 eV. If the temperature were 1 keV or more, then the transmitted x-ray intensities through the foils in this range would be roughly identical. The data shown in Figure IV-6 are averaged from 26 laser shots where the incident energy ranged from 5 to 8 Joules and no detectable scaling of the temperature with laser energy was observed. This is partly due to the fact that below 5 Joules incident laser energy no observable signals are transmitted through the thinnest Be foils. Similar behaviour in x-ray radiation was observed by Buchl⁹² using a neodymium glass laser on a deuterium ice target. A striking feature of the beryllium foil data is that their signal intensities are not radically different from those seen through the thinnest Al foils. This

Figure IV-6 Relative transmitted x-ray intensities from beryllium foils.

suggests that there is a substantial fraction of suprathermal electrons characterized by a 2 keV temperature.

The data from all the foils are summarized in Figure IV-7 where the relative intensities are all plotted versus the cut-off energy of the various foil absorbers. Such a plot of relative transmitted intensities versus the cut-off energy assumes a single temperature Maxwellian distribution of electrons. The data presented in Figure IV-7 are very similar in form to that obtained by Pepin et al.⁷³ and Mitchell et al.⁹³ for CO₂ laser produced plasmas.

The question arises whether line emission from highly ionized nitrogen may cause a serious error in the measurement of the electron temperature through the beryllium foils. Calculations shown explicitly in Appendix A reveal that the emission of the Lyman series of N VI and N VII ions has to be increased by a factor of 10^3 in order for the observed transmitted intensities to give an underestimate of the

Figure IV-7 Summary of x-ray transmission data.

temperature by a factor of two. Therefore, x-ray line emission from nitrogen is considered not to affect the temperature obtained from the beryllium foils.

In summary, the x-ray data suggests that the plasma can be characterized by two temperatures. The lower temperature of 300 eV is measured in the soft x-ray range of 1 - 2 keV while the 2 keV temperature is measured in the range of 3 - 18 keV. The measured intensities through the beryllium foils are mainly due to the soft photons emitted from the colder thermal electrons which should be considerably more abundant than any suprathermal electrons. The aluminum and copper foils transmit only the harder photons from the hotter plasma and thus signal intensities from these foils should be correspondingly smaller than those observed from the beryllium foils. The shape of the transmitted intensities versus cut-off energy thus contains

information of the relative populations of electrons of different temperatures. This topic will be further treated in Chapter V.

The plausibility of the observed electron temperatures can be ascertained by the knowledge of the thermal energy content of the plasma. Thus, several techniques were used to obtain a figure for the energy absorption of the plasma and these results are reported in the following sections of this chapter.

IV-2: Optical Investigation of the Plasma Expansion

This section describes the results obtained from time resolved studies of the plasma expansion using the image converter streak camera described in Chapter III. Observations of the chronology of the luminous front is reported and the results are compared with standard blast wave theory to obtain a figure for the absorbed energy. To investigate the expanding plasma density step without ambiguity, the shadowgraph method is subsequently applied.

An encompassing series of optical investigations were inspired by the feasibility of extracting information on the pertinent physical parameters of the plasma namely temperature, energy absorption and electron density. Comparison of the plasma expansion history with results predicted by the theory of self-similar flow leads to an estimation of the absorbed laser energy in the plasma (Ariga and Sigel⁹⁴, Leonard and Mayer⁹⁵, Basov et al.⁹⁶, Hughes and El-Ferra⁹⁷). Once the absorbed energy is determined and the number of electrons in the plasma are known

by interferometric means, the electron temperature may be calculated through an energy balance (Geunther and Pendleton⁹⁸).

The expansion in the axial direction of a laser-produced plasmas can be characterized, depending upon the experimental conditions, by a laser driven detonation wave¹ with a front velocity proportional to $(J_0/\rho_0)^{0.33}$, where J_0 is the instantaneous absorbed flux density and ρ_0 is the initial gas density. The transverse thermally driven radially symmetric expansion of a laser produced plasma is expected to be describable in terms of a simple blast wave theory which relates the front velocity to the absorbed energy.

Initial observations were time resolved studies of the radial expansion of the luminous emission front of the plasma using the RCA image converter streak camera mentioned in Chapter III. Since this unit has a small optical gain of 4, a fairly wide slit of 2.0 mm had to be used. This restricts the temporal resolution of the streak records to about 1 ns. Preliminary time resolved interferograms and shadowgraphs taken by Hilko for the 110 μ m throat jet revealed a roughly spherical plasma confined within the gas jet boundaries and evolving into an ellipsoid. Accordingly, a 0.5 mm slice of the center of the jet target was imaged onto the streak camera slits. The sweep rate used for the visible emission front histories was 6.9 ns/cm allowing the luminous front to be followed for approximately 50 ns. A representative streak photograph of the luminous front and its accompanying r-t trace is shown in Figure IV-8. All streak photographs taken of the luminous front show that when the plasma first becomes luminous it already has a finite initial

(a)

(b)

Figure IV-8 Luminous front streak photograph
 (a) photograph
 (b) r - t plot for (a)

size of about $400\text{ }\mu\text{m}$. The periodic intensity modulations in the streak photograph are an artifact of the streak camera apparatus and are thought to be due to be some "ringing" noise superimposed upon the applied voltage pulses. The observed initial velocities are constant for the first 10 ns or so and are not strongly correlated with the incident laser energy. The average value for the initial velocity is $6 \times 10^6\text{ cm/sec}$. By 20 ns or earlier after the start of the visible streak, the radial growth of the luminous front slows dramatically, limiting the expansion to 2 mm radius or less. Ahmed et al.⁹⁹ obtained similar streak photographs for ruby laser produced sparks in 4 atmospheres of He gas. The radial growth there also stopped completely after 50 ns of initial breakdown. Incorporation of the effects of ionization into the theory for the temporal radius of cylindrical blast waves resulted in theoretical profiles that also showed the limit to the radial growth.

In the strong shock region of plasma expansion, the well

known blast wave scaling applies (Zel'dovich and Raizer¹⁰⁰)

$$R(t) = A(\gamma) (E_0/\rho_0)^{0.2} t^{0.4},$$

where $R(t)$ is the radius of the blast wave, E_0 is the energy deposited into a point at $R = 0$ and $t = 0$, ρ_0 is the ambient gas density and

$$A(\gamma) = \frac{75}{16\pi} \left[\frac{(\gamma+1)}{(\gamma-1)(3\gamma+1)} \right]^{0.2},$$

where γ is a constant = 1.4, an approximation valid for small Mach numbers. This relationship can be used to obtain space-time curves for various deposited laser energies to which the experimental streak records of the plasma expansion can be compared.

However, most investigations of laser produced plasma expansions have shown better agreement with the theory of Sakurai¹⁰¹. This theory treats the differential hydrodynamic mass, energy and momentum conservation equations and finds solutions in terms of a second order power series in the reciprocal Mach number. This theory is especially applicable for low Mach number shocks $M < 4$. Also, the dimensionality of the blast waves has been considered, spherical blast waves starting from an initial point source, and cylindrical blast waves from a line source. This latter property makes the theory appropriate to laser plasmas formed in gases since they in general, tend to be ellipsoidal in shape. Sakurai gives the general solution for blast waves in the form:

$$(C/U)^2 (R_0/R)^{\alpha+1} = J \left\{ 1 + \lambda (C/U)^2 + 0.5\lambda (C/U)^4 + \dots \right\},$$

where $\alpha = 0, 1, 2$, corresponding to plane, cylindrical and spherical geometry respectively, U is the shock front velocity, C is the speed of sound in the undisturbed fluid, R_0 is a characteristic length related to the energy of the explosion and J_0 and λ are constants. For the 5 Torr N_2 jet the above equation reduces to the convenient form:

$$R(t) = (\gamma/0.88)^{0.25} (E_0/\rho)^{0.25} t^{0.5},$$

where E_0 is the energy deposited per unit length. With the assumption that the laser energy is deposited in a length equal to the depth of focus ≈ 1 mm, the above function was plotted as a function of the deposited laser energy and time in Figure IV-9. For comparison, a spherical blast wave curve is superposed in Figure IV-9. Plotting log-log plots of the r - t traces of the luminous front streak photographs indicates that the time dependence of the radial growth is $t^{0.5}$ rather than $t^{0.4}$. The radius time histories of the plasma for various laser input energies were compared with the theoretical blast wave curves to yield a relationship between the input laser energy and absorbed laser energy. The results are shown in Figure IV-10. The fit of the theoretical curves to the actual time histories is quite rough but the general trend is that the absorbed energy scales linearly with laser energy and the fraction of absorbed energy appears to be about 12%.

Although continuum emission is proportional to the square

Figure IV-9 R-t curves from the Sakurai theory for various absorbed energies.

Figure IV-10 Absorbed laser energies as obtained from the comparison of luminous front streak photographs with Figure IV-9.

of the electron density, the luminous front does not necessarily represent the boundary of the plasma. The blast wave theory applies to the shock front thermally driven by the hot core plasma. Thus, the shock front would be characterized by a region

of high density in comparison to the region of ambient gas ahead of it and plasma behind it. The luminous front may very well be just an ionization front and not be representative of the true shock wave boundary. This is indeed the case for many laser plasmas which are initially bounded by a supersonic heat wave or thermal wave, which later as a consequence of the decrease in energy flux density, change to a subsonic heat wave with the associated leading shock wave.

That the luminous front may misrepresent the shock front of the blast wave has been observed by other investigators (see e.g. Daiber and Thomson¹⁰², Alcock et al.¹⁰³). Alcock et al. performed time resolved studies of laser sparks by schlieren photography, scattered ruby laser light from the plasma, and by luminous emission from the plasma as above. All three methods yield the same initial time dependence of the radial growth, but after the peak of the laser pulse, the luminous front has a time dependence of $t^{0.4}$ as predicted for a spherical blast wave, whereas the shock front as visualized by schlieren time resolved photography has a time dependence of $t^{0.21}$. Daiber and Thomson on the other hand find that the luminous front lagged behind the shock front for their experimental conditions.

Thus, considerable information can be obtained in time resolved shadography of the plasma. The shadowgraph method works on the principle that light passing through a gradient in an index of refraction is deflected to a sufficient extent to make density changes observable. The shadowgraph method measures the second derivative of the refractive index and therefore makes visible only those parts of the medium where the density

gradients through which the light traverses change very rapidly as in the case of shock waves. The arrangement for time resolved shadowgraphy has been shown in the previous chapter. Since the luminosity of the spark gap light source is comparable to that of the plasma, the spark light source is focussed through a small aperture whereas the plasma light is defocussed as implied in the rays of Figure III-1. Such a method has been used successfully by Cheng et al.¹⁰⁴ in the study of exploding wires.

The feasibility of the method was tested by using a freely expanding sonic He jet in air. When the overpressure in this jet nozzle is just slightly above atmospheric, a long steady stream of He gas emanates from the exit. Such a jet in air was observed to produce a comparable fringe shift in an interferometer set-up as the plasma itself. The streak shadowgraph of this jet showed itself as a dark band, so that this method was then known to be useful to visualize the plasma expansion.

Demonstrated graphically by Hilko³¹ is the fact that the size of the shadowgraph images are somewhat dependent on the position of the object plane of the imaging optics. The further the object plane is away from the refracting object, then the larger is the resulting shadow. This is because small ray deflections due to regions of comparatively low electron density become apparent in a much more distant object plane. However, these effects are only prevalent for those phase objects which have a smoothly varying density with radius. For the case of an expanding laser plasma bounded by a dense shock front this problem should not be significant. The object plane for the particular set-up used was 10 cm behind the plasma. The

sensitivity of the arrangement was estimated by geometrical optics considerations³¹ and the indication is that electron densities greater than $0.11 n_{cr}$ should be visible.

A sample streak shadow photograph and its accompanying r-t diagram are shown in Figure IV-11. The two dark bands across the

Figure IV-11 Time-resolved shadowgraphy
 (a) Sample shadow streak photograph
 (b) R-t plot for (a)

streak photograph are taken to represent the boundary of the expanding shock wave. The vertical striations are an artifact of the apparatus and are to be ignored. The fundamental difference between the shadowgraphs and the luminous front streak photographs is that in the shadowgraphs, the plasma boundary is seen to grow from a point. The shadowgraphs also show a rapidly approached limit to the growth of the plasma at about 20 ns. The shadowgraphs indicate a maximum radial extent about 40% smaller than those observed for luminous front streaks for a similar laser input energy.

The comparison of the blast wave curves with the shadowgraph r-t traces reveal the absorbed energy as a function of incident laser energy shown in Figure IV-12. The error bars

Figure IV-12 Absorbed energies obtained by analysis of the shadow streak photographs of the plasma expansion.

come from the judgement of the quality of the fit of the experimental to theoretical curves and measurement error from the streak photographs. The conclusion of this shadowgraph comparison with the blast wave curves is that the energy absorption is linear with incident laser energy but that the energy absorption now is about 5% of the input laser energy.

IV-3: Blast waves with heat transfer

Although the application of standard blast wave theory is relatively common practice in many analyses of laser spark experiments, it is rather limited in its usefulness in providing information on plasma parameters. In addition, its applicability to laser produced plasmas is somewhat questionable for reasons discussed below. Accordingly, a digression is made in this section where the theory of blast waves with heat transfer is introduced. This provides the background for a more useful analysis of the streak photographs of the radial expansion that allows the determination of the electron temperature as well as the absorbed energy.

Comparison of experimental time histories of plasma expansion with blast wave curves must be approached with some caution, since the blast wave theory makes several assumptions which may not pertain at all to the case of laser produced plasmas. Blast wave theory evolves from the self similar flow assumption which requires that the flow variables are all functions of (r,t) thus enabling the solution of the set of hydrodynamic differential equations of energy, mass and momentum conservation. This assumption corresponds to requiring that the profiles of the physical quantities retain their initial form and so remain self-similar to themselves. Secondly, adiabatic flow is assumed, implying that the energy transfer is only due to shock heating and the medium is not supposed to lose, transmit or retain energy by conduction or radiation. Therefore, this assumption limits the applicability of blast wave theory to

those cases where the center temperature is relatively low and much material is initially accelerated. Thirdly, the blast wave is bounded by a shock wave in the usual theory. However, as mentioned previously, many investigations have shown that the luminous front has no resemblance to a shock front initially and does not coincide with the shock front later.

Thus, the transport processes in the gas are neglected, and the specific heat ratio γ and the degree of ionization are assumed to be constant. Nonetheless, self-similar blast wave theory adequately describes the late stages of certain laser spark experiments (Guenther and Pendleton⁹⁸, Buchl et al.¹⁰⁵, Lampis and Brown¹⁰⁶). However, it is not applicable at the space or time origin of a real laser-produced blast wave since it assumes an infinite core temperature and zero core density.

For laser sparks in low pressure gas where a large amount of energy is deposited in a short span of time, the center core has a high initial temperature so that radiative and conductive heat transfer invalidate the above assumptions of self-similar flow blast wave theory. In such cases, the heat transfer dominates resulting in a supersonic heat wave which at a later time in the evolution of the laser spark changes to a subsonic heat wave driving a shock wave ahead of it. This sequence in the evolution of a laser spark is the consequence of a reduction of the energy flux density in the plasma due to the increase of its surface area with time.

The theory for blast waves with heat transfer has been developed by Ahlborn and Ariga¹⁰⁷. In this model, the blast wave expands due to heat transfer described by a surface radiation

intensity. Step wave profiles for all physical quantities by narrow plane fronts are assumed so that integral "jump" conservation equations of mass, momentum and energy apply across all the fronts and determine the type of flow.

A subsonic heat front will be preceded by a compression shock front and a supersonic heat front will be followed by a rarefaction wave. This result is realized from the physical reasoning that for heat waves, the front velocity is given by

$$v_a = W / h_b \rho_a ,$$

where W is the energy flux, h_b is the enthalpy of the region behind the heat wave front and ρ_a is the density of the medium into which the heat wave travels. The velocity of a strong shock front ahead of a heat front is (Ahlborn and Liese¹⁰⁸)

$$v_s = (\text{const } h_b)^{0.5} ,$$

so that no shock can precede the heat wavefront if $v_a > v_s$ or if

$$h_b < \text{const } (W/\rho_a)^{2/3} .$$

This means that once the enthalpy of the plasma has decreased to a certain level due to expansion and a consequent decrease in W , the shock wave will overtake the heat wave front. In the subsonic heat wave's frame of reference, particles enter the heating zone where they are heated and are simultaneously accelerated towards the energy source, leaving the front with

increased velocity as depicted in Figure IV-13. This causes a

Figure IV-13 Regions of a thermal blast wave.

thrust which is responsible for driving a shock wave ahead of the heat wave. The point in the plasma evolution where the shock compression front and the heat expansion front are coincident is known as a Chapman-Jouguet detonation. Therefore, the theory of blast waves with heat transfer models a space-time history as shown in Figure IV-14. Initially, when the energy flux is high compared to the enthalpy h_b there is a supersonic heat wave (i.e., the intake velocity v_a is supersonic with respect to the gas behind it. When the radius has increased to the point where $h_b = \text{const } (W/\rho_a)^{2/3}$, a Chapman-Jouguet detonation momentarily exists where the rarefaction wave head inside the hot core has caught up to the front and the heated particles are leaving the front with the speed of sound. As W decreases further, the subsonic mode of the heat wave then persists and the shock peels

Figure IV-14 Space-time history of a blast wave with heat transfer.

off. The classical blast wave theory is expected to describe the space-time characteristics of this heat wave in the late stages provided that the heat transfer is flux limited. By flux limited transfer it is meant that each particle traveling at its thermal speed $T^{1/2}$ carries its internal energy through the hot surface to provide the flux $W \propto T^{3/2}$.

Many experiments find the time dependence of the radial expansion to be $t^{0.2}$ instead of the standard predicted value of $t^{0.4}$ (see e.g. Alcock et al.¹⁰³). This can be explained by a blast wave that is driven by heat transfer due to blackbody radiation $W \propto T^4$ (Ahlborn¹⁰⁸).

The analysis by Ahlborn and Ariga shows that for the early stages of thermal blast waves, the radius r will vary as

$$r \propto t^{1/(3a-2)}$$

and at later times after the C-J point it should vary as

$$r \propto t^{1/(a+1)},$$

where a is the exponent that characterizes the heat "response" of the plasma i.e., $W = \text{const } h^a$. Thus a streak record of the radial expansion converted into a $\log r - \log t$ plot should reveal regions with two different slopes corresponding to different modes of plasma expansion. The instant of time that the slope changes $t=t_{cj}$, indicates a Chapman-Jouguet detonation. The identification of this point in time in the history of the plasma has a particular significance since the local pressure P_{cj} , sound speed a and energy content E_0 can be obtained from a measurement of the front velocity V_{cj} and radius r_{cj} . At $t=t_{cj}$ the following relations apply:

$$a_b = \left[\frac{g_b P_b}{\rho_b} \right]^{0.5} = \frac{g_b V_{cj}}{g_b + 1},$$

$$P_b = \rho_a V_{cj}^2 / (g_b + 1),$$

$$E_o = \frac{8\pi}{3} \left[\frac{g_b}{g_b - 1} \right] \frac{r_{cj}^3 \rho_a V_{cj}^2}{a_{cj}} = C \rho_a^2 V_{cj}^2 r_{cj}^3.$$

The subscript b denotes the region behind the front, the subscript a denotes the region ahead of the front, ρ is the density and g is the enthalpy coefficient defined by

$$g = \frac{h}{h - P/\rho} = 1 + \frac{P}{3/2P + \sum_{i,i} E_i}$$

Once a $\log r - \log t$ plot of a shadow streak record identifies the Chapman-Jouguet point, the velocity at $t=t_{cj}$, V_{cj} and the radius of the plasma at that time r_{cj} can be used to find the deposited laser energy and pressure of the plasma at this time and consequently the temperature. The streak photographs (such as shown in Figure IV-11), were enlarged sufficiently to minimize the measurement error and $r-t$ and $\log \log$ plots were made an example of which is shown in Fig IV-15. The analyzed streak photographs all showed the same general features indicating a break in the slope at $t = 3$ ns after the initial breakdown. The initial average slope up to 3 ns was 0.82 and after the CJ point the average slope was 0.54, in good agreement with the blast wave theory of Sakurai. The r_{cj} and V_{cj} values for each shot were found by the graphical method indicated and the absorbed energy was calculated using the expression for the C-J point. It should be noted that this calculation is subject to considerable inaccuracy since it contains the square of the velocity and the cube of the radius. In addition, the constant C contains the enthalpy coefficient g . For short time scales where full ionization is not reached even when the shock velocity is high, then $g = \gamma = 1.4$, the adiabatic exponent for a cold gas. If time scales are longer so that the shocked gas will have a time to reach an equilibrium degree of ionization that is consistent with the enthalpy jump

Figure IV-15 Chapman-Jouguet analysis of streak photographs
 (a) r - t plot from a shadowgraph streak
 (b) log-log plot for (a) showing the change in slope

across the shock, then the value for g can be given approximately. Calculations were performed by Ahlborn and Ariga and show that for shock velocities greater than 10^6 cm/sec, the value of the constant C is about 30. It is noted that if full ionization right at the shock front is assumed, then $g = 1.67$ and $C = 9$, whereas for a cold neutral gas $C = 12$. Thus, the greater part of the uncertainty lies in the degree of ionization attained at the shock front. The absorbed energy E_0 as obtained using this analysis is shown in Figure IV-16 taking the optimistic value of 30 for C . The error bars in this graph arise from the uncertainty in the r_{cj} and V_{cj} values which stem from the uncertainty in determining t_{cj} by graphical analysis.

Figure IV-16 Absorbed energy obtained by analysis of shadow streak photographs.

Nonetheless, the indication is that the upper bound for the absorption by this method of analysis is 17% of the incident laser energy.

The pressure determined at the C-J point is a more confident estimate since the expression used only depends upon the factor $1/(g+1)$ so that the uncertainty of the degree of ionization is unimportant compared to the accuracy of the measurements themselves. Of interest here is to calculate the temperature and compare the resulting values with that obtained using the x-ray diagnostics. Assuming the ideal gas law we have:

$$\begin{aligned}
 P &= nkT = n_e kT_e + n_i kT_i \\
 &= \frac{n_e k(Z+1)T}{Z} \quad \text{for } T_e = T_i = T \\
 &\approx n_e kT \quad \text{for } Z > 3.
 \end{aligned}$$

For hydrodynamic motion one usually expects that $T_e = T_i$, however, the electron-ion equilibration time t_{eq} as calculated in chapter II is of the order of 100 ns so that in the short time scales of the laser pulse and certainly up to the time of the C-J point, $T_e > T_i$. Therefore, the electron temperature for the plasma can be obtained from the C-J point velocities by the relation

$$T_e = \frac{P_{cj}}{n_k e B} = \frac{\rho_1 V_{cj}^2}{n_k e B (g + 1)}$$

The analysis does not incorporate the possible effect of hot electrons on the hydrodynamic motion. It is assumed that the hot electron number is small compared to the cold electron number density. The C-J velocities V_{cj} , are determined by the graphical method described previously, ρ_1 is the nitrogen gas jet density into which the blast wave expands and it is known from the isentropic flow relations for a Laval nozzle as shown in Chapter III. The electron density behind the blast wave is measured using a Mach-Zehnder interferometer to be described shortly. The enthalpy coefficient g for nitrogen of a few hundred eV is 1.5 based on LTE calculations¹¹⁰. The results for the measured V_{cj} velocities are shown in Figure IV-17, where the error bars represent the uncertainty in determining t_{cj} . The results thus indicate that temperatures of a few hundred eV are plausible thermal temperatures for the plasma. A scaling of electron temperature with incident laser energy is suggested by

Figure IV-17 Electron temperatures determined from C-J analysis of shadowgraph streaks.

the data of Figure IV-17 although there is too much scatter to make a good quantitative statement for this scaling.

IV-4: Time-resolved interferometric study of plasma expansion

A Mach-Zehnder interferometer, the geometry of which is shown in Figure IV-18 was set-up to obtain a series of time resolved interferograms from which the density profiles of the plasma could be obtained. Interferograms were digitized and Abel unfolded by a computer routine utilized by McIntosh⁸³. Time resolution was provided by using an 80 ps duration ruby laser probe beam that can be arbitrarily delayed with respect to the CO₂ laser pulse. Perhaps the most striking feature of the interferograms is the result that the plasma first forms at the two He-N₂ interfaces. These two initially separate plasma islands quickly merge by the time of the peak of the laser pulse, the front plasma expanding much more rapidly and to

Figure IV-18 Geometry of the Mach-Zehnder interferometer.

greater dimensions than the rear plasma island. The axial length of the plasma by the end of the laser pulse reaches several mm with the front plasma expanding into the ambient He. Typical axial electron densities at the time of the laser pulse are around $2 \times 10^{18} \text{ cm}^{-3}$.

The reason why two plasma islands form is a topic for some speculation. The obvious suggestion is focussing lens spherical aberration effects causing the focal region to have several regions of high intensity near the focal plane (Evans and Grey-Morgan^{'''}). However, laser produced burnmarks on aluminized mylar films in the vicinity of the focal plane showed no indication of such aberration. Other possible explanations may lie in the cascade theory of gas breakdown by laser radiation where the ionization thresholds are critically dependent on the collision frequency for momentum transfer between electrons and atoms. This might imply a Penning effect phenomena i.e., there is a lowering of the breakdown threshold at the He-N₂ boundary.

An interferogram of the plasma taken several nanoseconds

after the start of the laser pulse is shown in Figure IV-19 together with the calculated density profiles from Abel

Figure IV-19 A sample late-time interferogram
 (a) Interferogram at 9.3 ns after the start of the laser.
 (b) Abel unfolded density contours of (a)

inversion. From the Abel unfolded contours the significant result is that an off axis high density annulus forms which stays close to the expanding plasma boundary. This is a clear indication of shock wave formation. Several series of unfolded interferograms were selected and the electron densities at various axial positions were plotted as a function of time as shown in Figure IV-20. The two axial positions used for these plots are the front He-N₂ interface toward the focussing lens and the middle of the gas jet target. This interferometric study reveals that strong shock formation is largely prevalent and restricted to the region of the front He-N₂ boundary, where the

Figure IV-20 Radial plasma density profiles at various times indicated.
 (a) at the front He-N₂ interface
 (b) at the jet center

expansion of the plasma is evidently more rapid than in the middle of the interaction volume.

To confirm the two mode nature of plasma expansion, the outer edge of the plasma obtained from the interferograms was plotted on log-log plots shown in Figure IV-21. Although there is some scatter in the points due to data being taken from different laser shots of slightly varying incident laser energy, the break in the slope indicating a Chapman-Jouguet condition is quite dramatic. The slope after the C-J point is 0.21 as expected for radiative transfer. The application of the same procedure to the middle jet region corroborates the observations of the shadow streak photographs, namely that the slope after the C-J point is 0.4 and the break in the slope occurs at about 3 ns.

Figure IV-21 Log-log r - t plots of the plasma radius.
 (a) at the front He-N₂ boundary region
 (b) at the jet target center

IV-5: C-J Analysis of the Front Plasma Island

The radial expansion of the front He-N₂ interface region was investigated more closely using the Hamamatsu streak camera system. Due to the high gain of this device the spatial resolution could be made as low as 5 μ m of the jet width and the first 3-4 ns of the plasma expansion streaked. In this case, the plasma was backlighted by the 7 ns output from a cavity dumped ruby laser. A sample shadow streak taken of the plasma at the front He-N₂ interface is shown in Figure IV-22. The initial velocities at this interface are higher than those observed at the jet target center, the average being 2.2×10^7 cm/sec.

Schlieren streak photography was also performed using the arrangement described in the previous chapter. The schlieren

Figure IV-22 Shadow streak photograph of the expanding plasma at the front He-N₂ interface.

image was scanned across the streak camera slit by calibrated adjustment of the turning mirror for the probe ruby laser beam. Only at the region of the front He-N₂ interface did a striking two front structure appear consisting of two separating luminous bands that correspond to the expanding cylindrical shock wave density front. The log-log space-time plots of shots taken in this region reveal a sharp C-J break in the slope with the slope being 0.2 after the C-J point as shown in Figure IV-23. C-J wave analysis must be applied somewhat differently for this data since the exact density into which the expansion occurs in the region of the front He-N₂ boundary is uncertain. The electron temperature can still be obtained by the following means.

At the time of the C-J point, the rarefaction wave head inside the hot plasma has caught up to the front and the heated particles are leaving the front with the speed of sound

$$\frac{a_b^2}{a_b} = \frac{v_b^2}{v_b} = g_b \left[\frac{p_b}{\rho_b} \right] = g_b \left[\frac{v_{cj}^2}{g_f + 1} \right] .$$

Figure IV-23 Schlieren streak photography of the front He-N₂ interface region.
 (a) R-t plot of a sample schlieren photo (inset)
 (b) Log-log r-t plot for (a)

The pressure/density term in the above expression can be approximated as:

$$\begin{aligned}
 \frac{2}{a_b} &= g_b \left[\frac{P_b}{\rho_b} \right] = g_b \left\{ \frac{Z n_i kT_e + n_i kT_i}{n_i m_i + n_e m_e} \right\} \\
 &= g_b (Z kT_e + kT_i) / m_i \\
 &= g_b kT_e / 2m_p \quad \text{for } Z > 1,
 \end{aligned}$$

where the last step arises from the fact that the mass of the

ion m_i is $2Zm_p$ for full ionization, m_p being the mass of the proton. The combination of this last expression with the previous one results in a simple expression for the temperature:

$$kT_e = \left(\frac{Zm_p g_b}{(g_b + 1)} \frac{V_{cj}^2}{c} \right)^2.$$

The above equation thus only depends on one experimental parameter V_{cj} and is relatively insensitive to discrepancies in g_b . The temperature as a function of the incident laser energy is plotted in Figure IV-24 using this equation and the V_{cj} velocities obtained from the graphical analysis of the Hamamatsu shadow and schlieren photographs. The average observed electron

Figure IV-24 Electron temperatures obtained by C-J analysis of the streak data for the front He-N₂ interface region.

temperature of 360 ± 80 eV is in good agreement with that obtained for expansion in the middle of the jet target. In summary, it has been shown that an electron temperature of 300-

500 eV is obtained from a C-J wave analysis of streak photographs which is based on the identification of the instant of time the plasma expansion changes its time dependence. Standard blast wave analysis of the streak photographs yields a rather poor fit of the observed space-time histories of the plasma and a low energy absorption. The C-J analysis presented here suggests a higher energy absorption than is suggested by the blast wave analysis and is about 17% of the incident laser energy. Which value of absorption is more acceptable can be decided by measurement of the energy absorption by an Ulbricht integrating sphere, the results from which are reported subsequently.

IV-6: Energy Absorption Measurements

The absorption of incident laser radiation was investigated by time integrated measurements of the scattered, backscattered, transmitted and reflected infrared radiation from the plasma using the experimental arrangement described in Chapter III. The backscattered, transmitted and incident energies are routinely monitored on every laser shot. For absorption measurements, the scattered-refracted laser radiation is measured by the Ulbricht sphere. The signal levels from the photometer's Gen-Tec energy meter were quite strong, ranging from 0.2 - 0.015 volts and free from electrical noise. The scattered energy detected by the Ulbricht sphere was found to be linear as a function of the incident laser energy and is shown in Figure IV-25 where it is seen that half of the laser energy is thus decoupled from the

plasma.

Figure IV-25 Scattered-refracted laser energy measured by the Ulbricht photometer.

The fraction of laser energy that is backscattered and collected back through the focusing optics is measured to scale linearly with incident laser energy and saturates at 10% for 10 Joules incident energy. The transmitted laser energy falls exponentially with increasing laser energy. Subtracting the backscattered energy E_b together with the refracted-scattered E_s and transmitted energy E_t from the incident laser energy E_{in} for each shot produces a figure for the absorbed laser energy E_0

$$E_0 = E_{in} - (E_b + E_s + E_t)$$

The results of the measurements using the above equation are shown in Figure IV-26. The absorbed energy is $(20 \pm 3)\%$, which is in good agreement with average absorption obtained by

C-J analysis of the plasma expansion. The abscissa intercept is

Figure IV-26 Absorbed laser energy determined by Ulbricht sphere.

a consequence of the fact that the breakdown threshold to form a plasma requires about one Joule of laser energy.

Comparison of the above absorption results to those obtained in other laboratories is difficult since relatively little has been published on the absorption of high power CO_2 laser radiation. The experiments of Villeneuve et al.¹¹², Garban-Labaune et al.⁶⁵, Lindman¹¹³, Nishimura⁶⁷ are concerned with the absorption of CO_2 laser radiation on solid targets as a function of the angle of incidence of polarization to verify the presence of resonance absorption. They find typical absorption figures of 30-45% for laser intensities between 5×10^{11} and $3 \times 10^{14} \text{ W/cm}^2$. The amount of radiation scattered into a 4π solid angle is about 60%. In these experiments the scale lengths and electron temperatures are such that inverse bremsstrahlung

absorption accounts for less than 5% of the observed absorption in these cases.

Fabre and Stenz¹¹⁴ investigated an underdense preformed plasma in the intensity range of 10^8 to 10^{11} W/cm² and found the absorption to be anomalously high at 50% at 10^{11} W/cm². Inverse bremsstrahlung absorption becomes less efficient with higher intensities as the electron temperature increases. Offenberger³³ at intensities of 10^{12} W/cm² purports a 90% absorption of which only 40% can be accounted for by inverse bremsstrahlung.

Therefore, the observed absorption of 20% is low compared to other experiments elsewhere, but it is also by the same measure an entirely plausible figure. The issue of how the measured degree of absorption can be accounted for in terms of the electron temperature will be addressed in the following chapter.

Chapter V: Analysis and Discussion

The results presented in the previous chapter will now be discussed and interpreted in order to obtain a coherent self-consistent picture of the laser plasma interaction investigated in this work. By unifying all the results obtained, some degree of confidence can then be attached to the values for the measured plasma parameters. The measurements will be interpreted on the basis of an energy inventory to determine how the laser light is absorbed by the plasma.

V-1: Overview

It is useful to summarize the important results of the measurements reported thusfar. Preliminary x-ray measurements on a 2-Torr jet target in the intensity range of $0.3 - 2 \times 10^{12}$ W/cm² produced low electron temperatures under 200 eV that scaled with incident laser energy as expected by comparison with other similar experiments. However, when the CO₂ laser beam was focussed onto a 5 Torr gas jet target at higher intensities, two electron temperatures were found that were characterized by two different ranges of x-ray emission spectra. A suprathreshold temperature of about 2 keV was found to be independent of the incident laser energy while a low 300 eV temperature can be measured only with thin beryllium foils. The two temperatures as presented by the x-ray data of Figure IV-5 suggest that 300 eV is a thermal temperature which characterizes the majority of the

electrons while the 2 keV temperature is ascribed to a suprathermal component comprising of a minority of the electrons.

Which temperature characterizes the plasma as a bulk temperature is ascertained by analysis of streak photographs of the radial plasma expansion. The Chapman-Jouguet detonation point in the plasma expansion history is identified by the graphical method described in the previous chapter. Relations for the local sound speed and pressure that apply to the C-J point then give the electron temperature of the plasma near the time of the peak of the laser pulse. Although there are different rates of expansion at different points along the focal axis, the average temperatures in the incident laser energy range available are 460 ± 200 eV for the middle of the interaction region and 360 ± 160 eV for the region of the front He-N₂ interface nearest the focussing lens. These are the averages of all the streak photographs analyzed for both regions of the target, and the quoted deviation is the standard mean deviation of 10 or more shots. These two measurements affirm the statement that the bulk temperature of the plasma as characterized by its hydrodynamic motion is of the order of 300 - 500 eV. A plasma of a few hundred eV is also plausible based on the calculations and computer code results presented in Chapter II.

A chronological series of interferograms clearly show the expected change in the temporal behaviour of the plasma expansion as it evolves from the supersonic to the subsonic mode. Density profiles obtained from Abel inversion of the

interferograms show the formation of a density step at the edge of the plasma as expected for a blast wave. The interferometric results provide justification of the application of the C-J wave analysis to the streak photographs.

An estimate of the absorbed energy is also provided by the C-J wave analysis, the fraction of absorbed energy being $(17\% \pm 9)\%$. This figure for the absorption is higher than that estimated by the comparison of the streak photographs with hypothetical blast wave space-time curves obtained from the theory of Sakurai or Taylor and Sedov. However, an Ulbricht photometric sphere used to measure the scattered and refracted radiation from the plasma showed that the amount of laser energy decoupled from the plasma is about 50% of the incident laser energy. Upon calculating the energy inventory by subtracting all infrared losses from the plasma the fraction of absorbed energy is determined to be $(20 \pm 3)\%$, in agreement with the estimate obtained by C-J wave analysis.

V-2: Inverse Bremsstrahlung Absorption

The degree of absorption as deduced from the Ulbricht sphere and streak photographs can be used to determine the electron temperature of the plasma, if it is assumed that the absorption is by inverse bremsstrahlung. The measured degree of absorption η , is then related to the electron temperature through:

$$1 - \exp(-k_{ib} L) = \eta \quad ,$$

where k_{1b} is the temperature dependent inverse bremsstrahlung coefficient (Johnston and Dawson⁵) as given in Chapter II. The plasma is fairly uniform in density by the time of the peak of the laser pulse so that the absorption described by the above equation is that which occurs for a laser beam traversing a uniform plasma of average ion charge Z , average electron density n_e of length L . The density contours from the interferograms series provide information on the path length L . A ray tracing computer routine conceived and implemented by Bernard³⁵ shows that the central rays of the laser beam propagate through typically about 2 mm of plasma at a density of $2 - 3 \times 10^{18} \text{ cm}^{-3}$.

Solving for the electron temperature results in the single expression

$$T_e = \left\{ \frac{9.74 \times 10^{-35} Z n_e^2 L}{1 - \left(\frac{n_e}{n_{cr}} \right)^{0.5} (-\ln(1-\eta))} \right\}^{2/3}.$$

The above expression was evaluated for several interferograms with L and n_e determined by the ray tracing routine assuming full ionization $Z=7$ and the measured absorption of $\eta = 0.2$. The results of this calculation are given in Table V-1. The interferometric data reveal why the absorption is linear with incident energy since the length of the plasma

INTERF. NUMBER	INCIDENT ENERGY (J)	TIME (NS) INTO PULSE	n /n e cr	L(cm)	T (eV) e
36	2.5	2.0	0.25	0.18	250
38	2.5	2.5	0.20	0.29	250
45	2.7	3.3	0.20	0.23	220
39	2.9	1.9	0.30	0.17	320
35	2.9	3.8	0.20	0.22	210
18	3.8	1.5	0.25	0.20	270
40	3.6	2.1	0.25	0.23	300
9	4.6	3.0	0.20	0.57	390
31	6.3	4.3	0.20	0.35	290
16	6.5	1.6	0.20	0.45	610
6	6.5	2.4	0.22	0.53	430
11	7.8	3.2	0.16	0.56	290

Table V-1 Electron temperatures calculated by using inverse bremsstrahlung.

increases with incident laser energy. The increase in length of plasma with higher input laser energy is compensated by a correspondingly higher electron temperature and so the absorption remains linear. The on-axis electron density remains fairly constant throughout the duration of the laser pulse. The average electron temperature calculated in this manner is 280 ± 40 eV, which is in good agreement with the average temperature obtained from the x-ray diagnostics and streak photograph analysis. The above calculation shows that linear inverse bremsstrahlung accounts for the observed electron temperature. However, the laser intensities on the target are in the regime where non-linear absorption can be expected to occur. The ratio of the electron oscillatory to thermal velocity is:

$$\left\{ \frac{eE}{m_e \omega_L} \right\} / \left(\frac{kT_e}{m_e} \right)^{0.5}$$

$$= 6.49 \times 10^{-6} (I/T_e)^{0.5} \quad (I \text{ in W/cm}^2, T_e \text{ in eV})$$

$$= 1 - 3, \text{ for } 300 \text{ eV, } 1 - 10 \text{ Joules}$$

$$= 0.4 - 1.2, \text{ for } 2 \text{ keV, } 1 - 10 \text{ J}$$

Therefore, since the oscillatory velocity is comparable or greater than the thermal velocity throughout the range of laser intensities on target, the absorption should be less than the linear result by the factor

$$\left(1 + \left(\frac{\langle v_{os} \rangle}{3v_{th}} \right)^2 \right)^{-3/2},$$

the physical reasons for this are presented in Chapter II. For the range of vacuum intensities incident on the plasma, the inverse bremsstrahlung absorption coefficient is lowered by a factor ranging from 0.79 to 0.25 which in the temperature calculation shown above results in a lower temperature of about 100 eV at moderate intensities of $3 \times 10^{13} \text{ W/cm}^2$.

A more detailed calculation allows the electron temperature to vary with time and thus the absorption becomes a time varying function. Recent evidence¹¹⁵ of Thomson scattering from excited TPD modes indicates that the temperature at the peak of the pulse is around 600 eV. In this calculation, the temperature was allowed to follow the laser pulse shape temporally with the peak temperature being 600 eV. The resulting calculated absorption was found to have the linear dependence

$$\eta(t) = 0.14t - 0.12,$$

with t in ns. The resulting absorption integrated over the period of the laser pulse is

$$\frac{\int \eta(t) E(t) dt}{\int E(t) dt} = 0.21$$

Therefore, this scenario of a time dependent temperature and absorption fits the measured time integrated absorption provided by the Ulbricht photometer.

Due to the large amount of beam refraction measured it is apparent that the on-target intensity is a function of time and not equal to the vacuum focal intensity of 10^{14} W/cm². This effect may very well be a mechanism that limits the intensity to levels such that collisional absorption remains linear. For a 300 eV plasma, this statement implies that the on-target intensities are $< 10^{13}$ W/cm². However, at the time of the peak of the laser pulse, the dimensions of the plasma observed from the interferograms are on a sub-millimeter scale which are insufficient to cause severe beam refraction. Thus the intensities early in the pulse must be near to 10^{14} W/cm².

The saturation of inverse bremsstrahlung absorption has not yet to date been conclusively confirmed. For example, Gray and Kilkenny¹⁴ incorporated this effect into their one dimensional computer code for simulating a CO₂ laser interaction with an underdense Z-pinch plasma at the modest intensities of 10^{11} W/cm². They found that the absorption had to be increased by factor of 1.1 to get a good fit of the simulation to the experimental results.

Langdon⁹ applies his correction to the absorption in the experiments of Shay et al.¹¹⁶ which involved intensities of 10^{14} W/cm² at a wavelength of 1.06 μ m, and electron temperatures of 400 eV. The calculated reduction in absorption is 40% which is comparable to other refinements invoked in the computer modeling of the experiment. More recently, Herbst et al.¹¹⁷ for similar experimental conditions to those of Shay et al. used a two dimensional cylindrical hydrodynamic code with collisional absorption and classical thermal conduction and found the experimental values of the absorption to be consistent or slightly greater than the calculated values, but no more than by 10% of the incident energy. Thus the presence of saturation of inverse bremsstrahlung at high laser intensities is a topic of some controversy.

Since the bremsstrahlung emissivity and absorption are directly related to one another it is appropriate to consider what the effect of the high intensity non-linearity is on the bremsstrahlung emission. This effect may introduce an error into the determination of electron temperature by the method of x-ray absorption as used in this thesis. This effect has been considered by Pert¹¹⁸. The bremsstrahlung emission cross-section for x-rays is $\sigma \propto v^2$ and the electrons possess a velocity which is the sum of the thermal and oscillatory components $v = v_{os} + v_{th}$. The time averaged emission rate for a Maxwellian distribution of electrons is then $R = \langle 1/(v_{os} + v_{th}) \rangle f(v_{th})$

Provided that the oscillation energy $0.5mv_{os}^2$ is much less than the emitted photon energy $h\nu$ then the resulting emitted intensity is calculated to be augmented by the factor of

$$\left(1 + \frac{(2h\nu/kT_e - 3)m\langle v_{os}^2 \rangle}{6kT_e} \right) .$$

To investigate the significance of this effect, the above expression was incorporated into the calculation of x-ray intensity transmitted through foil filters. The results are shown in Figure V-1 where it is seen that the electron temperature measurements should not be greatly in error.

Figure V-1 Electron quiver velocity effect on the electron temperature measurement. The dashed curves are the x-ray foil transmissions with the incorporation of the effect.

Although the non-linear bremsstrahlung emission is substantially increased over the linear emission, the relative transmitted intensities through the foils is not changed significantly. The reason for this is that the spectral distribution of the x-ray radiation is little changed by the laser field and the foils

transmit a limited wavelength band.

It is of interest to note that although the electron temperature increases due to bremsstrahlung absorption at an energy gain of $0.5m\langle v_{os}^2 \rangle = 1300$ eV per collision, the rapid expansion along with thermal conduction cools the plasma sufficiently quickly so that the actual temperatures are typically no greater than a few hundred eV in laser plasmas.

Therefore, linear inverse bremsstrahlung absorption can account for the electron temperature of 300 eV determined by two independent techniques. Although the vacuum focal intensities suggest that saturation of inverse bremsstrahlung can occur, beam refraction may limit the on-target intensities to the linear regime on a time integrated basis. The plausibility of the electron temperature can be checked further by a detailed energy balance calculation and this follows in the next section.

V-3: Energy Balance and Electron Temperature

The results presented thusfar now enable the energy balance to be calculated. Once the energy balance is determined, the thermal electron temperature of the plasma can also be estimated and compared with the values obtained by the soft x-ray diagnostics and the C-J analysis of the streak photographs.

By means of a detailed energy inventory, the efficiency of energy coupling of the incident laser beam to the plasma it produces can be assessed. The electron temperature can be determined once the thermal energy content of the plasma is known and thus serves as a check of the plausibility of the

results presented thusfar.

Both streak photograph analysis and Ulbricht photometric measurements indicate an energy absorption of 20% of the incident laser energy. This absorbed energy is converted into other forms inside the plasma. Parts of it are used to ionize the plasma, and to do work in expansion of the plasma. Some of the energy appears in the kinetic energy of hot electrons, and in the potential energy of the excited modes that occur in the observed parametric instabilities. Some absorbed energy is also lost in radiation. In this section, each of these energy sinks will be evaluated to determine what remains as the thermal energy content of the plasma. The energy balance is written as:

$$E_o = 0.20E_{in} = \frac{3}{2} k_B T_e N_e + \frac{3}{2} k_B T_i N_i + \sum_i N_i E_i + \int P dV + \sum_{rad} P dt + E_h N_h,$$

where P_{rad} is the radiation power, $E_h N_h$ is the energy contained in the hot electrons, the thermal and hot electron number are represented by N_e , N_h and N_i is the ion number. Measurements of the hot electron emission from the plasma (Meyer et al.⁸⁹) indicate that about 1% of all the electrons in the interaction volume are ejected at the mean energy of 85 keV. The focal volume is $\pi r^2 l = \pi \times 50 \mu m^2 \times 300 \mu m$. The energy contained in the Langmuir wave accelerated electrons is $\pi r^2 l n_e \times 0.01 \times k_B T_e$, $k_B T_e$ being 85 keV and is of the order of a millijoule. Therefore the 85 keV electrons are an insignificant sink for the absorbed laser energy.

The energy expended in ionizing the plasma can be crudely

estimated by assuming that all the electrons present are from fully stripped ions. This will be an upper bound to the ionization energy content of the plasma. To fully strip nitrogen of electrons requires a total of about 1500 eV/ion. The calculations for the various interferogram results are presented in Table V-2.

INTERF. NO.	INPUT ENERGY	TIME INTO PULSE (NS)	TOTAL # ELECTRONS 10^{15}	% ENERGY IN IONIZATION
39	2.9	1.9	0.6	0.7
36	2.5	2.0	1.3	1.7
8	2.5	2.5	2.1	2.9
45	2.7	3.3	3.2	4.0
35	2.9	3.8	3.1	3.6
28	4.2	1.0	0.013	0.01
19	4.2	1.2	0.23	0.02
18	3.8	1.5	0.65	0.58
40	3.6	2.1	1.6	0.02
9	4.6	3.0	5.1	3.8
26	6.5	0.9	0.028	0.02
16	6.5	1.6	1.6	0.84
6	6.5	2.4	5.5	2.8
31	6.3	4.3	6.3	3.4

Figure V-2 Energy contained in ionization using interferometric data.

From the table it is seen that by the end of the laser pulse < 4 ns, the maximum amount of ionization energy the plasma can contain is about 4% of the incident laser energy. At the peak of the laser pulse, there are fewer electrons so that the energy in ionization is small, less than 1%.

The energy consumed in expansion work of the plasma against its ambient surroundings is now estimated as follows. From the C-J wave analysis of the streak photographs, the pressure of the

plasma at the time of the C-J point is known. In addition, the volume of plasma as a function of time is known by contour integration of the interferogram data. Therefore the expansion work is calculated by

$$\int P dV = P_{cj} \Delta V,$$

where ΔV is the volume change of the plasma during the laser pulse. Since streak photographs and interferograms are not taken simultaneously the data for laser shots for similar input laser energies are combined. A typical plasma volume at the time of the C-J point is $2 \times 10^{-3} \text{ cm}^3$, while the pressure at this time is about $1 \times 10^9 \text{ dynes/cm}^2$. This results in expansion work of the order of several hundred millijoules or 3 - 5 % of the incident laser energy.

Radiation powers include the sum of free-free, free-bound and line radiation. The details of this calculation are shown in Appendix B where it is seen that radiation losses can be neglected in the duration of the laser pulse.

Since the electron-ion equilibration time for a plasma is (Spitzer²⁴)

$$\tau_{ei} = \left\{ m_i / m_e \right\}^{1/2} \frac{3}{8\pi} \frac{k_B T_e^{3/2}}{n_e^4 \ln \Lambda}$$

$$= 2.84 \times 10^8 \frac{Z_i T_e^{3/2}}{n_e \ln \Lambda}$$

$$= 300 \text{ ns}, \quad (n_e \text{ in cm}^{-3}, T_e \text{ in eV}),$$

for our plasma conditions, the collisional transfer between

electrons and ions is insufficient to impart an appreciable amount of energy to the ions and so ion heating can be neglected in this estimate. The energy balance can now be written as

$$\begin{aligned}
 E_0 &= 0.20E_{in} = \frac{3}{2} N_e k_B T_e + \frac{3}{2} N_i k_B T_i + \sum_i N_i E_i + \int P dV + E_h N_h \\
 &= \frac{3}{2} N_e k_B T_e + 0.04E_{in} + 0.04E_{in} .
 \end{aligned}$$

The electron temperature can be calculated from

$$T_e(t) = \frac{2}{3} \left\{ \frac{(0.12 + .04)E_{in}(t)}{N_e(t)} \right\} ,$$

where N_e is the total electron number which is obtained from the integration of the Abel unfolded density contours from the inteferograms. It is emphasized in the above equation that the energy input E_{in} and electron number N_e are both functions of time. The laser pulse is actually described by a triangular pulse shape of the form

$$I(t) = I_0 0.83t , \quad \text{for } 0 < t < 1.2 \text{ ns}$$

$$= I_0 \{ 1 - (t-1.2)0.385 \} , \quad \text{for } 1.2 < t < 3.8 \text{ ns} .$$

The fractional energy input at a time t into the laser pulse duration is obtained by integration and normalization of the above expression for the pulse shape. The temperature as determined from the energy balance is calculated and presented in Table V-3 using the interferometric data.

Temperatures calculated in this manner are quite high for

INTERF. SHOT #	TOTAL INPUT ENERGY	TIME (NS) INTO PULSE	TOTAL ENERGY AT TIME t	TOTAL # ELECTRONS $\times 10^{15}$	T (eV) e
39	2.9	1.9	1.8	0.6	1500
36	2.5	2.0	1.7	1.3	760
8	2.5	2.5	2.1	2.1	500
45	2.7	3.3	2.6	3.2	410
28	4.2	1.0	0.92	0.013	3510
19	4.2	1.2	1.33	0.23	2900
18	3.8	1.5	1.80	0.65	1620
40	3.6	2.1	2.50	1.6	910
9	4.6	3.0	4.3	5.1	420
26	6.5	0.9	1.2	0.028	2140
16	6.5	1.6	3.3	1.6	1030
6	6.5	2.4	5.2	5.5	550
31	6.3	4.3	6.3	6.3	580

Table V-3 Electron temperatures calculated from the energy balance.

those interferograms taken at times near the peak of the pulse. This arises because the plasma first appears near the peak of the pulse and there is too much energy for too few electrons. The high temperatures at early times are also a consequence of using time integrated measurements of absorption on a time dependent phenomenon. Nonetheless, shots 45, 35, 9, and 31 all demonstrate that there is sufficient absorption to attain temperatures from 400-500 eV since these are all interferograms taken at the end of the laser pulse. Therefore, the 300 eV temperature measured by the soft x-ray diagnostics as inferred from the streak photographic study is further corroborated by the simple energy balance calculation shown above.

V-4: The Nature of the suprathermal temperature

As shown in the preceding sections, the 300 eV thermal temperature determined by the x-ray diagnostics is corroborated by the temperatures derived from a modified blast wave analysis of the streak photographs and also from an energy balance calculation. However, the x-ray diagnostics also show the conspicuous presence of a 2 keV component in the plasma. In this section, the nature of this suprathermal temperature will be discussed in terms of its possible origin and meaning.

Smoothing Effects of the X-ray Diagnostics

In Chapter III it is mentioned that due to the small temporal and spatial extent of the laser produced plasma, the x-ray diagnostics are both time and spatially integrating measurements. Since the foil filters used sample a wide x-ray spectrum ranging from 1 - 18 keV, the spatial and temporal smoothing of the x-ray detection system has important consequences for the interpretation of the data. The information of where and when the 2 keV electrons occur remains obscured by the smoothing effects. However, several scenarios may be postulated and their validity tested by considering the transmission of x-rays through various foils resulting from a time varying electron temperature and density.

In general terms, any input signal $S_i(t)$ can be considered as a linear combination of delta function impulses (e.g.

Helstrom¹⁹⁾

$$S_i(t) = \int S_i(t') \delta(t-t') dt' .$$

The output of a linear device to a Dirac delta function input is the impulse response $K(t)$. For an arbitrary input signal, the output signal must be a linear combination of delayed input responses

$$\begin{aligned} S_o(t) &= \int S_i(t') K(t-t') dt' \\ &= \int K(\tau) S_i(t-\tau) d\tau . \end{aligned}$$

For simplicity and expediency, the impulse response function is taken to be the Heaviside unit step function

$$\begin{aligned} K(\tau) &= 0 , \quad \tau < 0 \\ &= 1 , \quad 0 < \tau < 10 \text{ ns} , \end{aligned}$$

where the limit of 10 ns is the FWHM of the measured x-ray detector impulse response function. The input signals $S_i(t)$ for which the convolutions were performed were the transmitted intensities of two pairs of foils as functions of a time varying temperature. The transmission of x-rays through any given foil is in general a monotonically increasing function of the electron temperature and is thereby readily amenable to a convenient polynomial fit suitable for subsequent integration for the temperature as a function of time. For these calculations the transmitted intensities through a 15 μm and 30 μm Be foil

pair as well as that for a 30 μm and 55 μm Al foil pair were convoluted with the step function to determine what difference in measured temperature occurs for various time dependent temperatures. These particular foil pairs were selected since the beryllium pair transmits the softer x-ray radiation in the vicinity of 1 - 2 keV while the aluminum pair transmits x-ray radiation in the vicinity of 6 -7 keV. The aluminum foil pair is thus selective in the sense that it transmits only the harder radiation from a comparatively hotter plasma. In addition to the hard radiation, the beryllium foils will detect the softer radiation from cooler parts of the plasma which may persist for a longer period of time than the hotter components.

The simplest possibility of temporal temperature variation is that the plasma is heated to a 2 keV temperature and then slowly cools. Since temporally the cooler plasma predominates, it may be possible this would result in the observed two temperatures. Experiments where time resolved measurements were made on laser heated plasmas (e.g. Gray and Kilkenny¹⁴) show that the temperature temporal profile follows the laser pulse and that after the maximum temperature is reached, the temperature falls exponentially with time. A slow exponential decay of the temperature can be somewhat justified on the basis that the visible emission from the plasma has been observed with a Hamamatsu vacuum photodiode to have an $1/e$ folding time of 50 ns. The temperature temporal profile was assumed to be of the form:

$$T(t) = 1667t e^{-t} \quad 0 < t < 1.2 \text{ ns}$$

$$= 2000 \left\{ \exp - (t-1.2)/\tau \right\} \quad \text{for } t > 1.2 \text{ ns} ,$$

where τ , the $1/e$ folding time, varied from 1 to 20 ns. It is also necessary to consider the time variation of the x-ray emission due to the growth of the emitting volume of plasma, particularly since there is relatively little plasma at the time of the peak of the laser pulse. A study of the interferogram series shows that the average density of the plasma is constant but the volume is a rapidly varying function:

$$\begin{aligned} V(t) &= 0 , \quad \text{for } t < 0.5 \text{ ns} \\ &= At^4 , \quad 0.5 < t < 2.5 \text{ ns} \\ &= A(2.5)^4 , \quad \text{for } t > 2.5 \text{ ns} . \end{aligned}$$

This function for the plasma volume was incorporated into the convolution. The integrations were numerically performed by means of the UBC library routine of DCADRE. The results for the temperature history described are shown in Table V-4.

The results in Table V-4 indicate that the beryllium foils should measure a temperature that is 60% within the peak temperature and the aluminum foils within 70% provided that τ , the fall time is around a few nanoseconds as is expected. Therefore, the temperature measured by the beryllium foil would always be about 1 keV if the plasma peak temperature is 2 keV.

The other possibility is that the suprathermal temperature persists only for a short time in comparison with the laser pulse duration and it characterizes only a fraction of the total electron distribution. This is suggested by the fact that time

τ (NS)	Be FOIL TEMP. (eV)	% OF MAX. Te	Al FOIL TEMP. (eV)	% OF MAX. Te
1	640	32	1240	62
2	780	39	1280	64
5	920	46	1300	65
10	1160	58	1420	71
15	1300	65	1520	74
20	1440	72	1580	79

Table V-4 Results from the convolution calculations with $T_{\max}$
= 2000 eV

resolved Thomson scattering observations (Bernard³⁵) of the excited TPD (two-plasmon decay parametric instability) and SBS modes indicate that these phenomena exist for times 0.7 to 1.5 ns after the start of the laser pulse. This hypothesis was tested by superimposing upon the low temperature history, a hot temperature "spike" whose duration was 0.8 ns as indicated below:

$$\begin{aligned}
 T_e(t) &= 250(t) && \text{for } 0 \leq t \leq 1.2 \text{ ns} \\
 &= 300 \exp\{-(t-1.2)/10\} && \text{for } 0 \leq t \leq 10 \text{ ns} \\
 &= T_e(t) + T_{\text{hot}} && 0.7 \leq t \leq 1.5 \text{ ns} .
 \end{aligned}$$

The contribution of transmitted x-ray intensity from T_{hot} was weighted by the factor

$$\left(\frac{n_{\text{hot}}}{n_{\text{total}}} \right)^2 ,$$

since the continuum emission is proportional to the square of the electron density. The use of a step function for the temporal profile for T_{hot} is justified on the basis that for the plasma densities observed, the mean free path for a kilo-electron volt electron exceeds the dimensions of the plasma.

Several possible values of T_{hot} were considered. Immediately apparent is the effect of the 100 keV electrons that arise as a result of the SRS and TPD instabilities. For the hot electron fraction of 15% the result of 2500 eV was obtained using the Al foil pair. A result of 1500 eV was obtained when the fraction was lowered to 10%. The time-smoothed transmitted intensity through the beryllium foils is unaffected by the short T_{hot} pulse, the temperature obtained being around 270 eV.

Resonance absorption if present, is predicted to produce a hot temperature of around 10 keV (Forslund¹²) for $n_e = 0.25 n_{cr}$ a cold temperature of 300 eV, and laser intensities of 10^{14} W/cm². The calculations for this T_{hot} show that if the hot electron fraction is 30% then the temperature inferred from the aluminum foil transmitted intensities is 1600 eV, while the Be foils as before imply a temperature of 250 eV.

Ion acoustic turbulence can be expected to produce a suprathermal temperature of a few keV maximum. Only by allowing all the electrons to have a 3 keV temperature for the interval 0.7 to 1.5 ns could a temperature of 1600 eV be obtained from the aluminum pair while a temperature of 300 eV is obtained from the beryllium foils.

Finally, a 2.5 keV temperature was allowed to have an

exponential fall time along with the cold temperature to test the effect of a continuous hot T_h component instead of a short spike as described above. A temperature of 1600 eV is obtained from the aluminum foils if the hot electron fraction is 40 % and both hot and cold temperatures have a decay time of 10 ns. However, in this case, the temperature obtained from the beryllium foils is 800 eV. A reduction in the decay time of the hot temperature to 2 ns lowered the beryllium foil temperature to 420 eV but then the aluminum foil measured temperature is only 1 keV.

The effect of spatial smoothing of the x-ray detection system was also considered. The interferograms in general, reveal that the plasma has a constant average density up to a radius of 300 - 200 μm for times up to about 2 ns after the start of the laser pulse. Thus a spatial radial temperature profile was assumed of the form:

$$T_e(r) = T_{\text{max}} - \left[\frac{T_{\text{max}} - T_{\text{min}}}{R_f} \right] r ,$$

where R is 300 μm . This is modeling the plasma by a hot central axial core surrounded by progressively cooler cylindrical shells all of the same electron density. If T_{max} and T_{min} were set to 500 and 100 eV respectively then both foil pairs measure a temperature that is within 10% of the maximum. For the extreme case of a 2 keV core temperature falling off radially to 100 eV at R both foil pairs measure 2 keV to within 40%. Therefore, on the basis of this calculation, the foil temperature measurement

spatially obtains the maximum temperature from a uniform density region provided the temperature gradients are not too steep.

In summary, the observation of a two component electron temperature by the x-ray absorber method can be best explained by the hypothesis that a hot electron fraction persists for a short time interval. The three unknown parameters are the hot electron fraction, the hot temperature and its time duration. The scenarios that yield the experimental measurements all have the bulk of the electrons characterized by a temperature of 300 eV which peaks at the laser pulse and decays exponentially with a fall time of several nanoseconds. This produces the cold temperature measured through the beryllium foils. The hot electron temperature, because it occurs only for a comparatively short time contributes a small background signal to the beryllium foils but a much more significant signal through the aluminum and copper foil absorbers due to their bias toward transmission of the harder x-ray radiation.

There are several combinations of hot electron temperature and fraction that produce a 2 keV temperature measurement from the aluminum foil pair. The presence of 100 keV electrons in the time interval specified produced 2 - 2.5 keV when the fraction of hot electrons was between 10 and 15%. This is however greatly in excess of the fraction of hot electrons (1%) estimated by the experiments described by Meyer et al.^{8,9}

A temperature of 10 keV, which may arise if resonance absorption is present, requires 30% of the electrons to be at this temperature to obtain a measurement of 2 keV. Resonance absorption requires that a critical density layer be present. An

overcritical density can be achieved if the jet target were to be instantaneously fully ionized. However, all interferograms taken indicate an electron density $< 0.4 n_{cr}$, so that resonance absorption as a mechanism for producing the hot electrons is unacceptable.

Ion acoustic turbulence and filamentation (Ng et al.¹²⁰) are known to produce temperatures of a few keV. If a 3 keV temperature is allowed to persist for the 0.8 ns interval, then only by allowing all the electrons to be at this temperature can the experimental time and space integrated x-ray data be duplicated.

As indicated in Chapter II and verified by Offenberger et al.³³, the presence of significant levels of ion acoustic turbulence requires that the absorption of the laser radiation be quite high. With the measured absorption limited to 20%, ion-acoustic turbulence is thereby an unlikely candidate for the origin of the observed suprathermal temperature.

This leaves filamentation as the remaining origin of the suprathermal temperature. Ng et al.¹²⁰ demonstrated experimentally a correlation of the intensity of hard x-rays attributed to filamentation and saturation of the stimulated Brillouin backscatter. In this gas jet experiment, saturation of SBS and SRS were both observed with a peak of SRS activity occurring at about 6 Joules of incident energy. A similar behaviour of the transmission of x-rays through thick aluminum foils was observed. The definitive proof of filamentation would be x-ray pinhole photographs (Fong¹²¹) of the focal volume showing small regions of intense x-ray emission. Unfortunately,

the combination of plasma density and temperature did not allow sufficient x-ray intensities to produce convincing single shot x-ray photographs. In the few instances where pinhole photographs were obtained by several integrated exposures, the x-ray emitting volume was comparable if not greater than the focal volume itself. However, the saturation of the SRS electrons and $3/2 \omega_0$ reflectivity would suggest that wavebreaking of the SRS electron waves provides the source of suprathermal electrons.

Wavebreaking, in the sense used in this thesis, refers to the phenomenon observed in the computer simulations of Forslund et al.¹²². The simulations demonstrated the disruption of the stimulated Brillouin and Raman waves for incident laser powers where $v_{os}/v_{th} > 1$. In these circumstances, the plasma becomes compressed into a wave form by the ponderomotive force pattern of the pump and backward light waves. The density perturbations may be as much as an order of magnitude above the background density. However, the plasma eventually reacts against this squeezing and this reaction manifests itself in the computer simulations as a wave breaking in the electron or ion phase space. At the point of breaking, the particle energy density is equilibrated with the radiation pressure. At wavebreaking, the pressure balance per unit volume is

$$\frac{E_0^2}{8\pi} = n_e kT_e,$$

where E_0 is the electric field of the laser beam. For 5 joules of laser energy delivered in 2 ns into a plasma of density $2.5 \times$

10^{18} cm^{-3} , the above condition yields a $T_e = 2.5 \text{ keV}$.

V-5: Deconvolution of the X-ray Data

Thusfar, the hot electron population, as evidenced by the suprathermal electron temperature, has not been considered in the energy balance. This section attempts to estimate the number of these hot electrons by a deconvolution procedure applied to the x-ray transmission data, and by a simple energy balance calculation.

A common procedure in x-ray diagnostics of laser produced plasmas is to obtain a spectrum of the x-ray emission from the transmission data. The intensity spectrum provides information of the temperature of the plasma from the slope of the fit of the data much the same as the methods used in Chapter IV. Furthermore, the intensity spectrum is used to obtain an estimate of the electron distribution, e.g. Brueckner¹²³. Directly applicable to the data presented in this thesis are the analytical methods presented by M.H.Key¹²⁴, where it is shown that subject to some approximations, knowledge of the continuum emission spectrum in a given photon range as contained in x-ray transmission data allows deconvolution of the electron energy distribution in the same energy range. The reasoning behind the convolution procedure is outlined below.

The spectral intensity per unit frequency interval due to electrons of number density $n_e(v)dv$, in the velocity range v to $v+dv$, interacting with ions of density N_i and charge Z is

$$dI(\nu) = N_i v n_e(v) p(\nu) dv ,$$

where $p(\nu)$ is the cross section for bremsstrahlung emission and is $\propto 1/\nu^2$. The total emission spectrum due to any distribution of electron velocities is then the integral :

$$I(\nu) = AN \int_{E_0}^{\infty} n(E) E^{-1/2} dE ,$$

where the variable change $\nu \rightarrow E$ was made and $E_0 = h\nu$. If $n_e(E)$ decreases rapidly with increasing energy i.e., $n_e(E) \propto \exp(-E/kT)$ then $E^{-1/2} \approx E_0^{-1/2}$ in the integration so that

$$n(E_0) = \frac{E_0}{AhN_i} \cdot \frac{dI(\nu)}{d\nu} ,$$

subject to the condition that $kT/2E_0 \ll 1$. The above equation represents the simplest form of deconvolution. Refinements involve the inclusion of Gaunt factors (Karsas and Latter¹²⁵) to adjust the classical results to quantum mechanically corrected values. However, for the photon energies and electron temperatures considered, the Gaunt factors are slowly varying and for free-free radiation are equal to 1 within 15% throughout the range of parameters considered. Therefore, for expediency the Gaunt factors are set equal to one.

To determine the electron energy distribution, the continuum intensity $I(\nu)$ must be deduced from the transmitted x-ray intensity as a function of absorber thickness $T(t)$. The absorption coefficient of an x-ray absorber is taken as $\mu(\nu) = k\nu^{-3}$ so that the transmitted flux is

$$T(t) = \int I(\nu) \exp(-k t \mu^{-3}) d\nu .$$

The introduction of the cut-off frequency by $\nu_c = (k t)^{1/3}$ and subsequent differentiation of the above expression gives:

$$dT(t)/dt = \nu_c / t \int I(\nu) (\nu / \nu_c)^3 \exp(-(\nu / \nu_c)^3) d(\nu / \nu_c) .$$

Key considers this integral for the case of the spectra encountered in x-ray continuum emission from low atomic number laser produced plasmas. The integrand above shows a maximum close to $\nu = \nu_c$ with a spectral spread of the order of $\Delta\nu/\nu_c = 0.3$. This integral is approximated in terms of the spectral intensity $I(\nu_c)$ in terms of a decaying exponential and the result is that

$$I(\nu_c) = \frac{1}{3} \frac{dT(\nu_c)}{d\nu_c} \cdot \frac{1}{X} ,$$

where X represents the value of the integrand convoluted with the approximation for $I(\nu)$ in terms of $I(\nu_c)$ and is a slowly varying function which to a first approximation may be assumed to be constant. Key applied this method of deconvolution to the experimental results of Buchl et al.¹⁰⁵ who observed a two temperature plasma in the x-ray photon range of 1 - 15 keV. The deconvolution method described is considered to have a 25% accuracy which is sufficient for an estimate of the hot electron fraction as desired here.

X-ray transmission data such as presented in Figure IV-6 generally show monotonically decreasing functions with increasing cut-off energy. However, the experimental data

obtained indicate that $dT/dE \rightarrow 0$ for the energy range 5 - 10 keV. This is a nonphysical result since $dT/dE \propto E/kT$ i.e., the slope of the x-ray spectrum is inversely proportional to the electron temperature. Accordingly, the best fit hyperbolic curve through the data was used as the starting point for the deconvolution procedure. The accuracy of the fit may be judged from Figure V-1. The hyperbolic fit to the data was subsequently

Figure V-2 Deconvolution of the x-ray transmission data.
 (a) Curve fit (dashed) of the data.
 (b) Intensity Spectrum obtained from (a)

differentiated and the result being in the second graph of Figure V-1. Differentiation of the intensity spectrum curve multiplied by $E^{1/2}$ leads to the electron distribution curve shown in Figure V-2. Maxwellian energy distributions of the form:

$$f(E) = A E^{0.5} \exp(-E/kT) ,$$

were fitted to the distribution with T_e being 300 eV and 2000 eV. The ratio of the areas underneath the two Maxwellian curves then gives the hot electron fraction as being $(12 \pm 5)\%$.

Figure V-3 Electron energy distribution from x-ray data.

Another estimate of the hot electron population is obtained by reexamination of the energy balance. As previously calculated in Table V-3 there is sufficient energy absorption of the input laser energy to heat all the electrons in the plasma to several hundred eV. Assuming the thermal plasma temperature is 300 eV and the suprathermal temperature is 2000 eV, the fraction of hot electrons can be simply estimated by:

$$E_{th} = (0.12 \pm .03)E_{in} = \frac{3}{2} N_{ec} k_B T_{ec} + \frac{3}{2} N_{eH} k_B T_{eH} .$$

The total number of electrons is

$$N_{ec} + N_{eH} = N_{Total}$$

where N_{Total} is measured by interferometry. The hot electron fraction was calculated by the solution of the previous two equations using data from late time interferograms and the results are presented in Table V-5. The average hot electron

INTERF. SHOT #	TOTAL INPUT ENERGY	TIME INTO PULSE	TOTAL ENERGY AT TIME t	TOTAL ELECTRON # $\times 10^{15}$	HOT ELECTRON FRACTION
36	2.5	2.0	1.7	1.3	0.21
8	2.5	2.5	2.1	2.1	0.12
45	2.7	3.3	2.6	3.2	0.063
35	2.9	3.8	2.9	3.1	0.10
40	3.6	2.1	2.5	1.6	0.28
9	4.6	3.0	4.3	5.1	0.11
6	6.5	2.4	5.2	5.5	0.10
31	6.3	4.3	6.3	6.3	0.12

Table V-5 Hot electron fraction calculated by energy balance.

fraction calculated in this manner appears to be $(12 \pm 3)\%$ in agreement with the estimate obtained by the deconvolution procedure described previously. It should be noted that the scenario to which these estimates apply are those where the hot electron fraction is always present throughout the laser irradiation of the plasma.

V-6: Energy Contained in the Plasma Waves

The previous calculations involving the energy balance do not consider the energy transferred to the ion and electron waves that arise from the SBS and SRS parametric instabilities. This additional possible energy sink will be considered in this section.

Computer simulations of SBS by C.E. Max¹²⁶ indicate the creation of ion heated "tails" in the Maxwellian distribution. Physically, the energy source for ion heating is from the momentum balance as expressed by the Manley-Rowe relations

$$\frac{E_L^2}{\omega_L} = \frac{E_W^2}{\omega_W},$$

where E_L , ω_L and E_W and ω_W are the electric fields and frequencies of the laser and plasma wave respectively. Since energy is proportional to E^2 , the fraction of incident light energy given to the plasma wave is ω_W/ω_L . For the case of SBS this fraction is

$$\frac{\omega_W}{\omega_L} = \frac{2c_s}{c},$$

here c_s is the ion sound speed velocity and c is the speed of light. The assumption of an electron temperature of 300 eV for the calculation of the ion sound speed shows that only 0.04% of the incident laser light can be contained in the ion acoustic wave.

The same reasoning as above leads to 50% of the laser

energy being absorbed into the electron plasma waves at $0.25 n_{cr}$. However, as demonstrated by the theoretical work of Coffey¹²⁷, electron plasma waves saturate when the electrostatic field of the wave becomes large enough to trap and accelerate electrons. The parameter of relevance is $\beta = 3v_{th}/v_{ph}$ where v_{th} is the thermal velocity of the electrons and v_{ph} is the phase velocity of the electron plasma waves which for Raman scattering at $0.25 n_{cr}$ is $c/\sqrt{3}$ (Hiob and Barnard¹²⁸). The result from Coffey's work is the saturation of the electric field of the electron plasma wave by the factor:

$$\frac{E'}{W} \rightarrow \frac{E}{W} \left(1 - \frac{1}{3} \beta - \frac{8}{3} \beta^{0.25} + 2\beta^{0.5} \right)^{0.5}.$$

For a sub-kiloelectron-volt plasma, the reduction in the electric field of the electron plasma wave due to electron trapping is about 50%. This means that by the Manley-Rowe relations, the maximum fraction of laser energy that can be converted into the electron wave energy is $(0.5)^2 \times 0.5$ or about 13% of the input pump energy. From the energy balance considerations presented earlier in this chapter, it is obvious that the energy contained in the electron plasma waves can only be at most, a few percent of the incident laser energy.

Chapter VI: Conclusions

The interaction of a CO₂ laser beam with an underdense plasma has been investigated in this work. The detailed energy balance of the laser-plasma interaction was obtained in order to explain the observed plasma electron temperatures. The diagnostics used by the author are soft x-ray absorption, streak photography, Ulbricht sphere photometry and supplemented by time resolved ruby laser interferometry. The information obtained from these diagnostics was used to determine the energy coupling of the laser beam to the plasma it produces in a low pressure nitrogen gas jet target.

Electron temperatures of a few hundred eV were obtained, consistent with numerical estimates based on inverse bremsstrahlung absorption and thermal conduction. The observed electron temperatures do not indicate a saturation of collisional absorption as is theoretically predicted for focal intensities where $v_{os}/v_{th} > 1$. An energy inventory of the infrared radiation leads to the figure of absorption of the incident laser energy of 20% which is accountable by collisional absorption. Of this fraction, a few percent has been calculated to go into ionization of the plasma and expansion work. The small number of electron plasma wave accelerated electrons constitute a negligible energy sink. The energy inventory allows an estimate of the thermal energy content of the plasma which upon assuming that all the observed electrons can be characterized by the same temperature, yields temperature estimates that account for heating to several hundred eV.

Shadow and schlieren techniques were used in the time resolved studies of the radial plasma expansion. These investigations were used initially to determine the energy deposition by comparison of the space-time chronologies of the plasma radius with that predicted by the conventional blast wave theory. This method of analysis of the streak records showed that the fractional absorption of the input energy is a comparatively small 5% to 12% depending on the particular optical method used. Closer inspection of the time dependence of the radial growth of the plasma indicates that there is a sudden change during the laser irradiation. This is suggestive of the two mode expansion observed in other laser spark experiments and electrothermal shock tubes. Ascribing Chapman-Jouguet wave detonation conditions at the time of change of the temporal behaviour of the plasma expansion allows the straightforward calculation of the electron temperature and absorbed energy subject to a few justifiable assumptions. The electron temperature calculated by this method yields several hundred eV in good agreement with that obtained by the x-ray diagnostics. The absorbed energy is found to be about 20% using the C-J wave analysis. The thermal temperatures measured and calculated in the course of this work are all summarized in the plot of Figure VI-1.

The optical study also revealed the complex structure of the laser produced plasma. The plasma initially consists of separated high density islands which have distinctly different growth rates due to expansion in different ambient conditions. This behaviour was verified by time resolved interferometry

Figure VI-1 Summary of measured and calculated electron temperatures.

which nonetheless indicated the unmistakable presence of the two mode temporal expansion of the plasma. Particularly striking is the observation of the radial propagation of shock fronts by a schlieren technique which showed a time dependence of $t^{0.2}$. This is indicative of radiative heat transfer driving the expansion.

The application of the Chapman-Jouguet wave conditions to the analysis of streak records is a novel technique used successfully to obtain the electron temperature and absorbed laser energy. The agreement of the calculated values with experimentally measured quantities illustrates the usefulness of this method to determine the parameters of a plasma from streak photographs alone. The temperatures obtained from this analysis also establish the x-ray diagnostics measured temperature of 300 eV as being the thermal temperature of the plasma.

The x-ray diagnostics also measure consistently a 2 keV

electron temperature in a harder x-ray energy range than in which the 300 eV was observed. Accounting for this temperature in terms of an energy balance is not straightforward since the x-ray and energy absorption measurements are time integrated. This leads to the postulation of several scenarios of time dependent hot electron temperatures attributed to various fractions of the electron distribution being comprised of hot electrons. Convolution calculations were performed to determine what plasma conditions would simulate the obtained results from the x-ray diagnostics. From the many possible combinations of hot electron temperature and fractional density, one plausible combination is that the plasma is at 3 keV for a short interval of 0.8 ns during which SRS and TPD modes have been observed to exist by time resolved Thomson scattering. This sudden temperature jump is superimposed upon a temperature history that follows the laser pulse and has a peak temperature of a few hundred eV.

Possible mechanisms for this 2 keV suprathermal temperature are ion acoustic turbulence, filamentation and wavebreaking of the excited electron plasma waves. The simple balance between radiation and plasma pressure that is thought to occur in wavebreaking does lead to a calculated electron temperature of 2.5 keV.

If the hot electron temperature were present throughout the laser-plasma interaction, then simple energy balance calculations indicate that about 12% of the electrons are at 2 keV the remaining being at 300 eV. A graphical deconvolution procedure of the x-ray data also suggests that 12% of the

electrons can be ascribed to be suprathermal.

The ambiguity that arises from the spatial and temporal smoothing of the x-ray diagnostics does not allow a conclusive statement to be made concerning the nature of the observed suprathermal temperature. However, several reasonable speculations can be proffered as illustrated above.

In conclusion, linear inverse bremsstrahlung absorption is shown to account for the observed thermal temperature. The observations are shown to be self-consistent and are corroborated by the analysis of streak photographs and from the energy balance determined by integrating sphere photometry. The suprathermal electron distribution determined by the x-ray diagnostics is ascribed to be due to the wavebreaking of the observed parametric instabilities.

Further Work

Considerable more insight may be gained about the laser plasma interaction if time resolved x-ray diagnostics were to be made available. The time integrated measurements although useful in providing some information cannot give a detailed account of the chronology of the electron temperature. For example, the experiments of Lee and Rosen¹²⁹ using an x-ray streak camera have shown that while the thermal temperature of the plasma from a laser irradiated gold disk target remains fairly constant throughout the laser pulse, the hot electron temperature has a temporal behaviour which follows the laser pulse intensity.

A problem with CO₂ laser plasmas is that because of their

low density, the x-ray fluence is several orders of magnitude smaller than that from Nd laser plasmas thus high gain x-ray measurement devices are necessary. Commercial x-ray streak cameras with microchannel plate amplification have become recently available. The use of such a camera may provide temporal resolution on the evolution of the electron temperature.

As this work is part of an overall program dedicated to the study of parametric instabilities, knowledge of the spatial and temporal distribution of the thermal electron temperature is essential for the calculation of instability thresholds for the various parametric processes. Time resolved ruby laser Thomson scattering can be in principle be implemented for this purpose. The inherent difficulty here is the fact that only a factor of 10^{-12} of the incident light is scattered and elimination of stray light from the collection optics becomes a major technical problem.

A possible circumvention to the problem of the low levels of Thomson scattered light from thermal fluctuations is to observe the scattering from the enhanced density fluctuations arising from the pumped instabilities. Electron plasma waves are subject to the dispersion relation:

$$\omega_e^2 = \omega_p^2 + \frac{3}{2} k^2 v_{th}^2,$$

where ω_e and ω_p are the electron wave and plasma frequency respectively. Two plasmon decay modes occur at the density of $0.25 n_{cr}$ so that ω_p is known. The scattered spectrum should show

a peak at a frequency corresponding to the electron wave frequency, by the frequency matching condition:

$$\omega_L = \omega_e + \omega_{\text{scat}}.$$

By the measurement of the scattered spectrum and subsequent identification of ω_{scat} the electron temperature may then be calculated from the dispersion relation. Once again, time resolved spectroscopy is desirable to discriminate the scattered ruby laser spectrum from nitrogen lines from recombining ions.

The remaining questions of the presence of beam filamentation or ion acoustic turbulence also remain topics that warrant further investigation. The unambiguous finding of the filamentation instability would require state-of-the-art x-ray imaging optics coupled to a high gain picosecond resolution x-ray streak camera. The investigation of ion turbulence however, can be done with the laboratory resources presently available by using an appropriate Thomson scattering geometry to probe submicron wavelength fluctuations.

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Appendix A

Nitrogen line emission effect on T_e Determination

For a nitrogen plasma of a few hundred eV temperature, the dominant radiation is line radiation from the highly ionized ions. This exceeds the recombination radiation and bremsstrahlung radiation combined by as much as four orders of magnitude (Jacobs et al.¹³⁰) at 10 eV. It is appropriate to consider the effect of the line radiation on the x-ray measurements particularly when thin foils of low cut-off energy are being used. The comparatively soft line radiation may be appreciably transmitted through these foils causing a background signal which upon taking ratios of transmitted x-ray intensities may lead to an erroneous figure for the temperature.

The calculations performed to consider this effect of the line emission are presented in this appendix. For simplicity, the steady state corona model is used for the ion populations. In a laser produced plasma, the population densities of various ion species is a time dependent problem and requires the solution of a system of $Z+1$ coupled rate equations (Magill²⁹). Here, the electron density, temperature and hence collisional ionization coefficients are all time dependent quantities of the electron temperature and density determined by the physics of the laser plasma interaction. Nonetheless, the steady state corona model is an often used simplification applied to laser plasmas (see e.g. Pepin et al.⁷³, Donaldson et al.⁷²). As shown by Galanti and Peacock¹³¹, the steady state corona model

overestimates the population densities of the higher charge states but much less so than the LTE Saha equilibrium model.

The line emission, bremsstrahlung and recombination continua and their transmitted intensities through various foils are calculated using the computer program of Fong¹² with minor modifications as follows. The absorption properties of the scintillator have been taken into account by convoluting the bremsstrahlung and recombination emissivities with the factor $(1 - \exp(-\mu_{sc}(\lambda)x))$ x being the thickness of the scintillator in units of mg/cm² and $\mu_{sc}(\lambda)$ is the wavelength dependent mass absorption coefficient of the scintillator. The analytic expressions for the mass absorption coefficients of the scintillator and the beryllium foils had to be incorporated into the program. These expressions are:

for the scintillator:

$$\mu_{sc}(\lambda) x_{sc} = 0.929 \lambda^{2.88} + 0.120 \lambda^{2.38},$$

for beryllium foils:

$$\mu(\lambda) = 0.269 \lambda^{3.063},$$

for copper foils:

$$\mu(\lambda) = 118 \lambda^{2.68},$$

where λ is in Angstroms.

An expression for the intensities of the line emission in the x-ray region was obtained and summed to the value of the continua intensity in the wavelength interval of 0.1 - 50 Å.

In general terms, the intensity of a line transition $p \rightarrow q$

of an ion with charge Z is (McWhirter¹³²):

$$\epsilon_{\lambda} = hc/4\pi\lambda A(p,q) N(z,p) ,$$

where $A(p,q)$ is the transition probability and $N(z,p)$ is the population of an ion of charge Z in state p . In the corona model, the population densities of the excited levels are determined by a balance between the rate of collisional excitation from the ground level balanced by the rate of spontaneous radiative decay:

$$N_e N(Z,g) X(T_e, g, p) = N(Z,p) \sum_{q < p} A(p,q) ,$$

where $X(p,q)$ is the collisional excitation rate of level q to level p . In the calculations, only the transitions to the ground level are considered to be sufficiently intense so that the sum over q disappears. Therefore, the line intensities are given by:

$$\epsilon_{\lambda} = hc/4\pi\lambda N_e N(Z,g) X(T_e, g, p) ,$$

and

$$X(T_e, g, p) = \left[\frac{6.5 \times 10^{-4}}{E(p,g) T_e^{1/2}} \right] f(g,p) \exp(-E(p,g)/kT_e) ,$$

where $E(p,g)$ is the excitation potential from the ground state (in eV), $f(g,p)$ is the oscillator strength.

In units appropriate to the incorporation of the expression into the computer program, the expression for the line intensities is written as:

$$\epsilon_{\lambda} = 6.38 \times 10^8 \left[\frac{f(g,p) N(Z)}{\lambda(A) E_p T_e} \right] \exp(-E_p / kT_e) .$$

The relevant parameters (Allen¹³³) are given in the Table where only the first few Lyman series lines are considered:

N VI			N VII		
WAVELENGTH	E (eV) p	f	WAVELENGTH	E (eV) p	f
23.771	521.55	0.052	19.826	625.3	0.0290
24.898	497.94	0.144	20.910	592.9	0.0791
28.787	430.67	0.674	24.781	500.3	0.4162

Table A-1 Excitation energies and oscillator strengths for N ion lines.

The contributions of the line intensities to the total radiated intensity as a function of the electron temperature is shown in Figure A-1.

Since the shortest wavelength of the line series is 19 Å the transmitted intensity even through the thinnest foils is not a significant effect. Calculations for the case of x-ray transmission through beryllium foils are presented in Figure A-3 where it is evident that the line intensity must be increased by a factor of 10^3 for the electron temperature measurement to be in error by 50%.

Figure A-1 Line emission intensity as a function of T_e (eV).

Figure A-2 Change in relative transmitted x-ray intensities due to x-ray line emission.

Appendix B

Calculation of Radiative Power Loss from the Plasma

The radiation from the plasma is radiated as bremsstrahlung, recombination and line radiation. In this appendix, the radiative power loss is calculated explicitly for a plasma characterized by a 300 eV electron temperature. The equations used here may be found in Puell⁶⁰. Bremsstrahlung arises from the deceleration of electrons in the Coulomb field of ions, and photons are emitted corresponding to the change in energy. The total bremsstrahlung power P_{ff} emitted per unit volume is

$$P_{ff} = 1.5 \times 10^{-32} T_e^{1/2} n_e \sum_Z n_Z Z^2, \quad (\text{W/cm}^3)$$

where the Gaunt factors are taken to be one and T_e is in eV. The sum is over the charge states. A steady state coronal equilibrium calculation shows that at 300 eV, 91% of the N atoms are fully ionized, the remaining being N VII ions.

Recombination radiation arises from free electrons being captured by ions. For the nitrogen plasma, recombination radiation exceeds bremsstrahlung (Appendix A). In the frequency range $h\nu > X$ where X is the ionization potential of an ion with charge $Z-1$ for electrons in the n th shell, recombination radiation and bremsstrahlung show the same spectral behaviour. At $h\nu = X$ the recombination radiation spectrum shows a discontinuity, since this case corresponds to the capture of a

resting free electron by an ion.

The total power emitted per unit volume as recombination radiation is

$$P_{fb} = 2 \times 10^{-31} n_e T_e^{1/2} \sum_Z n_i(Z) \sum_n \left(\frac{x}{x_H} \right)^2 \frac{\zeta_n}{n},$$

where T_e is in eV, $n_i(Z)$ is the density of ions of charge state Z , n is the principal quantum number and ζ_n is the number of unoccupied sites in the n th shell. The relevant parameters are given in Table B-1

$Z+1$	NVIII	NVII	NVI	NV	NIV	NIII	NII
Z	7	6	5	4	3	2	1
x	667	552	97.9	77.5	47.4	29.6	14.5
n	1	1	2	2	2	2	2
ζ	2	1	8	7	6	5	4

Table B-1 Relevant quantities for recombination radiation from a nitrogen plasma.

When there are ions with one or more bound electrons present in the plasma, the excitation of these electrons and the corresponding line radiation must be considered. Shown in Appendix A is that the line radiation exceeds the bremsstrahlung and recombination radiation combined for a plasma of a few hundred eV temperature. As in Appendix A, only the contribution of Lyman series lines of the N VI and N VII are considered as being sufficiently intense. The power emitted as line radiation

per unit volume is

$$P_1 = 3.5 \times 10^{-25} T_e^{-1/2} n_e \sum_Z n_i(Z) \exp(-E(Z)/kT_e),$$

where the excitation energies are given in Appendix A.

Assuming an electron density of $2 \times 10^{18} \text{ cm}^{-3}$ and a temperature of 300 eV, the radiation power loss from the plasma is

$$P_{ff} + P_{fb} + P_l = (7 \times 10^6 + 2 \times 10^7 + 5 \times 10^8) \text{ W/cm}^3$$

The volume of the plasma at the end of the laser pulse is typically $4 \times 10^{-3} \text{ cm}^3$ so that the power loss through radiation is 2×10^6 Watts. In contrast, the input laser power is typically 5×10^9 Watts. Therefore, the radiation power from the plasma is up to 3 orders of magnitude smaller than the laser input power. This has important consequences in the Ulbricht sphere measurements and in the calculation of the energy balance as described in the text.