Investigation of Mercury's Magnetospheric and Surface Magnetic Fields

by

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Abstract

This thesis is devoted to the study of Mercury's magnetic field environment, to reveal the nature of the interaction between a weak planetary magnetic field and the interplanetary medium. Due to the lack of orbital spacecraft observations at Mercury prior to the MErcury Surface, Space Environment, GEochemistry, and Ranging (MESSENGER) mission, work in this thesis presents some of the first analysis and interpretation of observations in this unique and dynamic environment.

The bow shock and magnetopause define the boundary regions of the planet's magnetosphere, thereby representing the initial interaction of the planetary field with the solar wind. We established the time-averaged shapes and locations of these boundaries, and investigated their response to the solar wind and interplanetary magnetic field (IMF). We found that the solar wind parameters exert the dominant influence on the boundaries; we thus derived parameterized model shapes for the magnetopause and bow shock with solar wind ram pressure and Alfvén Mach number, respectively.

The cusp region is where solar wind plasma can gain access to the magnetosphere, and in Mercury's unique case, the surface. As such, this area is expected to experience higher than average space weathering and be a source for the exosphere. Using magnetic field observations, we mapped the northern cusp's latitudinal and longitudinal extent, average plasma pressure and observed its variation with the solar wind and IMF. From the derived plasma pressure estimates we calculated the flux of plasma to the surface.

Mercury's internal dipole field is not centered on the planet's geographic equator but has a significant northward offset. We developed the technique of proton-reflection magnetometry to acquire the first measurements of Mercury's surface field strength. Proton loss cones are evident in both the north-

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ern and southern hemispheres, providing confirmation of persistent proton precipitation to the surface in these regions. We used the size of the loss cones to estimate the surface magnetic field strength, which confirm the offset dipole structure of the planetary field. With additional proton-reflection magnetometry observations, we generated a global proton flux map to Mercury's surface and searched for regional-scale surface magnetic fields in the northern hemisphere.

Preface

This thesis is based on four papers: three have been published and one is in preparation for publication. Consequently, some background information is repeated in the introductory section of each chapter.

A version of Chapter 2 has been published. I identified all bow shock and magnetopause crossing times analyzed in the paper, which were checked by Dr. Brian Anderson. I also performed all the calculations and analyses, made the plots and wrote the manuscript. Dr. Catherine Johnson and Dr. Brian Anderson provided guidance throughout and all co-authors commented on the manuscript. In the Journal of Geophysical Research paper associated with this chapter there was an error in the production of Figure 13; the figure and the associated text have been corrected in Chapter 2.

A version of Chapter 3 has been published. I identified all the cusp crossing times in the residuals to the magnetic field data, after the magnetospheric model was removed; residual magnetic field files were produced by Dr. Johnson. I conducted all the analyses, performed all the calculations, made all the plots and wrote the manuscript. I received comments on the manuscript from all the co-authors. Dr. Johnson and Dr. Anderson provided guidance throughout.

A version of Chapter 4 has been published. I developed the idea of using proton-reflection magnetometry, and adapted the technique of electron reflectometry to protons. Normalization files for the proper analysis of FIPS observations was provided by Dr. Dan Gershman and Dr. Jim Raines. I performed all the analyses and calculations, and received guidance throughout from Dr. Johnson and Dr. Anderson. I also wrote the manuscript which received comments from all the co-authors.

A version of Chapter 5 is in preparation for publication in a major space

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physics journal. The usefulness and possibility of deriving a global proton flux map to Mercury's surface and establishing regional-scale surface fields with additional proton-reflection magnetometry observations was noticed independently by Dr. Johnson and I. I developed the methods and techniques to perform the analyses, made all the calculations and wrote up the results. I am currently drafting the journal manuscript based on Chapter 5 of the thesis.

Journal Papers

Chapter 2:

R. M. Winslow, B. J. Anderson, C. L. Johnson, J. A. Slavin, H. Korth, M. E. Purucker, Daniel N. Baker, S. C. Solomon (2013), Mercury's magnetopause and bow shock from MESSENGER Magnetometer observations, *Journal of Geophysical Research: Space Physics*, 118, 2213-2227.

Chapter 3:

R. M. Winslow, C. L. Johnson, B. J. Anderson, H. Korth, J. A. Slavin, M.
E. Purucker, S. C. Solomon (2012), Observations of Mercury's northern cusp region with MESSENGER's Magnetometer, *Geophysical Research Letters*, 39, L08112.

Chapter 4:

R. M. Winslow, C. L. Johnson, B. J. Anderson, D. J. Gershman, J. M. Raines, R. J. Lillis, H. Korth, J. A. Slavin, S. C. Solomon, T. H. Zurbuchen, M. T. Zuber (2014), Mercury's surface magnetic field determined from proton-reflection magnetometry, *Geophysical Research Letters*, 41, 4463-4470.

Chapter 5:

R. M. Winslow, C. L. Johnson, B. J. Anderson, D. J. Gershman, J. M. Raines, R. J. Lillis, H. Korth, S. C. Solomon (2014), Regional-scale surface magnetic fields and proton fluxes to Mercury's surface determined from proton-reflection magnetometry, *in preparation*.

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List of Abbreviations

- EPPS Energetic Particle and Plasma Spectrometer
- ER Electron Reflectometry
- ESA Electrostatic Analyzer
- FIPS Fast Imaging Plasma Spectrometer
- FOV Field of View
- IMF Interplanetary Magnetic Field
- MAD Median Absolute Deviation
- MAG Magnetometer
- MBF Mercury Body Fixed
- MCP Micro-Channel Plates
- MESSENGER MErcury Surface, Space Environment, GEochemistry, and Ranging
- MHD Magnetohydrodynamic
- MSO Mercury Solar Orbital
- PAD Pitch Angle Distribution
- RMS Root-Mean-Square
- SEA Superposed Epoch Analysis
- UTC Coordinated Universal Time

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Chapter 1

Introduction

Mercury, because of its close proximity to the Sun, has long eluded direct in situ observations due to the technical challenges of observing, navigating, and operating so close to the Sun. Prior to 2008, our only information of Mercury came from Earth-based radio astronomy [1–4], and three Mariner 10 flybys of the planet between 1974 and 1975 [5–9].

The first and third Mariner 10 flybys made the surprising discovery that Mercury has a primarily dipolar intrinsic magnetic field, with a surface magnetic field strength of ~ 1% of Earth's dipole field [5, 9]. Such a predominantly dipolar internal field may imply an active dynamo process, which was unexpected for Mercury due to predictions of early solidification of the core (e.g., [10, 11]). The weak strength of the field was also thought to be difficult to explain if Mercury's dynamo is driven by thermo-chemical convection, as is the case for Earth's dynamo (see review in [12]). Remanent crustal magnetic fields were not originally favored for Mercury since crustal fields are usually dominated by small-scale structure. However, Aharonson et al. [13] showed that spatial variations in solar insolation on Mercury could give rise to long-wavelength variations in the depth to the Curie isotherm for the dominant magnetization carrier, that in turn would allow long-wavelength structure in the crustal field.

From the limited Mariner 10 flyby observations alone it was not possible to determine whether the dominant source of Mercury's magnetic field was of core or crustal origin and more observations were needed to characterize the spatial structure of Mercury's magnetic field, crucial for determining the source(s) of the field. More recently, Earth-based radar observations have provided support for a dynamo field because the measured amplitude of Mercury's forced librations suggests decoupling of the core and mantle, favoring the existence of a liquid outer core [4].

Our knowledge and understanding of the innermost planet has significantly improved in the last few years because of the MErcury Surface, Space Environment, GEochemistry, and Ranging (MESSENGER) mission, which on the 18th of March 2011 became the first spacecraft to enter orbit around Mercury. MESSENGER first reached Mercury's orbital distances in 2008, made three subsequent flybys of the planet, and was then inserted into orbit with a 200 km periapsis altitude, 82.5° inclination, 15,300 km apoapsis altitude, and 12-h period (reduced to 8-h in March 2012). The highly elliptical orbit of the spacecraft is due to the absorption and re-emission of solar radiation from Mercury's surface, which if endured for too long would increase spacecraft temperatures above safe operational limits.

MESSENGER has seven scientific instruments and a radio science experiment to probe Mercury's surface, its tenuous exosphere (a volume surrounding the planet containing neutral atoms and ions that are collisionless due to their low number density), and its magnetic field environment. Recent MESSENGER results have confirmed that Mercury's intrinsic magnetic field is a weak, global, dynamo generated dominantly dipolar field [14-16]. They have also revealed that the dipole field is aligned with the rotation axis but is offset northward from the planetary equator by 0.196 $R_{\rm M}$ (where $R_{\rm M}$ is Mercury's radius) [17–19]. Mathematically, such a northward offset of the dipole means that in a spherical harmonic expansion of the field, even though the dipole term dominates, there is a significant quadrupole term, with a ratio of 0.4 for the quadrupole to dipole terms. A number of core structures and dynamo regimes have been suggested for Mercury that can reproduce its weak dipole moment (see discussion in [18]). However, dynamo models still have difficulty reproducing all aspects of Mercury's magnetic field in a time-averaged sense, particularly the field's high axial alignment, the large value for the ratio of the quadrupole to dipole term, and the low upper bound on the ratio between the octupole and dipole field [18]. Research on this topic has shown that there is promise in models that invoke a non-conductive layer above a deep dynamo that preferentially attenuates the highly time-varying higher-degree components of the field [20, 21].





Figure 1.1: Schematic diagram of Mercury's magnetosphere. From Zurbuchen et al. (2011), Science, 333, 1862-1865 [22]. Reprinted with permission from AAAS.

A global dipolar planetary magnetic field is capable of dominating the space environment in its near vicinity by ordering and directing charged particles and thereby partially shielding the region called the magnetosphere from the solar wind. The magnetosphere is a magnetic cavity that forms due to the interaction of the solar wind with the dipole field and is described in more detail in Sections 1.1 and 1.4. Figure 1.1 shows a schematic diagram of Mercury's magnetosphere; here we briefly introduce aspects most relevant to this thesis and define them in more detail in the coming sections. Surrounding the planet and its magnetic field are the bow shock, magnetosheath, and magnetopause. The bow shock is a shock wave that forms as the solar wind transitions from supersonic to subsonic speeds as it encounters the obstacle of the planet's magnetic field. The magnetosheath, just downstream of the bow shock, is the region in which the shocked solar wind, compressed and heated at the shock, flows around the magnetopause. The

magnetopause marks the boundary of the magnetosphere, and is a current layer that acts to confine the planet's magnetic field inside the magnetosphere. Detailed studies on Mercury's magnetopause and bow shock will be presented in Chapter 2. At high latitudes are the magnetic cusp regions (one in the north, studied in detail in Chapter 3, and one in the south), where the magnetopause currents nearly cancel the planet's internal field, leaving behind a weakened field region where solar wind plasma can gain access to the magnetosphere. The plasma sheet in the nightside magnetosphere is also a region of high plasma density, whose spatial extent largely coincides with the tail current sheet, and separates the magnetotail's north and south lobes. The plasma populations found in the cusp region and in the plasma sheet make our proton reflection magnetometry observations in Chapters 4 and 5 possible.

MESSENGER has revealed that the interaction of Mercury's magnetic field with the solar wind is unique in our solar system. The combination of the weak planetary field and strong solar wind and interplanetary magnetic field (IMF) generates a small, highly dynamic magnetosphere around Mercury. At Mercury's orbital distances from the Sun, the solar wind density is on average an order of magnitude higher than at Earth, and the interplanetary magnetic field is a factor of five higher than at Earth [23]. Because Mercury's orbit is eccentric (0.31 AU < r < 0.47 AU, where r is heliocentric distance), the planet is also subjected to different solar wind conditions along its orbit; both the solar wind density and pressure increase significantly near perihelion. The north-south component of the IMF, which has a large influence on planetary magnetospheres (see Section 1.1), has been shown to change on a timescales of a few minutes at Mercury [23]. Variations in Mercury's magnetospheric conditions thus occur on timescales of a few minutes to a Mercury year (88 Earth days) [24].

Mercury's magnetosphere is strongly coupled to the exosphere and the planet's surface. Due to the lack of atmosphere and ionosphere (ionized upper atmosphere) the high-energy solar wind plasma can gain access to the magnetosphere in regions where the shielding is incomplete and bombards the surface of the planet. Ions sputtered from the surface move under the influence of magnetospheric electric and magnetic fields; some ions will be lost through reconnection down-tail or through collision with the magnetopause, while others will be returned back to the surface. This coupling between the magnetosphere – exosphere – surface influences the dynamics within the system.

This thesis is devoted to shedding light on the interaction of the interplanetary medium with this magnetospheric system, particularly as it pertains to three specific regions of the magnetosphere. We also show that the nature of the interaction between the magnetosphere and the solar wind can reveal information about the intrinsic planetary field at the surface and can be used to make indirect surface field measurements. The initial interaction of the solar wind with the magnetosphere occurs at the magnetospheric boundaries, the bow shock and the magnetopause. In Chapter 2 we discuss how we characterized these boundaries through observations and empirical models. In the magnetic cusp regions, shielding of the magnetosphere from the solar wind is incomplete, and so charged particles can gain access to the magnetosphere and precipitate down to the surface. We investigated this solar wind – magnetosphere – surface interaction in the northern cusp region of Mercury, and we describe this in detail in Chapter 3. Finally, in Chapters 4 and 5, we develop a novel method to remotely sense the surface planetary magnetic fields. The motion of solar wind protons, which have gained access to the magnetosphere near the cusp and the magnetopause boundary, allows for the measurement of the magnetic field strength at the surface of Mercury through the technique of proton reflection magnetometry. We describe this technique and show that this method confirms the offset dipole structure at Mercury's surface and particle precipitation to the surface.

1.1 Interaction of the solar wind with Mercury's magnetosphere

The interaction of the solar wind with a magnetized planet produces a magnetic cavity around the planet, the magnetosphere, which confines the planetary magnetic field and is largely devoid of solar wind plasma. For details on the formation of the magnetosphere, see Section 1.4.

Although the interaction between the solar wind and all magnetized planets produces a magnetosphere, the nature of this interaction (i.e. characteristics of the magnetopause current layer and of the shock wave, or bow shock, in front of the magnetopause as well as the rate of exchange of plasma between the solar wind and magnetosphere through reconnection) varies in our solar system. Solar wind conditions vary significantly with heliocentric distance; the solar wind density decreases as $1/r^2$ with distance from the Sun on average, the magnetic field strength also decreases with distance with a more complicated scaling law as do the solar wind electron and proton temperatures (see Table 1.1).

At Mercury, the dominant magnetic field component of the interplanetary magnetic field (IMF) is parallel or anti-parallel to the +x direction in a Mercury solar orbital (MSO) coordinate system, for which +x is sunward, z is normal to the orbital plane and positive northward, and +y completes the right-handed system. The angle of the IMF (also referred to as the Parker spiral angle) with respect to the radial (-x) direction is only $\sim 20^{\circ}$ for typical solar wind conditions at 0.3–0.5 AU. This is about a factor of 2 smaller than the Parker spiral angle of 45° at the Earth. The dominance of the IMF B_x component on average at Mercury has a significant effect on the magnetosphere, which will be discussed further in Chapter 3.

Other solar wind parameters that vary with heliocentric distance are the solar wind Mach numbers, which increase with distance from the Sun. The Mach number is defined as the solar wind flow speed divided by the speed of a fundamental wave mode (e.g. sound waves, Alfvén waves). The solar wind speed is approximately constant with heliocentric distance, and if the speed of the wave mode decreases with heliocentric distance, the Mach number increases correspondingly. The Alfvén Mach number, M_A , will be used in Chapter 2 of this thesis; Alfvén waves are magnetohydrodynamic waves, and have an associated speed of $v_A = \frac{B}{\sqrt{\mu o \rho}}$, where μ_0 is the permeability of free space, and ρ is the plasma density. For example, M_A is expected to be ~ 5 at Mercury and ~ 13 at Saturn [25]. Such a large difference in

			D	T	T
	v_{sw}	n_{sw}	B_{IMF}	T_p	T_e
Scaling	r^0	r^{-2}	$r^{-1}(2r^{-2}+2)^{1/2}$	$r^{-2/3}$	$r^{-1/3}$

Table 1.1: Scaling laws of solar wind speed, v_{sw} , number density, n_{sw} , interplanetary magnetic field magnitude, B_{IMF} , proton temperature, T_p , and electron temperature, T_e , with distance from the Sun [26].

 M_A implies that the bow shock in front (upstream) of Mercury's magnetosphere is much weaker than the equivalent bow shock in front of Saturn's magnetosphere [25].

The merging of the magnetic field lines of the planet's intrinsic magnetic field with those of the draped IMF, through a process called reconnection, is responsible for the circulation of flux in the magnetosphere (see Figure 1.2) and some fraction of the injection of solar wind plasma into the magnetosphere. This circulation of the magnetic flux and associated plasma from the dayside to the nightside and back, called the Dungey cycle [27], is also dependent on heliocentric distance and is much more dynamic at Mercury than at any other planet in our solar system. The large-scale circulation of magnetic flux and plasma in the magnetosphere is driven by the solar wind at Mercury and Earth. However at for example, Jupiter and Saturn the solar wind is much weaker and the magnetospheres of these planets are influenced to a much higher degree by the planetary rotation and the large amount of plasma originating within the magnetosphere. At Mercury, the Dungey cycle time is ~ 2 minutes [28, 29], whereas this time scale is much longer at Earth, ~ 60 minutes.

Reconnection is expected to proceed most efficiently when the IMF direction is oppositely directed to the planetary field in the subsolar region; on the dayside at Earth and Mercury this occurs when the IMF has a southward component, i.e., a negative B_z component. The rate of reconnection has also been shown to depend on M_A and the plasma β , the ratio of the plasma thermal pressure to the magnetic pressure [31]. Recent observations [32] suggest that a low value of M_A on the solar wind side of the magnetopause current layer, and a low β on either side of the current layer as well as a low value of the absolute difference in plasma β across the current layer



1.1. Interaction of the solar wind with Mercury's magnetosphere

Figure 1.2: Magnetospheric convection due to reconnection. The circulation of flux proceeds from numbers 1 through 8 in the figure. At 1, the planet's closed field lines merge with the southward directed, open field lines of the solar wind at the subsolar point where the fields are in opposite direction. The newly opened field lines are convected through from the dayside to the nightside by the flow of the shocked solar wind in the magnetosheath. At 7, the field lines are stretched out far enough that they are in opposite direction again and they reconnect in the tail, to close the magnetic flux circulation. Newly connected planetary field lines in the tail move planetward due to magnetic tension, while completely open IMF field lines move away from the planet. Figure from Basic Space Plasma Physics, Wolfgang Baumjohann and Rudolf A. Treumann, Copyright 1997, Imperial College Press [30]; used with permission.

 $(|\Delta\beta|)$, promotes reconnection onset because under these conditions reconnection is possible over a wider range of shear angles (i.e., angle between the magnetic field direction in the IMF and just inside the magnetopause). Furthermore, Swisdak et al. [33] showed that a diamagnetic drift, which is produced when a pressure gradient is present across the magnetopause

1.2. Charged particle motion

current sheet, can disrupt the reconnection process when this drift velocity is super-Alfvénic (i.e. the solar wind flow speed is greater than the Alfvén speed). Thus reconnection is more likely to be suppressed for cases of high M_A , high β , and high $|\Delta\beta|$. The Alfvén Mach number scales approximately as r with heliocentric distance, and the magnetosheath β has been observed to increase with heliocentric distance [31]; they are thus expected to be smallest at Mercury compared to other planets in the solar system. Work by Dibraccio et al. [34] has shown that as expected, reconnection rates are very high at Mercury, ~ 3 times higher than at Earth, and reconnection occurs over a much broader range of IMF orientations than at Earth. All this implies that Mercury's magnetosphere is more strongly driven by the solar wind and IMF than Earth's magnetosphere because of the interaction of the weak planetary dynamo field with the strong solar wind environment that it is embedded in. The dynamic nature of Mercury's magnetosphere is further explored in Chapters 2 and 3.

1.2 Charged particle motion

Basic insight into the motion of charged particles in electric and magnetic fields must be gained in order to lay the foundations for much of the work that is presented in this thesis. Charged particle motion is well documented in a variety of textbooks in electricity and magnetism as well as in plasma physics. We thus only present results here for scenarios most relevant to the work in this thesis, and for more details the interested reader is referred to texts such as [30, 35, 36]. We briefly describe here the motion of a single charged particle in specified magnetic and electric fields as well as the motion of charged particles trapped in planetary magnetospheres. Section 1.2.1 describes the gyrational motion of charged particles in uniform magnetic fields which applies to work presented throughout this thesis. In Sections 1.2.2 and 1.2.3 we describe the drift motion of charged particles under uniform electric fields as well as under changing magnetic fields, which pertain to results presented in Chapters 4 and 5 of this thesis (drift motions are responsible for populating closed magnetic field lines on the nightside with plasma from the plasma sheet). Finally, in Section 1.2.4 we outline the motion of charged particles in convergent (i.e. increasing) magnetic fields; this motion of charged particles is exploited by the proton reflection magnetometry technique in Chapters 4 and 5.

In most plasmas the electric field and the magnetic field depend on the positions and velocities of all the charged particles in the system, and are thus complicated functions. Magnetohydrodynamic (MHD) equations have been derived to deal with the plasma macroscopically as a fluid, instead of solving the kinetic equations of motion for all the particles individually. However, to gain physical insight into magnetospheric processes, it is instructive to consider the motion of a single charged particle in electric and magnetic fields that are independent of time and that are simple functions of position.

To determine the trajectory of a single particle in a force field as a function of position and time we solve the equation of motion

$$m\frac{d\mathbf{v}}{dt} = \mathbf{F}(\mathbf{x}, t), \tag{1.1}$$

where in the case of a charged particle, one of the forces acting on the particle is the Lorentz force given by

$$\mathbf{F} = q(\mathbf{E} + \mathbf{v} \times \mathbf{B}),\tag{1.2}$$

here q is the electric charge, **E** is the electric field, **v** is the particle's velocity, and **B** is the magnetic field.

1.2.1 Uniform magnetic field $(\mathbf{E} = 0)$

In a uniform magnetic field ($\mathbf{E} = 0$), the equation of motion simplifies to

$$m\frac{d\mathbf{v}}{dt} = q(\mathbf{v} \times \mathbf{B}). \tag{1.3}$$

The particle will have a motion both parallel and perpendicular to the magnetic field, and so it is useful to separate the velocity into its components

parallel and perpendicular to the magnetic field, $\mathbf{v} = \mathbf{v}_{\parallel} + \mathbf{v}_{\perp}$. By taking the dot product of equation (1.3) with the unit vector in the direction of \mathbf{B} we find that the equation of motion in the parallel direction is $m\frac{d\mathbf{v}}{dt} =$ 0. The solution to this is clearly $\mathbf{v}_{\parallel} = \text{constant}$, i.e. the particle has a constant velocity along the magnetic field line direction. By subtracting the equation of motion in the field aligned direction from equation (1.3) and using $\mathbf{v}_{\parallel} \times \mathbf{B} = 0$, we have

$$m\frac{d\mathbf{v}_{\perp}}{dt} = m\mathbf{v}_{\perp} \times \mathbf{\Omega},\tag{1.4}$$

where $\mathbf{\Omega} = qB/m \ \hat{\mathbf{b}}$ is the gyrofrequency, and $\hat{\mathbf{b}}$ is the magnetic field direction. Without any loss of generality, if we let the magnetic field be in the $\hat{\mathbf{z}}$ direction, then \mathbf{v}_{\perp} will lie in the x - y plane. Then the solution to equation (1.4) (showing only the real parts of the complex variables) is

$$v_{\rm x} = v_0 \cos(|\Omega|t + \delta),$$

$$v_{\rm y} = \mp v_0 \sin(|\Omega|t + \delta).$$
(1.5)

 v_0 is the magnitude of the perpendicular velocity and is independent of time, and the \mp sign corresponds to positive and negative particles, respectively. This equation shows that the particle follows a circular trajectory in $v_{\rm x} - v_{\rm y}$ space, with positively charged particles gyrating left while negative ones gyrate in a right-handed direction. The gyroperiod, or the time it takes the particle to complete a cycle of its circular motion is

$$T = \frac{2\pi}{|\Omega|} = 2\pi \frac{m}{|q|B} \tag{1.6}$$

The radius of the circle that the particle travels in is the gyroradius and is given by

$$r_g = \frac{v_\perp}{|\Omega|} = \frac{mv_\perp}{|q|B}.\tag{1.7}$$

The gyroradius of the particle represents a natural length scale in the system; particles with small gyroradii relative to the gradient length scale of the magnetic field, $B/\nabla B$, are highly affected by the ambient magnetic field and thus exhibit motions dictated by changes in the field. However, if the gradient in the magnetic field is high compared to the field strength, particles will have large gyroradii relative to $B/\nabla B$ and are referred to as unmagnetized, because their motion is largely unaffected by the magnetic field and changes in it.

As we saw before, the particle has a constant velocity in the field aligned direction, and the combination of this motion with the circular motion in the perpendicular direction yields a helical trajectory for the particle about the field line. The pitch angle of the helix, or the angle between the magnetic field direction and the particle's velocity is defined as

$$\alpha = \tan^{-1} \left(\frac{v_{\perp}}{v_{\parallel}} \right). \tag{1.8}$$

and will be used extensively in Chapters 4 and 5 of this thesis.

1.2.2 Uniform electric and magnetic fields

The motion of charged particles in the presence of both a uniform electric and magnetic field can be decomposed into a gyrating part about the field line and a uniform drift velocity part. This drift motion is perpendicular to the electric and magnetic fields and is given by

$$\mathbf{v}_E = \frac{\mathbf{E} \times \mathbf{B}}{B^2}.\tag{1.9}$$

This $\mathbf{E} \times \mathbf{B}$ drift is independent of the sign of the particle's charge, and both electrons and ions will drift in the same direction. This drift velocity essentially amounts to the gyroradius varying along the particle's trajectory, increasing it in the direction of \mathbf{E} and decreasing it in a direction opposite to \mathbf{E} .

1.2.3 Nonuniform magnetic field

We now consider the motion of charged particles in a magnetic field that is a function of position and with no applied electric field. In such a timeindependent field the kinetic energy of the particle remains constant, and the motion of the particle can be again decomposed into gyromotion plus a drift motion.

First, we briefly discuss gradient **B** drift. Suppose the magnetic field is in one direction only, $\mathbf{B}(\mathbf{x}) = B(\mathbf{x})\hat{\mathbf{z}}$ and that it has a gradient, ∇B , in the perpendicular direction to **B** along the *y* direction. The gyroradius of a particle in such a field configuration will be large wherever *B* is small, and will be small wherever *B* is large. Thus the drift motion of a positive charge will be in opposite direction to the drift of a negative charge; therefore a current will arise under a magnetic field gradient. The expression for the grad-*B* drift velocity is given by

$$\mathbf{v}_{\nabla B} = \pm \frac{1}{2} v_{\perp} r_g \frac{\mathbf{B} \times \nabla B}{B^2},\tag{1.10}$$

where \pm refers to positive and negative particles, respectively.

We will next consider particles moving along a curved magnetic field line. These particles will experience an outward centrifugal force in the frame of reference moving with the particle's parallel velocity. This force is given by

$$\mathbf{F}_{\mathbf{c}} = \frac{mv_{\parallel}^2}{R_c} \hat{\mathbf{R}}_{\mathbf{c}},\tag{1.11}$$

where $\mathbf{R}_{\mathbf{c}}$ is the curvature vector; it has a magnitude equal to the local radius of curvature of the field line and is pointing radially outward. This centrifugal force gives rise to a drift motion, called the curvature drift

$$\mathbf{v}_{cB} = \pm \frac{v_{\parallel}^2}{|\Omega| R_c} \mathbf{\hat{R}_c} \times \mathbf{\hat{b}}.$$
 (1.12)

A curved magnetic field cannot have a constant magnitude in a vacuum (otherwise $\nabla \times \mathbf{B} = \mathbf{0}$ would not hold), and thus the complete drift velocity of a particle in a nonuniform magnetic field is $\mathbf{v}_B = \mathbf{v}_{\nabla B} + \mathbf{v}_{cB}$. For a southward directed dipole field and prograde planetary rotation, positive ions undergo gradient and curvature drift opposite to the planet's rotation (westward), while electrons drift in the direction of the planet's rotation (eastward).

1.2.4 Motion of trapped particles: magnetic mirrors and bottles

Magnetic fields in the near-Mercury environment are non-uniform both spatially and temporally, but even in such fields, the proton motion is described by the first adiabatic invariant, the magnetic moment, μ_m of the particle. In general, the magnetic moment is given by the electrical current I in a current loop multiplied by the area A of the loop: $\mu_m = IA$. In the case of a single charged particle, the current is the charge q divided by the gyroperiod T. The gyroradius of the particle defines a circular area and so the magnetic moment is given by

$$\mu_m = \frac{q}{T}A = \frac{q}{T}\pi r_g^2 = \frac{\frac{1}{2}mv_{\perp}^2}{B}.$$
(1.13)

The magnetic moment is invariant for slow changes in a system (such that the process is adiabatic) and is thus very useful in interpreting certain elements of charged particle motion. In particular, it is useful in interpreting particle motion in magnetic mirrors and bottles.

A converging magnetic field is needed for magnetic mirroring to occur. We start with a magnetic field increasing in the z direction in a cylindrical coordinate system (Figure 1.3). In this scenario, the charged particle will move in the z direction with a parallel velocity, and as long as changes in the electric and magnetic field occur on length scales larger than the gyroradius, the magnetic moment of the particle will be conserved. We can see from equation (1.13) that if the magnetic field strength increases, then the perpendicular velocity of the particle must increase to keep μ_m constant. Thus the new perpendicular velocity of the particle will be given by

$$v_{\perp}^2 = \left(\frac{B}{B_0}\right) v_{\perp 0}^2,\tag{1.14}$$

where the symbol "0" indicates initial value. In a static magnetic field configuration the total energy of the particle must remain constant

$$v_{\perp}^{2} + v_{\parallel}^{2} = v_{\perp 0}^{2} + v_{\parallel 0}^{2}.$$
 (1.15)

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Figure 1.3: Schematic diagram of converging magnetic fields towards the planet's surface.

Thus as the magnetic field strength increases more and more so does v_{\perp} . There will be a point along the field line when all of the particle's velocity is in the perpendicular direction and thus it will not be able to penetrate further along the field line but will be reflected, or mirrored back. We can express equation (1.14) in terms of the particle's pitch angle, since $v_{\perp} = v \sin \alpha$. At the point of reflection, the particle's pitch angle is 90°, and so can rewrite equation (1.14) in terms of the particle's initial pitch angle and the ratio of the initial magnetic field to that at the reflection point,

$$\sin^2 \alpha_0 = \frac{B_0}{B_m}.\tag{1.16}$$

A particle with an initial pitch angle of α_0 will reflect at a point where the magnetic field is B_m , whereas a particle with a pitch angle less than α_0 (i.e. more field aligned) will reflect after reaching a magnetic field magnitude greater than B_m . The particle's trajectory upward is a mirror image of its trajectory downwards, and it returns to the same point with a pitch angle $180^\circ - \alpha_0$.

The physical explanation for the decrease in v_{\parallel} is that there will be a

component of the Lorentz force which is directed opposite to the gradient in B and decelerates the particles. The component of the Lorentz force in the z direction is

$$\mathbf{F}_{\parallel} = q(\mathbf{v} \times \mathbf{B})_{\mathbf{z}} = qv_{\perp}B_{\rho}, \qquad (1.17)$$

where the field B in the scenario depicted in Figure 1.3 is given by

$$\mathbf{B} = B_{\rho}\hat{\rho} + B_{z}\hat{\mathbf{z}}.\tag{1.18}$$

The field is primarily in the z direction $(B_{\rho} \ll B_z)$ and the component of the field in the ρ direction can be found from Maxwell's equation:

$$\nabla \cdot \mathbf{B} = \frac{1}{\rho} \frac{d}{d\rho} (\rho B_{\rho}) + \frac{dB_{z}}{dz} = 0.$$
(1.19)

In the case where $B_{\rho} \ll B_z$ and near the z-axis ($\rho = 0$) we have that

$$B_{\rho} \approx -\frac{1}{2}\rho \frac{dB_{\rm z}}{dz}.$$
(1.20)

This can be substituted into equation (1.17) to yield

$$\mathbf{F}_{\parallel} = -\frac{1}{2}qv_{\perp}r_g\frac{dB_{\rm z}}{dz},\tag{1.21}$$

where we are considering the case when $\rho = r_g$. This component of the Lorentz force can be rewritten in terms of the magnetic moment

$$\mathbf{F} = -\mu_m \nabla_{\parallel} B. \tag{1.22}$$

This force is thus clearly directed opposite to the gradient in the magnetic field, along $\hat{\mathbf{z}}$, and since μ_m is constant, the only variable in this relation is the magnetic field gradient. Therefore, this force will be higher as the particle moves toward larger values of the parallel magnetic field gradient, and will thus decelerate the particle.

In a magnetic bottle the situation is similar, except the magnetic field is convergent on both sides, i.e. it consists of two magnetic mirrors. In such a situation the particle bounces back and forth between the two convergent ends of the magnetic field, and is trapped. The maximum field strength of the bottle determines the minimum initial pitch angle required for confinement, i.e. whether a particle with a given initial pitch angle, α_0 will be trapped or not. The criterion for confinement in the bottle is given by

$$\sin \alpha_0 > \left(\frac{B_0}{B_{max}}\right)^{\frac{1}{2}},\tag{1.23}$$

where B_{max} is the maximum magnetic field strength reached in the bottle. Thus if initially there is an isotropic distribution of particles (i.e. equal particle fluxes in all directions), then some time later all particles that do not satisfy the confinement criterion above will be missing from the distribution, while the ones that do satisfy the condition will be trapped inside the bottle. The missing portion of the distribution is called the loss cone, which is directly relevant for Chapters 4 and 5 of this thesis.

Such a bottle magnetic field configuration exists on closed magnetic field lines of dipole fields, for example at Earth and Mercury as well. Particles trapped on closed magnetic field lines will bounce back and forth with a bounce period T_b given by the integral over a bounce cycle:

$$T_{b} = \oint \frac{dz}{v_{\parallel}(z)} = 4 \int_{z_{0}}^{z_{max}} \frac{dz}{v \left[1 - \frac{B(z)}{B_{0}} \sin^{2} \alpha_{0}\right]^{\frac{1}{2}}},$$
(1.24)

where v is the total particle speed, z_0 is the particle's starting position in the bottle while z_{max} is furthest the particle travels along the z direction. The bounce period for protons (with an average proton energy of 1 keV) in Mercury's inner dayside magnetosphere is ~ 30 seconds. Particles will also have a drift motion around Mercury due to the grad-B and curvature drift. The drift period of protons of the same energy is ~ 40 minutes. However, because Mercury occupies such a large fraction of its magnetosphere, most particles are not be able to complete a full drift orbit since drift paths on average either intersect the planet or the magnetopause, and the particles are thus removed from the system.

1.3 The bow shock

When the solar wind reaches the obstacle of Mercury's magnetic field it will be slowed down and deflected around it. A bow shock wave is generated as the solar wind hits the obstacle with supersonic speed, and a substantial fraction of the solar wind bulk flow kinetic energy is converted into thermal energy. The shock wave, which is a nonlinear wave, is an irreversible process that changes the state of the medium, for e.g., the temperature, pressure, and density. At the shock, the plasma density and temperature increase, implying a decrease in the flow speed to conserve mass and energy. Thus the shock wave is the mechanism by which the plasma transitions from supersonic to subsonic flow; upstream of the shock the solar wind flow speed is supersonic, while on the downstream side it is subsonic. The region right behind (downstream of) the bow shock where the solar wind is slowed and compressed is called the magnetosheath.

The shock can be thought of as a discontinuity, although in reality it has a finite thickness due to the kinetic processes at the shock. For a discontinuity to be a shock, there must be plasma flow across the shock surface accompanied by dissipation (i.e. the transformation of the kinetic energy of the flow to random thermal energy in the particles) and compression across the shock. Planetary bow shocks are fast shocks, corresponding to the fast mode wave in MHD. Across a forward fast-mode shock, the plasma pressure and magnetic field strength increase and the magnetic field bends away from the shock normal.

The shock structure is partly dependent on the upstream magnetic field direction relative to the shock surface. If the upstream magnetic field is perpendicular (or quasi-perpendicular) to the shock normal direction, particles cannot travel far back into the upstream region because their gyrational motion brings them back into the shock. Typical perpendicular shocks (Figure 1.4) show a shock foot where the magnetic field is gradually increasing in front of the main shock. Behind the main ramp where the magnetic field increases significantly, there is an overshoot with field values larger than the asymptotic downstream values. On the other hand, when the upstream



Figure 1.4: Schematic diagrams showing example bow shock magnetic field structure. Figure from Basic Space Plasma Physics, Wolfgang Baumjohann and Rudolf A. Treumann, Copyright 1997, Imperial College Press [30]; used with permission.

magnetic field is parallel (or quasi-parallel) to the shock normal direction, particles can travel back up into the upstream region along the field lines. In order for particles to outrun the shock and travel back upstream they must have a velocity higher than the escape velocity of the shock which is given by $\mathbf{v}_{esc} = \mathbf{v}_{nu}/\cos\theta$, where \mathbf{v}_{nu} is the upstream flow velocity normal to the shock and θ is the angle between the shock normal and the upstream magnetic field. Particles with higher velocities will move ahead of the slower escaping particles and thus higher energy particles will be observed further away from the shock. These foreshock particles drive instabilities such that the plasma and magnetic field properties both upstream and downstream

1.3. The bow shock



Figure 1.5: Schematic diagram showing regions of parallel and perpendicular shock. Just upstream of the parallel shock region is the foreshock. Figure from Basic Space Plasma Physics, Wolfgang Baumjohann and Rudolf A. Treumann, Copyright 1997, Imperial College Press [30]; used with permission.

show significant oscillations. Because the bow shock is a curved surface around the planet, for the same IMF direction the bow shock structure will be significantly different at different locations. Figure 1.5 shows an example bow shock schematic diagram at the Earth, where in regions where the IMF is perpendicular to the shock normal, the conditions will not be favorable for particles to travel back upstream and so no foreshock will form. However, as the angle between the shock and magnetic field gets smaller, more and more particles will be able to outrun the shock and form a foreshock region. The shock in this region will have quasi-parallel shock structure.

Using MHD equations, it is possible to derive the conservation relations (also known as the Rankine-Hugoniot shock-jump conditions) that prescribe the downstream plasma state in terms of the upstream parameters. How-
1.3. The bow shock

ever, the shock-jump conditions are a simplification because they are derived using a one fluid MHD approximation, thereby ignoring microscopic processes at the shock; thus they can be used as a guide and for comparison with observations of the flow upstream and downstream, and can test the assumptions that go into their derivation. These relations are derived in the shock frame, where the shock is stationary, and all involve parameters along the normal or transverse direction to the shock. Thus by measuring the magnetic field, flow speed, and density along the shock normal, it is possible to estimate how much the downstream conditions in the plasma have changed from that upstream. For the derivations of the Rankine-Hugoniot relations the interested reader is referred to [37].

The shock jump equations are given by

$$[\rho v_n] = 0, \tag{1.25}$$

$$\left[\rho v_n^2 + p + \frac{B^2}{2\mu_0}\right] = 0 \tag{1.26}$$

$$\left[\rho v_n \mathbf{v}_t - \frac{B_n}{\mu_0} \mathbf{B}_t\right] = 0 \tag{1.27}$$

$$\left[\rho v_n \left(\frac{1}{2}v^2 + \frac{\gamma}{\gamma - 1}\frac{p}{\rho}\right) + v_n \frac{B^2}{\mu_0} - \mathbf{v} \cdot \mathbf{B}\frac{B_n}{\mu_0}\right] = 0$$
(1.28)

$$[B_n] = 0 (1.29)$$

$$[v_n \mathbf{B}_t - B_n \mathbf{v}_t] = 0. \tag{1.30}$$

In these equations the square brackets denote subtraction of the downstream quantity in brackets from the upstream equivalent, subscript n denotes the component normal to the shock and t transverse to the shock, ρ is the flow density, **v** is the velocity, **B** is the magnetic field, μ_0 is the permeability of free space, γ is the polytropic index.

For the case of an exactly perpendicular shock, it is simple to derive the

1.3. The bow shock

expected maximum compression of the magnetic field, density, and velocity across the shock using the shock jump conditions. In such a case $B_n = 0$ and we examine the case where the upstream flow velocity is parallel to the shock normal, i.e. $\mathbf{v}_u = v_{un} \mathbf{\hat{n}}$. Across the shock, the mass flux must be non-zero, such that

$$\rho_u v_{un} = \rho_d v_{dn} \neq 0. \tag{1.31}$$

We define the compression ratio, $r = \rho_d/\rho_u$, which using equation 1.31 leads to the downstream velocity being defined in terms of the compression ratio as well as $v_{dn} = (1/r)v_{un}$. We can apply the third jump condition (equation 1.27), which due to the fact that $B_n = 0$ becomes

$$\rho_u v_{un} \mathbf{v}_{ut} - \rho_d v_{dn} \mathbf{v}_{dt} = 0, \tag{1.32}$$

implying that $v_{dt} = 0$ since v_{ut} is zero and ρv_n cannot be zero. We can use the jump condition in equation 1.30 and the fact that all the magnetic field is in the tangential component to substitute in for the compression ratio and get that $B_d = rB_u$. This shows that the magnetic field is compressed by the same amount as the flow density. Using equations 1.26 and 1.28 and substituting in for v_{dn} , B_d and eliminating p_d , it is possible to arrive at an equation that expresses the compression ratio in terms of the upstream parameters only:

$$(r-1)\left\{r^2\frac{2-\gamma}{M_A^2} + r\left(\frac{\gamma}{M_A^2} + \frac{2}{M_{ms}^2} + \gamma - 1\right) - (\gamma+1)\right\} = 0, \qquad (1.33)$$

where M_A is the Alfvén Mach number as defined earlier, and $M_{ms} = v_u \left(\frac{\rho_u}{\gamma p_u}\right)^{\frac{1}{2}}$ is the magnetosonic Mach number. In the high Mach number limit when M_A and $M_{ms} \gg 1$, equation 1.33 becomes $r = \frac{\gamma+1}{\gamma-1}$. γ is not dependent on the upstream parameters; for a monatomic gas $\gamma = \frac{5}{3}$, which yields a maximum compression of a factor of 4 for the magnetic field, density, and velocity. This number, although much quoted as the asymptotic limit for the magnetic field strength increase across the shock, is dependent on the polytropic index in the plasma and thus on the way that the plasma is heated. In Chapter 2, the example cases of Mercury's bow shock that are shown in Section 2.2 all have a compression factor of less than 4 (measured in the magnetic field strength). The upstream Mach number of the shock also defines the strength of the shock and is a measure of the amount of energy being processed by the shock; the higher the Mach number, the stronger the shock. In Chapter 2 we will explore how Mercury's bow shock varies with Alfvén Mach number.

1.4 The magnetopause and the magnetic cusps

The magnetopause demarcates the upper boundary of the magnetosphere and controls the flux of solar wind mass, energy, and momentum into the magnetosphere. Figure 1.6 illustrates how the interaction between the solar wind and a magnetized planet generates the magnetic cavity of the magnetosphere, marked by the magnetopause boundary. In this example case, in the equatorial plane of the planet the magnetic field is pointing northward. When charged particles in the solar wind approach the planetary magnetic field, the Lorentz force (see equation (1.2)) in the absence of electric fields, given by $q(\mathbf{v} \times \mathbf{B})$, deflects ions to the right and electrons to the left. The opposite motion of the charges produces a sheet current from left to right in the figure (dawn to dusk). The magnetic perturbations from this magnetopause current (or Chapman-Ferraro current as it was first proposed by [38]) cancels the planet's field sunward of the current and increases the field planetward. Thus if the magnetopause was a plane, the field strength just inside the boundary would be twice the dipole magnetic field. However, because the magnetopause is curved, the field immediately inside the near-equatorial dayside magnetopause is slightly greater than twice the dipole field.

The return current above the pole is from dusk to dawn, because the field points southward. Near the pole there is a singular point in the field, the neutral point, where this current sheet completely cancels the planet's field. These are called the magnetic cusps because here the field comes together



Figure 1.6: Currents on the magnetopause. By courtesy of Encyclopaedia Britannica, Inc., copyright 1994; used with permission.

at a cusp-like geometry. (Chapter 3 of this thesis will explore in great detail Mercury's northern cusp region.) The magnetopause current circulates in a sheet around the neutral point, and is symmetric about the magnetic equator, with a corresponding circulation about the southern neutral point. In this way the magnetopause current shields the magnetized planet from the solar wind, confining the planet's magnetic field to the magnetosphere.

By looking at a cross-sectional view of the magnetosphere (see Figure 1.7), more clarity can be gained about the magnetopause current, the cusps and other features of the magnetosphere structure. Inside the magnetosphere, the magnetic field is divided into two regions: 1) equatorward and 2) poleward of the cusp latitude. The division between these two regions is called the separatrix, which separates regions of closed field lines at mid-to-



1.4. The magnetopause and the magnetic cusps

Figure 1.7: Magnetospheric current systems. From: Hughes, W.J., The magnetopause, magnetotail, and magnetic reconnection, in Introduction to Space Physics, Eds M.G. Kivelson and C.T. Russell, Cambridge, 1995 [37].

low latitudes and open field lines at high latitudes. "Closed" here refers to field lines that have both ends map to the planet, whereas "open" refers to field lines that have one end map to the planet and the other end map to a tail lobe (or in the case of an open magnetopause, to the solar wind). In Figure 1.7, the magnetopause current flows out of the page equatorward of the cusps, and into the page poleward of the cusps. The magnetotail is generated by the fact that the magnetosheath flow convects open magnetic field lines toward the planet's nightside. In the magnetotail, the field direction is oppositely directed between the two hemispheres, which must be supported by the existence of a current sheet at the magnetic equator.

The location of the magnetopause around the planet corresponds to the surface across which the pressures of the internal magnetospheric magnetic fields and charged particles are balanced by the external solar wind particle and magnetic field pressure [39]. Thus the magnetopause is in pressure

balance, given by the equation

$$\left[P + \frac{B^2}{2\mu_0}\right]_{msh} = \left[P + \frac{B^2}{2\mu_0}\right]_{msp}$$
(1.34)

where msp stands for magnetosphere, msh for magnetosheath, P is a combination of particle thermal and dynamic pressure, and $B^2/2\mu_0$ is the magnetic pressure. On the left hand side we thus have the solar wind thermal and dynamic pressure as well as the IMF magnetic pressure, while on the right hand side we have the magnetic pressure of the planet's magnetic field and the plasma pressures. The dominant contributors to this equation are the solar wind dynamic pressure and the magnetic pressure of the planet's dipole field, although the solar wind thermal pressure becomes non-negligible at higher latitudes away from the subsolar point.

Using this pressure balance, we can obtain an estimate of the location of Mercury's magnetopause at the subsolar point. In Chapter 2, we describe how we determined the actual subsolar stand-off distance of Mercury's magnetopause using MESSENGER magnetic field observations as well as an empirical model of the boundary, but for now, we illustrate how a simple estimate can be achieved in order to build further intuition of the magnetopause. We assume the Newtonian approximation, an empirical relation from hypersonic flow theory [40], whereby the magnetosheath pressure at the subsolar point is given by

$$P_{msh} = \kappa \rho_{sw} v_{sw}^2, \tag{1.35}$$

where ρ_{sw} is the solar wind density, v_{sw} is the solar wind velocity, and κ is a constant that depends on how the solar wind flow is diverted around the planet's magnetic field. κ would be 1 in the ideal case where the bow shock and the magnetic field obstacle are close together and parallel, but is most often quoted as 0.88 if the flow is diverted in a more gradual manner [41].

To estimate the location of the magnetopause, we can rewrite equation

(1.34) in a simplified form in terms of the dominant components as

$$P_{msh} \approx 0.88 \rho_{sw} v_{sw}^2 \approx \frac{B_{msp}^2}{2\mu_0}.$$
(1.36)

We know that the magnetospheric magnetic field at the magnetopause should be slightly higher than twice the dipole field at that distance, and so we can express B_{msp} as

$$B_{msp} = aB_{dipole} = aB_{d0} \left(\frac{R_M}{R_{ss}}\right)^3, \tag{1.37}$$

where a = 2.4 has been deduced for Earth's magnetopause, B_{dipole} is the magnetic field strength of the dipole at the magnetopause, B_{d0} is the magnetic field strength of the dipole at the magnetic equator on the planet's surface, R_M is Mercury's radius, and R_{ss} is the distance from the planet's surface to the subsolar point on the magnetopause, i.e. the subsolar standoff distance. Substituting this into equation (1.36) yields that $R_{ss} \propto P_{msh}^{-\frac{1}{6}}$ and that

$$R_{\rm ss} \approx R_M \left(\frac{a^2 B_{d0}^2}{2\mu_0 \kappa \rho_{sw} v_{sw}^2} \right)^{\frac{1}{6}}.$$
 (1.38)

At the magnetic equator the magnetic field strength on the surface is given by $B = \frac{\mu_0 M_M}{4\pi r^3}$ and thus from having an estimate of the magnetic moment of Mercury, M_M , of 190 nT R_M^3 [19], we can calculate B_{d0} . We use the WSA-Enlil heliospheric model [42] predictions of the solar wind at Mercury to obtain averages of the solar wind density and speed, which are 51 cm^{-3} and 422 km/s, respectively. Substituting these values into equation (1.38), we estimate that Mercury's magnetopause stands off the solar wind at the subsolar point 1.35 R_M away from the planet's surface. This value is only 0.1 R_M less than the average R_{ss} established for Mercury in Chapter 2 using 3 Mercury years of observations confirming that the approximations made above are correct to first-order.

In Chapter 2 we also determine the shape of Mercury's magnetopause. We can however build up some intuition and derive the first-order 2-D shape of the magnetopause under the simple balance of forces used above. We start by revising the solar wind dynamic pressure relation to accommodate any point on the magnetopause surface. Since the dynamic pressure is exerted only in the anti-sunward direction, the magnetic pressure only needs to stand off this component everywhere around the planet. If \hat{n} is the magnetopause normal, the dynamic pressure incident on it away from the subsolar point is reduced to

$$P_{msh} = \hat{n} \cdot \hat{x} \rho_{sw} v_{sw}^2 = \cos \theta \rho_{sw} v_{sw}^2, \qquad (1.39)$$

where θ is the angle between the solar wind flow direction and the magnetopause normal. Using this revised dynamic pressure we arrive at the magnetopause distance from the planet's surface, $R_{\rm mp}$ away from the subsolar point

$$R_{\rm mp} \approx R_M \left(\frac{a^2 B_{d0}^2}{2\mu_0 \kappa \cos \theta \rho_{sw} v_{sw}^2} \right)^{\frac{1}{6}} = R_{\rm ss} (\cos \theta)^{-\frac{1}{6}}.$$
 (1.40)

We can see clearly from here that at $\theta = \pi/2$, $R_{mp} \to \infty$, which suggests the magnetopause extends far downtail. We can now derive a differential length, $d\vec{l} = dx\hat{x} + dz\hat{z}$, tangent to the magnetopause, where $dx = -\sin\theta dl$ and $dz = \cos\theta dl$. This leads to the ratio between the differential change in x and z

$$-\frac{dx}{dz} = \tan\theta = \frac{\sqrt{1 - \cos^2\theta}}{\cos\theta} = \frac{1}{\cos\theta}\sqrt{\frac{1}{\cos\theta} - \cos\theta} = \sqrt{\sec^2\theta - 1}.$$
 (1.41)

We can express equation (1.41) in terms of $R_{\rm mp}$ and $R_{\rm ss}$ through θ because

$$\sec\theta \approx \left(\frac{R_{\rm mp}}{R_{\rm ss}}\right)^6 = \left(\frac{x^2 + z^2}{R_{\rm ss}^2}\right)^3,$$
 (1.42)

which when substituted into equation (1.41) becomes

$$-\frac{dx}{dz} = \sqrt{\left(\frac{x^2 + z^2}{R_{\rm ss}^2}\right)^6} - 1.$$
 (1.43)

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We can gain some intuition about the shape by just choosing a few different values for x and z in equation (1.43). At $x = R_{ss}$ and z = 0, $\frac{dx}{dz} = 0$ showing that all of $d\vec{l}$ is in the \hat{z} direction as expected. For large negative x, when $R_{\rm mp} > R_{\rm ss}$ we find that $(R_{\rm mp}/R_{\rm ss}) \gg 1$ such that $\frac{dx}{dz} \approx (R_{\rm mp}/R_{\rm ss})^6$. This implies that $dx \gg dz$ and thus the magnetopause shape will become cylindrical past 2 $R_{\rm ss}$ down the magnetotail. In reality, the magnetopause tail can be more flared because the magnetic field in the tail is not dipolar (due to the stretched-out nature of the tail), while our assumption is based on a Newtonian approximation and a dipole field. In the far tail, the currents that form the tail become the dominant source of the magnetic field, not the dipole which is too far away to have the largest affect. From these physical arguments we thus expect Mercury's magnetopause to be curved (approximately hemispherical) on the dayside and approximately cylindrical on the nightside. In Chapter 2 we fit empirically motivated shape models to Mercury's magnetopause to establish its best-fit shape and location from nearly a year of observations.

1.5 The planetary magnetic field

There are four terrestrial planets in our solar system, but only Earth and Mercury have a currently active global dynamo field (although all the outer planets do possess dynamo-generated magnetic fields). In such dynamo operated magnetic fields, electric currents and magnetic fields are continuously induced by the movement of conducting fluid in the planet's interior. Three basic ingredients are needed for self-sustaining dynamo action to take place: a large volume of electrically conducting fluid material, planetary rotation, and a strong enough energy source to drive convective motions. These characteristics are in part set during planetary formation, but each one can change during the planet's lifetime. The purpose of this section is to introduce the basic idea behind planetary dynamos but not to delve into the field of dynamo theory, as this thesis is concerned with the results of planetary dynamos, namely the planetary magnetic field and its interaction with the solar wind. Larmor [43] first suggested that a dynamo process is responsible for creating and maintaining planetary magnetic fields. His hypothesis was that an initial magnetic field can be altered via interactions through fluid flow to amplify and regenerate the magnetic field. In 1955 Parker [44] presented the first modern model of the dynamo process, known as the $\alpha\omega$ model. In this model (which is still relevant today), the differential rotation of the fluid outer and solid inner core can generate toroidal (azimuthal) magnetic fields by stretching poloidal (non-azimuthal) magnetic fields threaded through the core (the ω effect). In order to maintain this generation of the toroidal field, the poloidal field must be generated through another process. The α effect accomplishes this task by relying on turbulent fluid convection to generate poloidal fields from the toroidal field, which closes the feedback loop.

The turbulent fluid motions and the planetary rotation act together to stretch out the magnetic field (which is "frozen in" to the core fluid). Due to the planetary rotation this generates magnetic field loops with approximately the same orientation and can thus give rise to a coherent global field. In such a way the dynamo process can generate and sustain magnetic fields of planetary bodies. Many variations on this basic dynamo model have been proposed, most of which require both planetary rotation and turbulent convection in a fluid core. Thus the fact that Mercury has a global dipole field is an indication (independent of the evidence from forced librations [4]) that the planet possess a liquid (or partially liquid) core in its interior.

Several dynamo models have been suggested to explain Mercury's unusually weak, equatorially offset, and axisymmetric magnetic field, although none have yet been able to closely match all aspects of the observed field. As described earlier, models with a stable layer at the top of the outer core [20, 21] show promise. Furthermore, a study currently in progress [45] has found that fields that are asymmetric about the geographic equator can be produced in dynamos driven by volumetrically-distributed buoyancy sources. The fact that Mercury's magnetopause lies very close to the surface and the large inferred size of the core (radius of 2020 km [46]), suggests that there may be a strong link between the internal dynamo generated field and the external fields of the magnetosphere. A feedback-dynamo model (in which the external fields of the magnetosphere penetrate to the core and act to weaken the dynamo generated field) [47] has been suggested for Mercury. However, this model predicts high odd harmonics in a spherical harmonic expansion of the field, which are not inferred from MESSENGER observations [18]. This does not however rule out the possibility of induction in Mercury's core, which may act to stiffen the dayside magnetopause against changes in solar wind pressure [48]. Better resolution of the magnetic field strength at the surface of the planet can provide further constraints for dynamo models of Mercury and this is one of the underlying motivations for Chapter 5 of this thesis.

1.6 Relevant MESSENGER instruments

In this section we summarise the details of the Magnetometer (MAG) instrument onboard MESSENGER, from which data were used in all chapters of this thesis, as well as the Fast Imaging Plasma Spectrometer (FIPS), data from which were used in Chapters 4 and 5. For a thorough description of MAG the reader is referred to [49], and for more details on FIPS the reader is referred to [50].

1.6.1 The Magnetometer

The MESSENGER Magnetometer is a tri-axial fluxgate magnetometer mounted on a 3.6-m-long boom. The MESSENGER fluxgate magnetometer consists of a permeable ring core driven in alternative saturation states by a toroidal winding at a frequency of 15 kHz. The ring core and winding are surrounded by a second set of pickup windings that sense any change in net magnetic flux. If there is a net field along the axis of the pickup winding coil then when the ring core switches saturation polarities there will be a brief pulse of net flux in the pickup winding since one side of the ring core will transition polarity before the other side. The sense of this net flux pulse is in the direction of the background field and is the same sense every half cycle of the 15 kHz drive. Synchronous detection at 30 kHz in the pickup winding is used to monitor any net flux pulses and a bias current is sent to the pickup coil to null the fluxgate response. The voltage applied to drive this nulling current is directly proportional to the component of the ambient magnetic field along the axis of the pickup coil. In a tri-axial fluxgate magnetometer, three such sensors are mounted in such a way as to measure the three vector components of the magnetic field. MAG outputs three analog signals which are converted simultaneously to digital signals by three independent 20-bit sigma-delta A/D converters at a rate of 20 conversions per second. Prior to conversion these signals are low-pass filtered using a 10-Hz anti-alias lowpass filter to limit the input bandwidth into the A/D converter.

MAG has a range of ± 1530 nT full scale with a 0.047 nT resolution during Mercury operation. In order to minimize variable and static spacecraftgenerated fields at the sensor, the MAG instrument underwent a strict magnetic cleanliness program prior to launch. Observations after boom deployment indicated that there is a fixed residual field, however, it is less than a few nT at the location of the magnetometer. A variable contamination field measured at the magnetometer is below 0.05 nT.

To meet the science objectives of the MESSENGER mission as well as to accommodate variable telemetry rates, MAG can be operated at variable sample rates (from 0.01 s^{-1} to 20 s^{-1}) adapted to various mission phases. A combination of digital filtering and sub-sampling is used to provide the wide range of output rates. MAG sampling inside the magnetosphere is at the highest rate (20 Hz) in order to resolve highly varying magnetospheric fields, while sample rates in the interplanetary medium are at least as high as 2 samples/s. In addition to the vector magnetic field samples, the average output amplitude of a 1-10 Hz bandpass filter is also evaluated for the zmagnetometer axis, although this can be commanded to be any of the three axes. This is a measure of the average 1-10 Hz ambient field fluctuations, called B_{AC} (used in Chapters 2 and 3 of this thesis), which is made regardless of the sample rate of the vector data in order to provide an uninterrupted measure of the field variability. The B_{AC} measure is also used to trigger high-time-resolution sampling in 8-min segments to capture events of interest when continuous high-rate sampling is not possible.

1.6.2 The Fast Imaging Plasma Spectrometer

FIPS is part of the Energetic Particle and Plasma Spectrometer (EPPS) package on MESSENGER, and it measures the energy, angular, and compositional distribution of low-energy ions (< 50 eV/charge to 20 keV/charge) in Mercury's space environment. The sensor was specifically developed and designed for MESSENGER. It is highly compact and lightweight, and houses a new electrostatic analyzer system geometry that enables it to have 1.4π steradians field of view (FOV). FIPS is a time-of-flight mass spectrometer, in which an ion's mass-to-charge ratio is determined through a time measurement.

FIPS is made up of an electrostatic analyzer (ESA), a time of flight detector, and a compartment that houses the instrument electronics (see Figure 1.8 for a cross-sectional view of the FIPS sensor). The electrostatic analyzer uses an electric field to filter and focus the particles that are allowed to pass through to the time of flight detector chamber: only particles of a specific energy-per-charge range pass through the analyzer. The ESA covers the energy-per-charge range in 64 logarithmically spaced steps. FIPS is normally operated in one of two stepping rates, one step per second (normal mode) or one step per 100 milliseconds (burst mode), which result in integration times of 64 s and 8 s, respectively. Inside the magnetosphere FIPS is operated in burst mode and measures energy per charge from 0.1 to 13 keV/e [51].

Inside the ESA, the entering ion's trajectory is bent by the deflector plates in the analyzer, after which the collimator filters out particles with certain trajectories. Next, the ion is post-accelerated by a potential drop of up to -15 kV to give low energy ions sufficient energy to penetrate the carbon foil. The ion then enters the time of flight detector and impacts the carbon foil, which causes the foil to eject secondary electrons. The electrons move ahead of the ion and their path will be bent to ensure that they collide with the start microchannel plate (MCP) detector, thereby recording the ion's position and providing a timing-start signal. On the other hand, the ion will pass straight through the time of flight chamber, hitting the stop MCP



Figure 1.8: Cross-section of the FIPS sensor showing major functional components. Figure from Andrews et al. (2012), Space Science Reviews, 131, 523-556 [50]. Reprinted with permission from Springer Science + Business Media.

detector, and thereby stopping the timer. The difference in time between the electrons and ion hitting the two detectors is used to calculate the initial ion speed. This can be obtained after a correction is made for the travel time of the secondary electrons from the foil to the start detector. The mass-percharge follows from the known energy-per-charge and the measured time of flight. This allows for the reconstruction of the distribution functions of different mass-per-charge species. The microchannel plate detectors are electron multipliers, which when struck by a particle, start a cascade of electrons that propagate through the microchannels amplifying the original signal of the particle by several orders of magnitude. These electrons exit the channels and are detected by measuring the total current on a single metal anode.

The angular resolution of FIPS, in terms of determining the incident direction of the ion is 15°. It has a nearly hemispherical $(1.4\pi \text{ sr})$ instantaneous FOV. The FOV has conical symmetry about the z-axis of the sensor's reference frame and is defined as the region of solid angle bounded by two

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Figure 1.9: Top view of a portion of MESSENGER which depicts the obstructions in the solar direction by the spacecraft sunshade and the FIPS FOV. The figure (from Gershman et al. (2012), Journal of Geophysical Research, 117, A00M02 [52]) also shows that FIPS cannot observe the solar wind, however, this is not a focus of this thesis. Figure reprinted with permission from John Wiley and Sons.

nested cones. Both of these cones have their vertices at the origin of the FIPS coordinate system and have angles with respect to the z-axis of 15° and 75° , respectively. The FOV is the region contained within the 75° cone excluding the region contained within the inner 15° cone. This excluded region shows up as a gap in Figure 1.9 in the centre of the unobstructed FOV of FIPS. Such a constrained field of view places limitations on the reconstruction of the 3–D ion distribution from this sensor, as will be shown in Chapter 4.

Chapter 2

Mercury's magnetopause and bow shock from MESSENGER Magnetometer observations¹

In this chapter, we establish the average shape and location of Mercury's magnetopause and bow shock from orbital observations by the MESSEN-GER Magnetometer [53]. We fit empirical models to midpoints of boundary crossings and probability density maps of the magnetopause and bow shock positions. The magnetopause was fit by a surface for which the position rfrom the planetary dipole varies as $(1 + \cos \theta)^{-\alpha}$, where θ is the angle between r and the dipole-Sun line, the subsolar stand-off distance $R_{\rm ss}$ is 1.45 $R_{\rm M}$ (where $R_{\rm M}$ is Mercury's radius), and the flaring parameter $\alpha = 0.5$. The average magnetopause shape and location were determined under a mean solar wind ram pressure, P_{Ram} , of 14.3 nPa. The best-fit bow shock shape established under an average Alfvén Mach number $(M_{\rm A})$ of 6.6 is described by a hyperboloid having $R_{\rm ss} = 1.96 R_{\rm M}$ and an eccentricity of 1.02. These boundaries move as P_{Ram} and M_{A} vary, but their shapes remain unchanged. The magnetopause $R_{\rm ss}$ varies from 1.55 to 1.35 $R_{\rm M}$ for $P_{\rm Ram}$ in the range 8.8 to 21.6 nPa. The bow shock $R_{\rm ss}$ varies from 2.29 to 1.89 $R_{\rm M}$ for $M_{\rm A}$ in the range 4.12 to 11.8. The boundaries are well approximated by figures of revolution. Additional quantifiable effects of the interplanetary magnetic

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2.1. Introduction

field are masked by the large dynamic variability of these boundaries. The magnetotail surface is nearly cylindrical, with a radius of ~ 2.7 $R_{\rm M}$ at a distance 3 $R_{\rm M}$ downstream of Mercury. By comparison, Earth's magnetotail flaring continues until a downstream distance of ~ 10 $R_{\rm ss}$. The modeled cylindrical shape of Mercury's magnetopause suggests that magnetic flux has a short residence time in the tail, and return convection of flux from the tail to the dayside proceeds rapidly, consistent with the expected ~ 2 minute Dungey cycle period at Mercury.

2.1 Introduction

The boundaries of Mercury's magnetosphere reflect fundamental processes of the solar wind interaction with the planet's dipolar magnetic field [24, 54]. These boundaries are the bow shock, across which the solar wind is compressed and deflected around Mercury, and the magnetopause, which is the current layer separating the shocked solar wind plasma and interplanetary magnetic field (IMF) from the planetary magnetic field. The bow shock is a fast magnetosonic shock wave that "stands" in the solar wind while diverting the solar wind around the planet's magnetospheric cavity [39]. The bow shock changes shape and stands closer or farther from the planet in response to variations in solar wind Mach number and, to a lesser extent, IMF direction [26, 39, 55]. The magnetopause location and shape are determined principally by the pressure exerted on the magnetopause by the shocked solar wind plasma, which scales with the solar wind ram pressure, balanced by the planetary magnetic field [39, 56]. Accordingly, the dynamic pressure of the solar wind and the magnetic pressure of the magnetosphere are the dominant factors determining the location and shape of the magnetopause.

Because the distribution of magnetic flux within a magnetosphere is determined both by the intrinsic planetary field and the external currents, magnetic reconnection, which drives some of these external currents, also affects magnetopause shape and position. Most important among these reconnection-driven effects are the inward "erosion" of the dayside magnetopause to lower altitudes by transfer of magnetic flux to the tail [57]

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and the outward "flaring" of the nightside magnetopause as the magnetotail is loaded with magnetic flux [58]. These expansions and contractions of the dayside and nightside magnetosphere take place on timescales comparable to the Dungey cycle that governs the circulation of magnetic flux and plasma within the magnetospheres of Earth-like planets. The characteristic timescale of the Dungey cycle is ~ 1 h at Earth and ~ 2 min at Mercury [28].

It has been suggested that reconnection may have a greater effect on magnetopause location at Mercury than at Earth [57]. This prediction has been supported by the extreme loading and unloading of Mercury's magnetotail observed during the third flyby of Mercury by the MESSENGER spacecraft [29, 59]. Further, it has also been predicted that magnetic fields associated with induction currents in Mercury's interior may act to oppose and limit the solar wind's ability to compress or expand the extent of the dayside magnetosphere [60-63]. Mercury has an internally generated, axially aligned dipolar magnetic field with a moment of 190 to 195 nT $R_{\rm M}^3$ (where $R_{\rm M}$ is Mercury's mean radius, 2440 km) that is offset northward from the geographic equator by 0.2 $R_{\rm M}$ [17–19]. The combination of a weak dipole moment, the absence of a conducting ionosphere (i.e., no magnetosphereionosphere coupling), and the predominantly quasi-parallel subsolar shock conditions (due to the small Parker spiral angle at Mercury's orbit) distinguish Mercury from other planets in our solar system that possess magnetic fields.

The Earth's bow shock has been studied extensively both observationally and theoretically (e.g., [55]). It is a highly dynamic boundary, controlled by temporal variations in solar wind characteristics. The general shape of the shock has been investigated with empirical models (e.g., [26, 64]), gas dynamic flow models (e.g., [65, 66]), and magnetohydrodynamic models (e.g., [67]) and is well described by a conic section. Formisano et al. [68] found that the subsolar bow shock position moves outward during conditions of low Alfvén Mach number (M_A). Peredo et al. [55] confirmed that M_A primarily controls the bow shock shape, but in contrast to the findings of the earlier study [68], they observed that the subsolar shock moves earthward and the flanks flare outward during times of low M_A .

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The average shape of the terrestrial magnetopause has been described with several empirical models (e.g., [69–71]) and magnetohydrodynamic models [72]. The dynamic pressure of the solar wind affects the size of the magnetospheric cavity as well as the shape of the magnetopause. The subsolar magnetopause stand-off distance is observed to decrease with increased dynamic pressure [73, 74], and the boundary shape is observed to flare with increased pressure [69]. Reconnection-related effects can also influence the dayside magnetopause location and the flaring of the tail. On the dayside, a southward IMF component will facilitate reconnection at low latitudes, which can erode the magnetopause near the subsolar point and flare the magnetopause on the night by adding flux to the tail. The magnetopause has been observed to move planetward by as much as 1-2 $R_{\rm E}$ (where $R_{\rm E}$ is Earth's radius) due to the erosion of the boundary under southward IMF [57, 73, 75]. However, a statistical study of these effects at Mercury with a large database of magnetopause crossings, such as the one reported here, would require knowledge of IMF direction on timescales of 1 min or less because of the very short Dungey cycle time at Mercury. The lack of an upstream monitor therefore limits the scope of our investigation with respect to the effects of reconnection on magnetopause position.

The solar wind and IMF play major roles in influencing the bow shock and magnetopause, and those roles must be understood quantitatively in order to model the internal magnetic field, the magnetosphere-solar wind interaction, and the access of solar wind and magnetospheric charged particles to the planetary surface. At Mercury the solar wind density is approximately an order of magnitude higher and the IMF magnitude a factor of 5 higher than at Earth [23], whereas the planetary magnetic moment is only 0.06% of Earth's. Mercury's orbit is eccentric, so the planet is subjected to different solar wind conditions at perihelion and aphelion. Mercury's magnetopause and bow shock have been studied from limited data obtained during flybys by the Mariner 10 [76] and MESSENGER spacecraft [77]. Russell [76] fit shape models to both the magnetopause and bow shock with data from the two Mariner 10 flybys, and Slavin et al. [77] updated these boundary shapes with MESSENGER flyby data. Slavin et al. [77] also looked at the boundary crossings during northward IMF from data taken during the first MESSENGER flyby. However, even the combined Mariner 10 and MESSENGER flyby dataset comprised only six pairs of bow shock and magnetopause crossing points.

In this chapter we present analyses of Mercury's bow shock and magnetopause shape obtained from orbital magnetic field observations by the MESSENGER spacecraft over a span of three Mercury years under a variety of IMF and solar wind conditions. The objective of this study is to characterize the underlying shapes of both the magnetopause and bow shock and, insofar as the data permit, to assess how these shapes are affected by the solar wind and IMF. We have analyzed the observations in two ways. First, from the locations of the inner and outer magnetopause and bow shock crossings we defined a mean crossing point, and we found the average boundary shape from the ensemble of crossing points. In the second method, the probability of spacecraft residence within the range of magnetopause or bow shock crossings on each pass has been used to build a probability density map of the two boundaries. In Section 2.2, we describe the magnetic field observations and how the boundaries were identified. In Section 2.3, we establish the general shape of both the magnetopause and the bow shock from the mean locations (Section 2.3.1) and the probability densities (Section 2.3.2). In Section 2.4, we assess how the boundaries respond to solar wind forcing, and the results and conclusions are given in Section 2.5 and 2.6, respectively.

2.2 Magnetic field observations: boundary identifications

The MESSENGER spacecraft was inserted into orbit about Mercury on 18 March 2011. The initial orbit had a 200 km periapsis altitude, 82.5° inclination, 15,300 km apoapsis altitude, and 12 h period. During each orbit, MESSENGER typically spent 1 to 2 h inside the magnetosphere; the rest of the time was spent in the magnetosheath and in the interplanetary medium. For this study we used three Mercury years of MESSENGER Magnetometer (MAG) [49] data starting on 23 March 2011 and extending through 19 December 2011, providing repeated coverage at all local times. To conserve data volume during this period, MAG operated at variable sample rates, with high-rate (20 samples/s) data collection primarily in effect during transits of the magnetosphere. Sample rates in the interplanetary medium were at least as high as 2 samples/s or higher, and a channel to record fluctuations at 1-10 Hz operated continuously to provide an uninterrupted measure of the field variability.

Magnetic field data were analyzed in Mercury solar orbital (MSO) coordinates. In MSO coordinates $x_{\rm MSO}$ is positive sunward, $z_{\rm MSO}$ is positive northward, $y_{\rm MSO}$ is positive duskward and completes the right-handed system, and the origin is at the center of the planet. To analyze boundary locations the spacecraft positions were translated into a system centered on the planetary dipole [18]. The vector components of the magnetic field in the planet-centered and dipole-centered systems are identical. Because the bow shock and magnetopause are ordered by the solar wind flow in the frame of Mercury's orbital motion, the spacecraft position data were transformed into an aberrated system such that the +x direction is anti-parallel to the solar wind velocity relative to Mercury. The average aberration angle at Mercury is about 7° toward dawn. However, because of Mercury's variation in orbital speed between perihelion and aphelion, as well as variability in the solar wind speed, the aberration angle varied by a factor of about 3, from 3.5° to 10.2° , during the period of our study. The aberration correction for both the magnetopause and the bow shock crossings was calculated from the planet's instantaneous orbital speed together with predictions of solar wind speed obtained from the WSA-Enlil heliospheric model [42] within four minutes of the times of the boundary crossings. Because the Fast Imaging Plasma Spectrometer (FIPS) on MESSENGER does not typically see the solar wind, we do not have in-situ estimates of solar wind properties, so we use WSA-Enlil model predictions of solar wind parameters in this study. Use of model predictions for the solar wind parameters introduces some level of uncertainty into the normalized boundary shapes derived in Section 2.4

of the paper, with the highest uncertainty most likely introduced for the cases of extreme events, when the solar wind pressure is predicted to be highest. Benchmarking of the WSA-Enlil model at Mercury has been initiated [52, 78]; however, the long-timescale (> 44 day) variations in the model outputs that are modulated by Mercury's orbital distances from the Sun are persistent over the three Mercury years of data analyzed here. Our aim here is not to use WSA-Enlil for event studies, but to capture the average annual variation in the solar wind ram pressure and density and use these variations to correct our average boundary shapes.

The IMF magnitude assigned to each crossing was evaluated as a 1 h average of MAG data upstream of the outermost bow shock encounter. The Alfvén Mach number, $M_{\rm A} = \frac{v_{sw}}{[B_{\rm IMF}/(\mu_0 \rho)^{0.5}]}$, where v_{sw} is the solar wind speed, B_{IMF} is the magnetic field magnitude in the IMF, μ_o is the permeability of free space, and ρ is the solar wind plasma density, was calculated using IMF values that were estimated from the 1 h IMF averages and WSA-Enlil model-generated solar wind speed and density. We have shown previously that the 1 h IMF averages are suitable for determining the IMF $B_{\rm x}$ direction, which is the dominant IMF component at Mercury, but they are not suitable for establishing the IMF B_y and B_z directions, which vary on timescales less than 1 h [Winslow et al., 2012]. The magnitude of the IMF is dominated by the $B_{\rm x}$ component and is thus also steady on the 1 h time scale, i.e., the average duration of a transit of the magnetosphere by the MESSENGER spacecraft. The resolution of these measurements is suitable for investigating the effect of solar wind pressure and Alfvén Mach number on magnetopause and bow shock positions, but not the reconnection-driven effects that depend on IMF orientation and take place on timescales of one minute to a few minutes [28, 29].

Magnetopause and bow shock crossings were identified on every orbit, both before and after the magnetospheric transit and denoted as the inbound and outbound crossings, respectively. On almost every orbit, multiple crossings of each boundary were observed as a result of motion of the boundary relative to the spacecraft. Although the repeated crossings were often difficult to distinguish individually, the first and last boundary encounters were readily identified. Thus rather than attempt uncertain identifications of every boundary crossing within each passage through the boundary region, we identified the times of the innermost and outermost crossing on each pass for the bow shock and magnetopause. Our approach has the added benefit that each pass corresponds to an independent sample of external IMF and solar wind conditions, whereas statistical analyses that count every crossing equally in passes with multiple crossings will overweight such passes. We aim to provide demarcations for the inner and outer limits of the boundaries, within and outside of which the spacecraft was clearly located in the magnetosphere, magnetosheath, or interplanetary medium. Wave characteristics (including foreshock waves, non-linear quasi-parallel shock phenomena, and magnetopause boundary waves) are beyond the scope of this paper.

All boundary crossings were picked by visual inspection with the following criteria. For the bow shock, the inbound outer limit was identified as the time at which the first sharp increase in the magnitude $|\mathbf{B}|$ of the magnetic field was observed. The inner limit was identified as the time of the last sharp increase in $|\mathbf{B}|$. These criteria worked well when the IMF was oriented somewhat oblique to the planet-Sun line, that is, for quasi-perpendicular shock conditions. A perpendicular shock forms when the shock-normal direction is perpendicular to the IMF direction, whereas a parallel shock occurs when the shock normal is parallel or anti-parallel to the IMF. For near-parallel shock conditions, there was often little or no increase in $|\mathbf{B}|$, but the bow shock boundary was marked by the onset of large variability in $|\mathbf{B}|$. Sometimes these modulations grew gradually, in which case we chose the outermost excursions in $|\mathbf{B}|$ that were distinctly larger than the upstream variability. The outbound bow shock was picked in a similar fashion, that is, a sharp decrease in the field magnitude marked the boundary.

For the magnetopause, the crossings were most readily identified when the shear angle between the direction of the magnetic field in the magnetosphere and that in the magnetosheath was larger than about 45° , because the field rotation is a direct signature of the magnetopause current layer. On the dayside and on the flanks, the shear is typically in B_y and B_z , whereas on the nightside the shear is mostly in B_x . The inner boundary was iden-

tified by the innermost substantial rotation away from the magnetospheric field direction and the outer boundary by the outermost rotation toward the ambient magnetospheric field direction, excluding the background magnetosheath variability. In many cases, however, the field direction in the magnetosheath was the same as (or close to) that inside the magnetosphere. For example, the magnetic shear can be low at low dayside latitudes when the IMF is northward and on the nightside when the IMF is anti-sunward such that the southern lobe field is parallel to the draped magnetosheath field. In such cases, although the local magnetic shear is low, there were other signatures in the magnetic field that indicated transition between magnetosheath and magnetosphere regimes. These signatures include an increase in magnetic fluctuations in the magnetopause layer (documented from Earth's magnetopause [79, 80]) and a change in the character of the low-frequency fluctuations on the magnetosheath side of the boundary. The higher-frequency magnetic fluctuations are recorded by the 1-10 Hz fluctuation channel (or $B_{\rm AC}$), which provides an average amplitude of the 1-10 Hz bandpass-filtered field. In addition, on the dayside the inner magnetopause boundary is also often indicated by either a step-wise increase in $|\mathbf{B}|$ or the onset of an inward gradient in $|\mathbf{B}|$ on the magnetosphere side of the boundary. However, these signatures were not always sharp, and in some cases it was difficult to identify the magnetopause boundary. Boundary crossing choices were made conservatively such that the inner edge was definitely inside the magnetosphere and the outer edge was definitely in the magnetosheath. Using plasma measurements from FIPS to identify the boundary crossings is beyond the scope of this paper and would require careful accounting of the FIPS look direction, since its field of view is 1.4π sr. Several dozen comparisons show excellent correspondence between the boundaries identified from MAG observations and abrupt changes in FIPS proton count rates.

Data from the first magnetospheric transit on 12 October 2011 (orbit 418) are shown in Figure 2.1 together with expanded views of the inbound and outbound boundary passages. In this case, the shock conditions were oblique (perpendicular shock geometry) and there was high magnetic shear

across the magnetopause on both the inbound and outbound crossings. The spacecraft entered the magnetosphere from the dayside and exited on the nightside. The spacecraft first grazed the shock at 3:43:59 Coordinated Universal Time (UTC) before passing through it at 3:45:29 UTC. On this transit, the shock was also marked by an increase in high-frequency fluctuations, $B_{\rm AC}$, although most of the fluctuations before the shock crossing are attributed to foreshock waves. The spacecraft then traveled through the magnetosheath and encountered the high-shear magnetopause on the dawn side, marked by a rotation in B_z and B_y and an increase in B_{AC} shortly before the magnetopause boundary. On the outbound portion of the orbit, the spacecraft first encountered the magnetopause at 5:19:56 UTC and was finally in the magnetosheath at 5:25:05 UTC. The rotation in the magnetic field is evident in $B_{\rm x}$ as the spacecraft exited the magnetosphere from the southern tail lobe and was also associated with a rise in $B_{\rm AC}$. The spacecraft then crossed the bow shock twice on its path back into the interplanetary medium.

Often the boundaries were less clear, and Figure 2.2 shows data from the first magnetospheric transit on 5 July 2011 (orbit 218) with such crossings. On this orbit the shock was quasi-parallel on the dayside, which caused the large modulations in $|\mathbf{B}|$ near the bow shock crossing. The magnetic field inside the magnetosheath was highly variable, with large, quasi-periodic rotations in B_y and B_z up to the magnetopause. The magnetic shear across the magnetopause was low, so the decrease in B_{AC} and the increase in the total field magnitude were taken to indicate the magnetopause crossing. On the outbound part of the orbit, the magnetic shear was again low across the magnetopause, with slight rotations visible in B_y and B_x , but the crossing was indicated by an increase in B_{AC} . The field magnitude decreased as the spacecraft crossed the magnetopause but increased inside the magnetosheath until it reached the oblique bow shock boundary and decreased abruptly at 7:45:54 UTC.

The times of the inner and outer edges of the recorded magnetopause and bow shock crossings for all the data presented in this study are given in the Supplemental Materials of [53]. There are 1,065 magnetopause and 1,084



Figure 2.1: (a) MESSENGER Magnetometer data for the first magnetospheric transit on 12 October 2011 (orbit 418). Left axes give the scales for B_x (red), B_y (light green), B_z (blue), and $|\mathbf{B}|$ and $-|\mathbf{B}|$ (black); the right axis is the scale for B_{AC} (dark green). Vertical lines denote the crossing times of the inner and outer edges of the bow shock (dashed) and magnetopause (dotdashed). (b) Close-up view of the inbound portion of the orbit. (c) Close-up view of the outbound portion of the orbit.

bow shock crossings altogether in our dataset. The number of bow shock crossings exceeds the number of magnetopause crossings because the Magnetometer was switched off to conserve power for parts of 19 orbits during



Figure 2.2: (a-c) MESSENGER Magnetometer data for the first magnetospheric transit on 5 July 2011, orbit 218, in the same format as in Figure 2.1.

MESSENGER's first long-eclipse season in orbit around Mercury. On these orbits both the inbound and outbound bow shock crossings were recorded, but typically only the inbound (and not the outbound) magnetopause crossings were observed.

2.3 Average boundary shapes

The first step in analyzing the crossing data was to determine the best boundary shapes for all crossings together, effectively averaging over IMF and solar wind conditions. This averaging was accomplished by fitting empirical models to the magnetopause and bow shock crossing locations. The boundary locations were specified with two different techniques. First, we used the mean locations of the boundaries and fit empirical shapes to these directly. Second, we used a probabilistic measure of residence within the boundary regions to identify the locations of maximum residence probability. Empirical shapes were then fit to these probability density maps, with the established probabilities used as weights in the fitting. In the analysis that follows, all positions are in aberrated coordinates (x, y, z), where x and y are the aberrated $x_{\rm MSO}$ and $y_{\rm MSO}$ coordinates, respectively, and $z = z_{\rm MSO}$. In addition, we have assumed that the boundaries are figures of revolution about the line through the dipole center that parallels the x axis; the validity of this assumption is quantitatively tested and confirmed in Section 2.5. The northward offset of the planetary dipole is included in the definition of the distance from the axis of revolution, given by $\rho = \sqrt{y^2 + (z - z_d)^2}$, where $z_{\rm d} = 0.196 R_{\rm M} [17-19]$.

2.3.1 Midpoint fits

In the first approach to determining boundary locations, model curves were fit to the average crossing points, that is, the midpoint between the inner and outer edge of the boundary location on each pass. The inner and outer limits of the boundaries were assigned as the uncertainty range. This method allows direct comparison of our results with approaches that have been used historically to determine boundary shapes (e.g., [69, 77]).

2.3.1.1 Magnetopause

For the magnetopause, we used a paraboloid conic section [81, 82] as well as the model shape proposed by Shue et al. [69] to fit our crossings.



Figure 2.3: Midpoints between the inner and outer magnetopause crossing positions identified from MESSENGER Magnetometer data from 23 March 2011 through 19 December 2011. Error bars show the distance between the inner and outer crossing. Curves show the best-fit paraboloid (blue) to the dayside crossings and the best-fit Shue et al. model (red), as well as models from previous studies by Slavin et al. [77] (green) and Russell [76] (yellow). The paraboloid has parameters given by $R_{ss} = 1.5 R_{\rm M}$ and $\gamma = 1$, whereas the Shue et al. model is given by $R_{ss} = 1.45 R_{\rm M}$ and $\alpha = 0.5$.

Figure 2.3 shows the midpoints of the magnetopause crossings from the three Mercury years of data analyzed in this study. To establish a time-averaged magnetopause shape from the crossing points, we modeled them in $\rho - x$ space. In our boundary fits, we used a grid search method that minimized the root-mean-square (RMS) residual of the perpendicular distance of the observed midpoints from the model boundary.

The paraboloid fit is motivated by the magnetospheric model of Alexeev et al. [82], which was derived with a parabolic parameterization of magnetopause shape. Past studies of the magnetopause shapes around other planets have also involved fits to conic sections (e.g., [76]). The paraboloid model shape is described by

$$x(\rho) = -\left(\frac{\gamma^2 + 1}{4R_{ss}}\right)\rho^2 + R_{ss}$$
(2.1)

where γ is a flaring parameter and R_{ss} is the subsolar magnetopause distance [81]. Any value of $\gamma > 1$ is physically unreasonable, because it gives a subsolar stand-off distance that is not the minimum distance between the magnetopause and the planet. Setting $\gamma = 1$, we find $R_{ss} = 1.25 R_{\rm M}$ as the best-fit paraboloid. This model does not provide a good visual fit, however, to data either on the dayside or in the distant tail region. Relaxing the constraint of $\gamma > 1$ gives a better fit to the crossings on the nightside, but the resulting model still does not fit the dayside points. We find that the paraboloid model represents the dayside magnetopause shape best when we exclude the tail crossings, yielding $R_{ss} = 1.5 R_{\rm M}$ and $\gamma = 1$ (Figure 2.3).

We also fit the magnetopause crossings with the functional form proposed by Shue et al. [69] (hereafter referred to as the Shue et al. model) and given by

$$R = \sqrt{x^2 + \rho^2} = R_{ss} \left(\frac{2}{1 + \cos\theta}\right)^{\alpha}$$
(2.2)

where R is the distance from the dipole center, $\theta = \tan^{-1} \left(\frac{\rho}{x}\right)$, and α is another flaring parameter that governs whether the magnetotail is closed $(\alpha < 0.5)$ or open $(\alpha \ge 0.5)$. This model has been used successfully to model the Earth's magnetopause [69, 83], as well as the magnetopause of other planets, such as Saturn [84]. The Shue et al. model that best fits our midpoint magnetopause crossings yields parameter values of $R_{ss} = 1.45 R_{\rm M}$ and $\alpha = 0.5$ (shown in Figure 2.3). As can be seen from Figure 2.3, the Shue et al. model provides a better representation of the magnetopause crossings than the paraboloid model. Even for the Shue et al. model, the best-fit parameters are not tightly constrained, as similar RMS values are achieved over a range of values for R_{ss} and α (Figure 2.4). This behavior is in part because of a trade-off between the R_{ss} and α parameters, imposed by the



Figure 2.4: RMS misfit between the midpoints of the magnetopause crossings and the Shue et al. model, as a function of the subsolar stand-off distance, R_{ss} , and the flaring parameter, α .

observation geometry, and in part because of the spread in the magnetopause crossing positions most likely caused by the dynamics of variable solar wind and IMF conditions.

To better constrain our best-fit parameters, we conducted analyses of the residuals (the perpendicular distance of our crossings from the model boundary) for a range of models with $R_{ss} - \alpha$ parameter pairs from the minimum misfit region in Figure 2.4. We find that as we depart from the best-fit parameters on either side of the misfit well, the residuals have distributions with a non-zero mean and are not Gaussian. Thus even though the RMS misfit is not very different from the absolute minimum value in the minimum misfit region, the models generated by the parameter pairs in that region describe the data less well than our best-fit model. In addition, models have residuals that vary systematically with $x_{\rm MSO}$. This situation can be visualized by taking, for example, a model curve lying within the minimum misfit region with $R_{ss} = 1.25 R_M$ and $\alpha = 0.6$. Such a boundary yields residuals that are systematically positive on the dayside (i.e., the boundary is too close to the planet relative to the observations) and are systematically negative on the nightside (i.e., the boundary is too flared and farther from the planet on average than the data). On the other hand, the residuals of the best-fit model with $R_{ss} = 1.45 R_M$ and $\alpha = 0.5$ have zero mean and a Gaussian distribution and show no systematic variation with x_{MSO} . Also, the RMS misfit (~ 590 km) for the best-fit Shue et al. model is a factor of 1.5 lower than that of the best-fit paraboloid model (which was fit only to the dayside data), and thus we use only the Shue et al. curve to model the magnetopause for the remainder of this discussion.

2.3.1.2 Bow shock

To characterize the shape of the bow shock, the midpoints of the bow shock crossings were modeled by a conic section given by [77]:

$$\sqrt{(x-x_0)^2 + \rho^2} = \frac{p\epsilon}{1 + \epsilon \cdot \cos\theta}$$
(2.3)

where the focus of the conic section lies at x_0 along the line through the planetary dipole that parallels the x axis at x_0 . The focus point, x_0 , the eccentricity, ϵ , and the focal parameter, p (which together with the eccentricity gives the semi-latus rectum, $L = p\epsilon$), are determined by a grid search method that minimizes the RMS misfit. The best-fit parameters to the bow shock midpoints are given by $x_0 = 0.5 R_M$, $\epsilon = 1.04$, and $p = 2.75 R_M$ (Figure 2.5). As mentioned above, bow shock identification is difficult for parallel bow shock conditions, and the outliers in bow shock locations are due to the conservative outer limits chosen for crossings during these conditions. Our best-fit model has an RMS misfit of ~ 1100 km between model boundary and bow shock position. The extrapolated nose distance for this best-fit model is 1.90 R_M , which yields an approximate magnetosheath thickness of 0.45 R_M from our midpoint magnetopause and bow shock fits. This magnetosheath thickness is comparable to that predicted at Mercury by magnetohydrodynamic and hybrid models (e.g., [85, 86]).



Figure 2.5: Midpoints between the inner and outer bow shock crossing positions. Error bars show the distance between the inner and outer crossings. Curves show the best-fit conic section to the data (red) and models from previous studies by Slavin et al. [77] (blue) and Russell [76] (green). Parameters for the best-fit model to the bow shock midpoints are $p = 2.75 R_M$, $\epsilon = 1.04$, and $x_0 = 0.5 R_M$.

2.3.2 Probabilistic fits

We also examined the magnetopause and bow shock positions in a probabilistic manner, by means of a method employed at Jupiter by Joy et al. [87]. As described in Section 2.2, for each crossing we identified an extended region in space within which the magnetopause or bow shock crossings occurred. The inner and outer limits of the boundaries that we identified span the portion of the spacecraft trajectory over which boundary encounters occurred during each pass. The data set therefore reflects locations where encounters with the magnetopause and bow shock boundary were probable, and so we used the crossing data to build a probability density map of these boundaries around the planet. The spacecraft trajectories between each inner and outer crossing limit were registered on spatial grids around the planet, and the number of crossings passing through each grid cell was used to build a probability density map of the region of space in which the magnetopause and bow shock are most likely to be encountered.

We divided space around the planet into grid cells as follows. For the magnetopause, we adopted a spherical coordinate system on the dayside and a cylindrical system on the night to match approximately the shape of the boundary. The dayside was split into cells in $r_{\rm MP}$, $\theta_{\rm MP}$, and $\phi_{\rm MP}$, where $r_{\rm MP} = \sqrt{(x^2 + \rho^2)}$ is the distance from the dipole center, $\theta_{\rm MP} =$ $\cos^{-1}(x/r_{\rm MP})$ is the angle measured from the axis of revolution, and $\phi_{\rm MP}$ is the azimuth about the axis of revolution defined as $\tan^{-1}[y/(z-z_d)]$. Grid cells were spaced every 50 km, 10°, and 30° in $r_{\rm MP}$, $\theta_{\rm MP}$, and $\phi_{\rm MP}$, respectively. The nightside was divided into grid cells spaced every 50 km, 680 km, and 30° in ρ , x, and $\phi_{\rm MP}$, respectively. Because the bow shock nightside data were not well matched by a cylindrical shape, we used the spherical coordinate system, $r_{\rm SK}$, $\theta_{\rm SK}$, $\phi_{\rm SK}$, but with an origin on the axis of revolution at $x = -4 R_M$ for all of the bow shock crossings. That is, $r_{\rm SK} =$ $\sqrt{((x+4R_M)^2+\rho^2)}, \ \theta_{\rm SK} = \cos^{-1}[(x+4R_M)/r_{\rm SK}], \ \text{and} \ \phi_{\rm SK} = \phi_{\rm MP}.$ The bow shock bins were 50 km, 5°, and 30° in $r_{\rm SK}$, $\theta_{\rm SK}$, and $\phi_{\rm SK}$, respectively. These coordinates are used only to bin the data, and all results are shown in $\rho - x$ space.

We evaluated the frequency with which the spacecraft trajectory between the inner and outer magnetopause (or bow shock) crossings passed through each bin. That is, for each orbit, each grid cell received a "hit" for every 1 s measurement point in that bin between the inner and outer limit of the magnetopause (or bow shock). The hits were summed in each cell over all orbits for the magnetopause and bow shock boundaries separately. The cells with the highest number of hits had the highest likelihood of falling between the inner and outer boundary limits. Probabilities were evaluated by dividing the number of counts in each cell by the sum of all the hits in all cells along a predefined direction (e.g., along each x bin on the nightside and along each θ bin on the dayside for the magnetopause). The normalization choice reflects the 100% probability that the boundary occurs at some x or θ position. This procedure resulted in three-dimensional (3-D) probability density maps of both the magnetopause and bow shock around the planet. For the analyses that follow, we assumed that both the bow shock and magnetopause are figures of revolution. We used a two-dimensional (2-D) probability distribution in the $\rho - x$ plane generated by summing the hits in bins with the same r and θ (dayside) or ρ and x (nightside), over all ϕ bins and then normalizing by the total number of counts along r (or ρ) at each θ (or x). The assumption of a figure of revolution for each boundary was tested quantitatively (see Section 2.5).

The 2-D probability density map (Figure 2.6) for the magnetopause shows that this boundary has a maximum probability of occurrence that lies within a narrow band on the dayside and within a more extended region on the nightside. We fit Shue et al. models to this probabilistic boundary by performing a weighted fit to the grid cell locations in $\rho - x$ space, such that each cell location was weighted by its probability. The best-fit curve yielded the same model parameters as the fit to the magnetopause midpoints (Table 2.1), but with a lower RMS misfit of ~ 96 km that reflects the use of the probabilities as weights.

The bow shock probability density map (Figure 2.7) shows a more extended spread in the boundary locations than that for the magnetopause, in particular on the nightside, where the highest-probability regions occur at large ρ values. This spread is the result of outlier crossings (Figure 2.5), most of which occurred during parallel shock conditions. A conic fit to the grid cell locations weighted by the probabilities yields slightly different model parameters than our midpoint fit. The least sensitive of the model parameters in our fits was the focus location, x_0 : varying this parameter from $-0.7 R_M$ to $0.7 R_M$ changed the RMS misfit by only a few percent. Due to the covariance between the parameters, this variation in x_0 was accompanied by large changes in p and ϵ , and the bow shock shape varied from an ellipse to an open hyperbola, but the RMS misfit changed only by a few percent. We thus fixed the focus location to $x_0 = 0.5 R_M$ in order to establish best-fit p and ϵ values that yielded bow shock shapes as close



Figure 2.6: Probability density map of aberrated magnetopause positions. The white space represents regions where the magnetopause is never observed, the blue regions where it is very rarely observed, and the dark red regions where it is observed most often. The red line represents the Shue et al. model that best fits the probability densities, with $R_{ss} = 1.45 R_M$ and $\alpha = 0.5$. The probabilities sum to 1 along each x bin on the nightside and along each θ bin on the dayside.

to hyperboloids as possible. The fit parameters to the bow shock were then given by $x_0 = 0.5 R_M$, $\epsilon = 0.99$, and $p = 3.2 R_M$, with a weighted RMS misfit of ~ 187 km (Table 2.2). With these parameters, the bow shock nose distance is extrapolated to be 2.09 R_M , which gives a magnetosheath width of 0.64 R_M from the probabilistic analysis. The shock distance from the dipole-Sun line in the y - z plane is given by 3.83 R_M at x = 0 and by 6.35 R_M at $x = -4 R_M$.
	M	idpoir	nt fit	Probabilistic fit		
	$R_{\rm ss}$	α	RMS	$R_{\rm ss}$	α	RMS
	$(R_{\rm M})$		residual	$(R_{\rm M})$		residual
			(km)			(km)
All P_{Ram}	1.45	0.50	587	1.45	0.50	96
$P_{\rm Ram}$ corrected	-	-	-	1.45	0.50	94
$< P_{\rm Ram} >= 8.8 \text{ nPa}$	-	-	-	1.55	0.50	96
$< P_{\text{Ram}} >= 11.5 \text{ nPa}$	-	-	-	1.50	0.50	106
$< P_{\text{Ram}} >= 15.2 \text{ nPa}$	-	-	-	1.45	0.50	155
$< P_{\rm Ram} >= 18.8$ nPa	-	-	-	1.40	0.50	120
$< P_{\text{Ram}} >= 21.6 \text{ nPa}$	-	-	-	1.35	0.50	115

2.4. Response of boundaries to solar wind forcing

Table 2.1: Summary of the best-fit Shue et al. model parameters to the magnetopause crossing points under different ram pressure conditions for a fixed α value of 0.50. Angular brackets denote mean value within each ram pressure bin.

2.4 Response of boundaries to solar wind forcing

2.4.1 Magnetopause

The magnetopause shape and location are expected to vary with solar wind and IMF conditions (e.g., [74, 83]). Figure 2.8 shows the solar wind statistics at Mercury at the time of the bow shock crossings, obtained from averages of MESSENGER MAG observations of the IMF and from WSA-Enlil model predictions of the solar wind density, speed, ram pressure, and Alfvén Mach number [42, 78]. We assumed that these statistics held at the times of both the magnetopause and bow shock crossings on a given orbit and used them to examine the response of the boundaries to the IMF and solar wind.

The magnetopause boundary is observed to be closer to the planet during times of increased ram pressure and farther out during times of low ram pressures (Figure 2.9). Baker et al. [78] showed that solar wind ram pressure values from the WSA-Enlil model order and organize the magnetopause stand-off distance. This behavior is expected, because the magnetic pressure of the planet's magnetosphere and the solar wind dynamic pressure



Figure 2.7: Probability density map of the aberrated bow shock crossings. The red line represents the best-fit conic section to the probability densities, with $p = 3.2 R_M$, $\epsilon = 0.99$, and $x_0 = 0.5 R_M$. The probabilities sum to 1 along each x bin on the nightside and along each θ bin on the dayside.

are the two largest contributors to the pressure balance that defines the magnetopause boundary. The ram pressure effect is the dominant factor influencing the magnetopause, because any IMF effects, if present, were not readily apparent in the raw data. We thus investigated any possible effect of the IMF direction on the magnetopause after removing the dependence on $P_{\rm Ram}$.

We assessed changes in the magnetopause shape and position under different ram pressure conditions. We binned the magnetopause crossing data into five P_{Ram} bins, such that each P_{Ram} bin contained one-fifth of the range of P_{Ram} values represented in the data. Thus the P_{Ram} bins did not contain

	Midpoint fit						
	x_0	p	ϵ	RMS	Subsolar		
	$(R_{\rm M})$	$(R_{\rm M})$		residual	distance		
				(km)	$(R_{\rm M})$		
All $M_{\rm A}$	0.5	2.75	1.04	1115	1.90		
	Probabilistic fit						
	x_0	p	ϵ	RMS	Subsolar		
	$(R_{\rm M})$	$(R_{\rm M})$		residual	distance		
				(km)	$(R_{\rm M})$		
All $M_{\rm A}$	0.5	3.20	0.99	187	2.09		
$M_{\rm A}$ corrected	0.5	2.90	1.02	149	1.96		
$< M_{\rm A} >= 4.12$	0.5	3.55	1.02	241	2.29		
$< M_{\rm A} >= 6.32$	0.5	2.95	1.02	195	1.99		
$< M_{\rm A} >= 11.8$	0.5	2.75	1.02	162	1.89		
Low cone angle	0.5	3.10	0.99	169	2.04		
$(\theta < 45^{\circ})$							
High cone angle	0.5	2.95	0.99	157	1.97		
$(\theta > 45^{\circ})$							

2.4. Response of boundaries to solar wind forcing

Table 2.2: Summary of the best-fit conic section parameters to the bow shock crossings under different Mach number conditions. In the fits, x_0 was fixed at a value of 0.5 R_M (see text). In the probabilistic fits to the three different M_A bins, p was the only parameter varied, and ϵ was fixed at the mean value between the best-fit probabilistic and best-fit midpoint result (see text). Angular brackets denote mean M_A value within each Mach number bin.

equal numbers of data points, but the range of P_{Ram} values was the same in each bin. We built 2-D probability density maps for each P_{Ram} bin and conducted fits of the two-parameter Shue et al. model to these separately (Table 2.1). The largest uncertainty in the best-fit model shape is most likely associated with the highest ram pressure bin, which will include the more extreme solar events that may not be properly captured by the WSA-Enlil model. The R_{ss} and α values for the best-fit curves for each of the ram pressure bins are plotted in Figure 2.10.

With the exception of the highest P_{Ram} bin, the R_{ss} parameter decreases overall with P_{Ram} ; in contrast, α shows no systematic variation with P_{Ram} . Thus we first removed the dominant R_{ss} versus P_{Ram} variation from the data



Figure 2.8: Solar wind statistics at the times of the bow shock crossings obtained from WSA-Enlil model predictions of the solar wind and from averages of MESSENGER MAG observations in the interplanetary medium for the IMF. Histograms of (a) solar wind number density (cm^{-3}) , (b) solar wind speed $(km \ s^{-1})$, (c) solar wind ram pressure (nPa), and (d) solar wind Alfvén Mach number. In (d), 15 observations with Mach numbers > 20 are not shown; these numbers reach a maximum value of 69.0.

and normalized the magnetopause crossings by the ram pressure. To do this, we estimated only the R_{ss} parameter for the data in each of the five P_{Ram} bins while keeping α fixed at its average best-fit value of 0.5 (i.e., a fixed magnetopause shape). We checked that the best-fit models for the fixed- α fits in all the five P_{Ram} bins yielded residuals that had Gaussian distributions with a nearly zero mean. The R_{ss} values for the best-fit curves for each



Figure 2.9: Midpoints of the magnetopause crossings color-coded by solar wind ram pressure. The solid black line through these data points is the best-fit Shue et al. model. During times of high solar wind ram pressure, the magnetopause is observed to move closer to the planet, as expected.

of these ram pressure bins are plotted in Figure 2.11 and given in Table 2.1. These were fit by a power law given by $R_{ss} = (2.15 \pm 0.10) P_{\text{Ram}}^{[(-1/6.75)\pm 0.024]}$, where R_{ss} is in units of R_M and P_{Ram} in nPa and the uncertainties are the 95% confidence limits on the exponent and the coefficient obtained from the fitting procedure. The magnitude of the exponent is only slightly less than the -1/6 dependence of R_{ss} on P_{Ram} expected for a simple pressure balance between the internal dipole magnetic field pressure (and small internal plasma pressure) and solar wind dynamic pressure. This result suggests that the effects of induction in Mercury's conductive interior, which "stiffens" the dayside magnetosphere against changes in solar wind pressure [60–63], may be present, but if so it is a secondary effect, at least at the altitudes over



Figure 2.10: (a) Magnetopause subsolar stand-off distance, R_{ss} , plotted as a function of solar wind ram pressure when both parameters were allowed to vary in the Shue et al. model fits. The blue stars represent R_{ss} values established from best-fit curves to five ram-pressure data bins. (b) Flaring parameter, α , corresponding to the R_{ss} values in (a) for the best-fit Shue et al. curves, plotted as a function of solar wind ram pressure.

which MESSENGER samples the magnetopause. A detailed study of induction signatures is beyond the scope of this paper. Refinement of the power-law above will require both additional observations and assessment of uncertainties in the WSA-Enlil model predictions.

The power law relationship above indicates that a ram pressure of 175 nPa would collapse the magnetopause to the planet's surface. We note that the minimum and maximum pressures consistent with the uncertainties of



Figure 2.11: Magnetopause subsolar stand-off distance, R_{ss} , plotted as a function of solar wind ram pressure for fits when the flaring parameter was kept fixed at a value of 0.5. The blue stars represent R_{ss} values established from best-fit curves to five ram-pressure data bins, and the red curve is a power law fit to these values with an exponent of -0.148 ± 0.017 .

our power law fit span a wide range from 65 nPa and 692 nPa. In the simulations conducted by Kabin et al. [88], a ram pressure of 147 nPa was found to collapse the magnetopause to the surface, which is within our uncertainty range. Using a -1/6 power law (i.e., the relation expected in the absence of induction) relative to the mean observed R_{ss} of 1.45 R_M yields a lower pressure for the collapse of the magnetosphere of 133 nPa, well within the range of uncertainty of our power law expression.

With the derived power-law relationship we normalized our magnetopause crossings as follows. From the WSA-Enlil data, we found the corresponding P_{Ram} value for the inbound and outbound portion of each orbit, and from those values we established an associated R_{ss} value for each orbit from the power-law fit to our P_{Ram} -binned data. The ram-pressure-independent magnetopause crossing locations were then determined by multiplying the x, y, and $(z - z_d)$ values by the mean R_{ss} for all crossings (i.e., $R_{ss} = 1.45$ R_M) divided by the R_{ss} value associated with each observation point. Figure 2.12 shows the probability density map of the magnetopause after the solar wind ram pressure dependence was removed. The P_{Ram} -independent magnetopause location is better constrained than the uncorrected locations (Figure 2.6), especially on the dayside, where the zones of high-probability regions lie in a narrow band. Although the model that best fits these data is still described by the same parameters as the P_{Ram} -uncorrected magnetopause (see Table 2.1), the RMS residual is lower than for the P_{Ram} -uncorrected magnetopause that at the 91% significance level the variance of the P_{Ram} -uncorrected magnetopause model is lower than the variance of the P_{Ram} -uncorrected magnetopause model.

After removing the first-order variation of the magnetopause position with ram pressure, we binned the data again into the same P_{Ram} bins as above, fixed R_{ss} to the best-fit P_{Ram} -corrected value of 1.45 R_M and left α to vary in Shue et al. model fits to these bins, to test the influence of ram pressure on the flaring of the magnetopause. Unlike the situation at Earth, there is no increase in flaring of the magnetopause with increased ram pressure. This result is expected if the magnetospheric currents and the ratio of the static solar wind pressure to the dynamic pressure remain constant [75].

We then assessed whether the IMF direction affects the magnetopause shape after the ram pressure dependence was removed. Since the MES-SENGER IMF averages 1 h upstream of the bow shock are not ideal for evaluating IMF B_z affects on the magnetosphere [89], we used the magnetic shear angle across the magnetopause to search for any overall dependence on magnetic reconnection. The magnetic shear was calculated by taking the dot product of the magnetic field unit vector 1 min inside the inner edge of the magnetopause crossing and the unit vector 1 min outside the outer edge of the magnetopause crossing. Figure 2.13 shows the distribution of magnetic shear angles across the magnetopause for our observations. We divided



Figure 2.12: Probability density map of the aberrated magnetopause crossings after the crossing positions were normalized by solar wind ram pressure. The best-fit Shue et al. model (red curve) yields the same R_{ss} and α values as before (Figure 2.6), but the spread in the data has decreased somewhat on the dayside, and the regions of highest probability are more spatially constrained.

the ram pressure corrected magnetopause crossings into low ($\theta < 80^{\circ}$) shear angle and high ($\theta > 100^{\circ}$) shear angle bins, and built probability density maps of these data. Shue et al. model fits to the highest probability regions do not yield a resolvable difference between fits to the magnetopause locations separated by magnetic shear angle. Further observations are needed to increase the signal-to-noise ratio in the boundary locations and potentially resolve any shear angle dependence on the magnetopause boundary.



Figure 2.13: Histograms of (a) the magnetic shear angle across the magnetopause and (b) the IMF cone angle, calculated from MESSENGER magnetic field data.

2.4.2 Bow shock

We used WSA-Enlil model data as well as MESSENGER IMF averages to assess how the bow shock is affected by solar wind and IMF conditions. To first order, the solar wind Alfvén Mach number is the dominant factor affecting the bow shock; the shock is closer to the planet during high M_A than during low M_A (Figure 2.14). This result is expected and in agreement with gas dynamic and magnetohydrodynamic simulations [85, 86, 90]. As M_A increases, the jump of the plasma flow speed transverse to the shock surface also increases, corresponding to higher plasma flow speed around the magnetopause that results in a thinner magnetosheath [91]. At greater downstream distances, the bow shock weakens as it asymptotically approaches



Figure 2.14: Midpoints of the bow shock crossings color-coded by solar wind Alfvén Mach number. The solid black line through the data is the conic section best fit to the midpoints.

its Mach cone [92, 93]. Although MESSENGER does not sample the bow shock at large downstream distances, an enhanced flaring of this surface with decreasing Mach number may be present (Figure 2.14).

To establish the bow shock position normalized by Alfvén Mach number, we adopted a procedure to remove the Mach-number dependence similar to that applied to the magnetopause to remove the dependence on ram pressure. We binned the bow shock crossings into three $M_{\rm A}$ bins and fit separate conic sections to each. To have sufficient data points in each $M_{\rm A}$ bin to perform model fits (see Figure 2.8d), the low- $M_{\rm A}$ bin was assigned $M_{\rm A} < 5$, the medium- $M_{\rm A}$ bin had values in the range $5 < M_{\rm A} < 8$, and the high- $M_{\rm A}$ bin had values of $M_{\rm A} > 8$ (see Table 2.2 for mean values of $M_{\rm A}$ in each bin). Fits in which all three bow shock parameters were varied revealed large trade-offs among the three model parameters, in a manner similar to the magnetopause fits. However, through these fits we were able to establish that the parameter most systematically affected by $M_{\rm A}$ is the focal parameter, p, which steadily decreased with increasing $M_{\rm A}$. The other two parameters did not show any systematic behavior with $M_{\rm A}$. Thus in order to normalize our bow shock boundary by $M_{\rm A}$ and look for higher-order dependencies in the data, we assumed that the bow shock shape does not change (by keeping x_0 and ϵ fixed at their average values) and fit separate conic sections to the bow shock crossings in the three $M_{\rm A}$ bins.

We again checked that these fixed-shape fits still yielded residuals with zero means and nearly Gaussian distributions. In our fits, the low- M_A bin had the highest RMS misfit, and the high- M_A bin had the lowest RMS misfit for all the fits conducted (Table 2.2), consistent with most of the spread in the bow shock location occurring during low M_A (Figure 2.14). From the fits of conic sections to the different M_A bins, we established a power-law relationship between p and M_A , given by $p = (4.79 \pm 2.54)M_A^{(-0.23\pm0.17)}$, where the uncertainties in the exponent and coefficient are the 95% confidence limits determined from the fitting procedure.

We scaled the bow shock crossing positions by p_0/p_i , where p_0 is the mean p value obtained from averaging the p values from the midpoint and probabilistic fit (Table 2.2), and p_i is the p parameter for the ith crossing point determined from the power-law relationship above. The resulting probability density map for the M_A -corrected bow shock positions (Figure 2.15) shows a decrease in the spread of the bow shock locations, as well as a marked decrease in the distance between the high-probability regions on the nightside and the best-fit model boundary. The best-fit model boundary to the M_A -independent bow shock crossings has a p parameter of 2.9 R_M , an eccentricity of 1.02, a focus point of $x_0 = 0.5$, and a minimum RMS misfit 20% less than that for the fit shown in Figure 2.7.

We also assessed the influence of $M_{\rm A}$ on the flaring of the bow shock after removing the first-order dependence of the bow shock position on the Mach number. By allowing ϵ to vary (and fixing $x_0 = 0.5$ and $p = 2.9 R_M$) in our conic sections fit to the different $M_{\rm A}$ bins, we found no statistically



Figure 2.15: Probability density map of the aberrated bow shock crossings after removing the dependence on Alfvén Mach number. The red line represents the best-fit conic section to the probability densities, with $p = 2.9 R_M$, $\epsilon = 1.02$, and $x_0 = 0.5 R_M$. The normalized bow shock is more spatially constrained, in a manner similar to that for the magnetopause normalized by ram pressure.

significant variation of ϵ with M_A . For the sake of completeness, we also conducted similar tests at a variety of other x_0 values, which yielded similar results. We conclude that variation in bow shock flaring with M_A is masked by the high variability of the crossing locations. More bow shock crossings are needed at high M_A values to establish whether ϵ varies systematically with M_A at Mercury.

The IMF cone angle is also expected to affect the bow shock, as the shock is anticipated to flare during quasi-parallel shock conditions [94]. We

2.5. Discussion

evaluated whether any dependence on IMF cone angle is observed in the bow shock location or shape after the Mach number dependence was removed. The IMF cone angle is given by $\theta = \cos^{-1} \left(\frac{B_{|\mathbf{x}|}}{B_{\text{Total}}}\right)$, and its distribution is shown in Figure 2.13. By dividing the data into bins of low ($\theta < 45^{\circ}$) cone angle and high ($\theta > 45^{\circ}$) cone angle, we found that there is no resolvable difference between conic section fits to the bow shock locations separated by cone angle (Table 2.2). The shock is more spatially spread out for the low-cone-angle bin, an expected result because a low cone angle signifies parallel shock conditions. The flaring of the bow shock does not appear to be affected by IMF cone angle.

2.5 Discussion

The observations of Mercury's magnetopause and bow shock presented here indicate that these boundaries are variable and dynamic. At Mercury, the solar wind ram pressure and the Alfvén Mach number are the two dominant external influences on the boundaries. The magnetopause is observed to move planetward during high P_{Ram} , and similarly the bow shock moves planetward during times of high M_{A} . Unlike at Earth, increased P_{Ram} does not increase the flaring of the magnetopause; the shape of the boundary is unchanged under variations in solar wind ram pressure. In a like manner, M_{A} does not appear to influence the flaring of the bow shock at Mercury despite the fact that at Earth the shock is more flared during times of low M_{A} .

An important result is that the average magnetopause becomes cylindrical at relatively small downstream distances of only $\sim 2-3 R_{\rm M}$ (Figure 2.3). At Earth, in contrast, the magnetotail does not cease flaring until a downstream distance of $\sim 100 R_{\rm E}$ [95]. Expressed in terms of subsolar magnetopause stand-off distances, the downstream flaring of Mercury's tail ceases by $\sim 2 R_{\rm ss}$ whereas at Earth this effect is not seen until $\sim 10 R_{\rm ss}$. The factors determining the location where tail flaring ceases are not well understood, but the position likely corresponds to the distance at which the plasma sheet is disconnected from the planet as a result of reconnection, and

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plasma flow in the tail is all anti-sunward. At Earth this position occurs at about $-100 R_{\rm E}$ [96], whereas for Mercury Slavin et al. [59] estimated a downstream distance of $\sim 3~R_{\rm M}$ from MESSENGER flyby observations of reconnection in the tail. Mercury's magnetopause is well fit by the Shue et al. model with $\alpha = 0.5$, which defines the transition from an open ($\alpha > 0.5$) to a closed ($\alpha < 0.5$) magnetospheric cavity on the nightside. From the best-fit Shue et al. model for the ram-pressure-corrected magnetopause we find that the magnetotail is on average nearly cylindrical with a radius of 2.05 $R_{\rm M}$ at the dawn-dusk terminator and 2.77 $R_{\rm M}$ at a distance of 4 $R_{\rm M}$ down the tail. In comparison, Earth's magnetopause is more flared, with the data fit well by Shue et al. models that have $\alpha > 0.5$ [69]. The overall nearly cylindrical shape of Mercury's magnetopause in comparison to Earth's may imply that magnetic flux has a short residence time in the tail, and thus return convection of flux from the tail to the dayside proceeds rapidly, consistent with the expected ~ 2 minute Dungey cycle period at Mercury.

The analysis of the magnetopause and bow shock boundaries in $\rho - x$ space was based on the assumption of rotational symmetry about the dipole-Sun line. We tested whether this assumption was justified both qualitatively and quantitatively. First, the boundary crossings plotted in x - y and $x - (z - z_d)$ space did not reveal systematic differences. The ram pressure and Mach number are not observed to cause any asymmetries in the shape, as the boundary crossings corrected for P_{Ram} and M_{A} have similar shapes to the uncorrected crossings in x - y and $x - (z - z_d)$ space. These comparisons indicate that variations from rotational symmetry can be treated as a perturbation to the figures of revolution. We then assessed the degree to which systematic deviations from figures of revolution are present. From the corrected crossing locations, we evaluated the best-fit models for both boundaries, and we calculated the perpendicular distances of each crossing from the model boundaries as a function of the azimuthal angle, i.e., $\phi_{\rm MP}$ or $\phi_{\rm SK}$. If either the magnetopause or bow shock were flattened or elongated in the north-south direction, such an effect would be evident as a sinusoidal variation in $\phi_{\rm MP}$ or $\phi_{\rm SK}$, respectively, relative to the mean boundary at

2.6. Conclusions

a period of 180°, i.e., two cycles. We did not observe any systematic departures from figures of revolution for either the magnetopause or the bow shock. By binning the deviations into 2° bins in $\phi_{\rm MP}$ or $\phi_{\rm SK}$ and fitting a sinusoid to these binned deviations, we find maximum sine-wave amplitudes of 62 and 3 km for the magnetopause and bow shock, respectively. These sine-wave amplitudes are more than an order of magnitude smaller than the variability in the deviations of the two boundaries about the models. We conclude that, to first order, the boundaries are figures of revolution. The scatter relative to the mean is high however, implying that dynamic variability in Mercury's magnetosphere and bow shock are large and that second-order structure could be present but masked by the large dynamic variability. It is thus possible that instantaneously these boundaries are not figures of revolution. However, on average, the departures from figures of revolution are small compared with the dynamics in the system. Because of the high variability of the crossing locations analyzed so far, we cannot yet resolve any average asymmetries in the boundary shapes. At Earth, the maximum departure from a figure of revolution at high latitudes is ~ 1 $R_{\rm E}$ [72], which corresponds to ~ 0.1 $R_{\rm ss}$ (subsolar stand-off distance is ~ 10 $R_{\rm E}$). In comparison, an equivalent 0.1 $R_{\rm ss}$ departure is only ~ 350 km at Mercury, which would be a sufficiently small departure to be easily masked by the variability in the available data. Such signatures may be resolved with additional observations.

2.6 Conclusions

We have established Mercury's time-averaged magnetopause and bow shock location and shape from MESSENGER Magnetometer data obtained during three Mercury years in orbit. We find that the magnetopause is well described by a Shue et al. model parameterized by a subsolar standoff distance of 1.45 $R_{\rm M}$ and a flaring parameter of $\alpha = 0.5$. The solar wind ram pressure exerts a primary control on magnetopause location; the boundary moves closer to the planet under higher $P_{\rm Ram}$ (giving a subsolar distance of 1.35 $R_{\rm M}$ for a mean $P_{\rm Ram}$ of 21.6 nPa) and farther away from the planet under lower P_{Ram} (with an R_{ss} of 1.55 R_{M} for a mean P_{Ram} of 8.8 nPa), while leaving the shape unchanged (Table 2.1). The paraboloid model of Belenkaya et al. [81] provides a substantially worse overall fit to the magnetopause crossings than the Shue et al. model, reflecting the absence of evidence for substantial flaring from observations on the nightside. This comparison suggests that future improvements in global models for Mercury's magnetosphere should use a ram-pressure-corrected Shue et al. model magnetopause. The observed low flaring of the magnetotail may imply that magnetic flux has a short residence time in the tail on average. This short residence time of the tail flux could also imply that return convection of flux from the tail to the dayside proceeds rapidly.

The shape of Mercury's bow shock corrected for Alfvén Mach number is that of a hyperboloid given by the parameters $x_0 = 0.5 R_M$, $p = 2.9 R_M$, and $\epsilon = 1.02$, and a subsolar stand-off distance of 1.96 R_M . The bow shock shape does not appear to vary with Alfvén Mach number, as there is no change in flaring. This is an unexpected result, since Earth's bow shock is observed to flare with decreasing Alfvén Mach number. At Mercury, the bow shock moves closer to the planet for high M_A and farther out for low M_A ; the extrapolated nose distance of the shock is at 1.89 R_M for a mean M_A of 11.8, and at 2.29 R_M for a mean M_A of 4.12 (Table 2.2). Both the magnetopause and bow shock boundaries are figures of revolution to first order, but the variability about the mean is large. With the current data available we do not resolve effects of IMF orientation on the magnetopause or the bow shock. As more data are acquired by MESSENGER, effects of IMF on the magnetopause or bow shock should be more readily discernible.

The variation of the bow shock and magnetopause location with dynamics is large at Mercury, and understanding the processes that these dynamics reflect is a key area of future study. The derivation of the average boundaries presented here provides a baseline with which to evaluate excursions of the system from its average state. Extensive progress has already been made in understanding boundary waves [97, 98] and reconnection at Mercury [34, 59, 99]. Extending that work to better understand those aspects of global magnetospheric dynamics that could lead to the large variations in boundary locations is a fruitful area of inquiry.

Chapter 3

Observations of Mercury's northern cusp region with MESSENGER's Magnetometer¹

The magnetic cusp of a planetary magnetosphere allows solar wind plasma to gain access to the planet's magnetosphere and, for Mercury, the surface. From measurements by the MESSENGER Magnetometer we have characterized the magnetic field in the northern cusp region of Mercury [89]. The first six months of orbital measurements indicate a mean latitudinal extent of the cusp of $\sim 11^{\circ}$, and a mean local time extent of 4.5 hrs, at spacecraft altitudes. From the average magnetic pressure deficit in the cusp, we estimate that $(1.1\pm0.6)\times10^{24}$ protons s⁻¹ bombard the surface over an area of $(5.2 \pm 1.6)10^{11}$ m² near the northern cusp. Plasma pressures in the cusp are 40% higher when the interplanetary magnetic field (IMF) is anti-sunward than when it is sunward. The influence of the IMF direction does not overcome the north-south asymmetry of Mercury's internal field, and particle flux to the surface near the southern cusp is predicted to be a factor of 4 greater than in the north. The higher particle flux impacting the surface in the south should lead to a greater exospheric source from the south and a higher rate of space weathering than in the area of the northern cusp.

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3.1 Introduction

The magnetospheres of planets with dipolar internal fields possess magnetic cusps, regions near the magnetic poles at which fields from magnetopause currents nearly cancel the internal field. For vacuum superposition of the magnetic fields of the dipole and magnetopause currents, the cusps are topological singularities where the magnetic field vanishes. The weak field near the cusp allows the shocked solar wind plasma of the magnetosheath ready access to the magnetosphere, and the magnetic field lines that thread the cusp are populated with this plasma.

Mercury's internal field is symmetric about the rotation axis but asymmetric about the geographic equator and can be represented by a dipole with a moment of 195 nT $R_{\rm M}^3$ (where $R_{\rm M}$ is Mercury's mean radius, 2440 km) offset 0.2 $R_{\rm M}$ northward from the planetary center [17]. The high-latitude field at the surface is predicted to be 4 times weaker in the southern hemisphere than in the northern hemisphere, leading to a correspondingly greater spatial extent of the cusp projection to the surface in the south than in the north. Solar wind sputtering of species from the planetary surface may be a substantial source of exospheric particles (e.g., [100]). It is not known whether asymmetric particle bombardment of the surface and corresponding differences in space weathering rates could have produced detectable hemispheric differences in surface color or reflectance.

Earth's cusps have been extensively studied at low and high altitudes (e.g., [101]). The position and size of the cusp areas at Earth depend on the solar wind pressure [102, 103] and the IMF (e.g., [100, 104–106]). Mercury lacks an ionosphere, the magnetosphere is a factor of ~ 8 smaller than Earth's relative to the planetary diameter, and the average solar wind density is an order of magnitude higher than at Earth [23], so the cusps at Mercury may be quite different from those at Earth. In addition, because the IMF is dominantly sunward or anti-sunward at Mercury, the IMF component in the Sun-Mercury direction may play a prominent role in the dynamics of Mercury's cusps (e.g., [107]).

Previous work on Mercury's cusps focused on magnetosphere-solar wind

interaction by means of analytic models [108–110], global magnetohydrodynamic models [88, 111], hybrid simulations [112–114], or semi-empirical models [100, 115]. These studies indicated that solar wind ions can reach the surface in the cusp region, but the spatial extent of the cusp and the particle fluence vary among the models. This variation is partly due to the different IMF and solar wind conditions assumed. From solar wind and IMF conditions at Mercury's aphelion and perihelion, Sarantos et al. [110] predicted that the largest flux of precipitating solar wind ions impacting Mercury's surface occurs at local noon between 40° and 60° latitude with an equatorward shift at perihelion. To date there have been no observations that quantify the total plasma pressure in Mercury's cusps or provide a basis for assessing its sensitivity to the sunward IMF component.

Orbital observations by the MESSENGER spacecraft's Fast Imaging Plasma Spectrometer (FIPS) have revealed that the flux of heavy ions in Mercury's magnetosphere peaks between 65° and 75° latitude, consistent with the predicted location of the northern magnetic cusp [22]. The ion flux peaks coincide with depressions in magnetic field strength [116] measured with the MESSENGER Magnetometer (MAG) [49]. In this chapter we characterize the northern cusp with MAG data from six months of orbital observations, calculate the corresponding surface precipitation, and investigate the influence of the sunward IMF and solar wind pressure on the mean cusp plasma pressure.

3.2 Observations

The MESSENGER spacecraft was inserted into orbit about Mercury on 18 March 2011. The initial orbit had a 200 km periapsis altitude, 82.5° inclination, 15,300 km apoapsis altitude, and 12 hour period. We use two Mercury years of MAG data starting from 23 March 2011, providing coverage at all local times. Data were analyzed in Mercury solar orbital (MSO) coordinates, for which +x is sunward, +z is northward, and +y completes the right-handed system.

The cusp was identified from depressions in the magnitude of 1 s av-



Figure 3.1: Example of a cusp observation on 21 August 2011, orbit 313. (top) Measured (black) and modeled (red) magnetic field magnitude in the cusp region. (middle) Magnetic depression in the residual $|\mathbf{B}|$ (black), residual data before and after cusp entry (red), and a third-degree polynomial fit (blue) to the red curve. (bottom) The calculated pressure deficit ($-P_{\text{Plasma}}$).

eraged total-field data from which a model field had been subtracted. The model incorporates the offset internal dipole field and the magnetopause and tail fields of the Alexeev et al. [82] paraboloid magnetospheric model, with

3.3. Data analysis

model parameters given in Anderson et al. [17]. An aberration correction was calculated from Mercury's orbital speed and a mean solar wind speed of 405 km s⁻¹. For each orbit exhibiting a dayside magnetic depression poleward of the magnetopause, we identified the times of the cusp outer and inner entry and exit points. Transits in the cusp were indicated by sustained depressions in the magnitude of the magnetic field **B** that exceeded typical variability and lasted several minutes. An outer cusp entry was identified at the point where the first transient decrease in $|\mathbf{B}|$ was seen, and the inner entry was picked where the sustained depression in $|\mathbf{B}|$ started. Similar criteria were used for the exit inner and outer points. Figure 3.1 shows the dayside depression in the total residual between the observed and model fields, given by $|\mathbf{B}|_{\text{res}} = |\mathbf{B}|_{\text{obs}} - |\mathbf{B}|_{\text{model}}$. The orbit does not always intersect the cusp, particularly when periapsis is on the nightside, and magnetic depressions were seen on 169 of the 279 orbits analyzed. Each entry and exit time and the aberrated MSO spacecraft positions are given in the Supplemental Materials of [89]. The field depressions were generally associated with enhanced magnetic fluctuations at 1-10 Hz frequency, consistent with greater intensities of local plasma instabilities. The proton gyrofrequency is 2 to 6 Hz for field strengths observed in the cusp (150 to 400 nT). The cusp entry and exit times changed by less than a few seconds for different magnetospheric model parameters.

3.3 Data analysis

We conducted superposed epoch analyses (SEA) of $|\mathbf{B}|$ and $|\mathbf{B}|_{\text{res}}$ in the cusp to derive an average magnetic depression signature (Figure 3.2). Individual profiles from different orbits were aligned in time on their respective cusp interval midpoints and averaged over a time span of six minutes on either side of this midpoint. We also conducted SEA of the 1-10 Hz fluctuations. The fluctuation intensity was evaluated from the 20 sample/s data by taking the root mean square (RMS) value over 1 s intervals in the direction parallel to and two components perpendicular to the 1 s averaged field direction, denoted by δB_{\parallel} , $\delta B_{\perp 1}$ and $\delta B_{\perp 2}$, respectively. We define $\delta B_{\perp} = \sqrt{(\delta B_{\perp 1}^2 + \delta B_{\perp 2}^2)/2}$, so that $\delta B_{\perp} = \delta B_{\parallel}$ if the fluctuations are equal in all components. These analyses confirm the depression in the magnetic field over the cusp and show that this signature is accompanied by an increase in the magnetic fluctuations. The ratio $\delta B_{\perp}/\delta B_{\parallel}$ is about 1.5 in the cusp and higher on either side of the cusp (Figure 3.2), indicating that although the fluctuations in the cusp are transverse, they are less so than the adjacent lower-amplitude fluctuations.

We calculated a plasma pressure that balances the magnetic field depression from $P_{\text{Total}} = P_{\text{Mag}} + P_{\text{Plasma}}$, where P_{Total} is the total pressure; P_{Mag} is the magnetic pressure, $B^2/(2\mu_0)$, where μ_0 is the magnetic permeability; and P_{Plasma} is the particle thermal pressure. We estimated P_{Total} from the magnetic field removed from the cusp magnetic field depression [116]. The unperturbed magnetic field was determined for each pass from the magnetospheric model field and a third-degree polynomial fit to the residuals one minute before and after but excluding the depression interval (Figure 3.1, middle panel). The boundaries of the depression intervals were taken as the average of the inner and outer cusp entry or exit times. In some cases the polynomial fit did not consistently remain above the residual field magnitude in the cusp. These fits were rejected, and new fits were obtained by increasing the time interval for the baseline fit. The polynomial fit was added to the magnetospheric model field to estimate the unperturbed total magnetic field, $B_{\rm U}$. We then evaluated $P_{\rm Total} = B_{\rm U}^2/(2\mu_0)$ and the magnetic pressure deficit, $P_{\rm B-deficit} = P_{\rm Mag} - P_{\rm Total} = -P_{\rm Plasma}$ (see Figure 3.1, bottom panel). This latter quantity gives the additional plasma pressure in the cusp relative to any background plasma pressure in the magnetosphere. In general, FIPS data do not show substantial proton counts adjacent to, but outside, the cusp, indicating that the background plasma pressure near the cusp is much lower than that in the cusp.

The limits of the northern cusp are 55.8° and 83.6° MSO latitude and 7.2 h and 15.9 h local time. On average the cusp is approximately symmetric about noon (Figure 3.3). Since the MESSENGER orbit is eccentric and periapsis is on the descending latitude portion of the orbit, the cusp is encountered at lower altitudes on the descending than on the ascending orbit



Figure 3.2: Superposed epoch analysis of all cusp observations. (top) SEA of observed $|\mathbf{B}|$ (blue) and model $|\mathbf{B}|$ (red) indicated by the scale on the lefthand ordinate, and SEA of residual $|\mathbf{B}|$ (green) indicated by the scale on the right-hand ordinate, for all 169 cusp profiles. (bottom) SEA of RMS 1-10 Hz fluctuations perpendicular and parallel to the local field, δB_{\perp} (blue), and δB_{\parallel} (black) (scale on left). The red curve shows $\delta B_{\perp}/\delta B_{\parallel}$ (scale on right).

track. At higher altitude the cusp is on average a few degrees equatorward of that seen at lower altitude. In the magnetosphere model, the magnetic field at the magnetopause vanishes near 62° N at noon, consistent with the expected shift in cusp latitude closer to the magnetopause.



Figure 3.3: Stereographic projections of the pressure deficit $(-P_{\text{Plasma}})$ along each cusp profile in aberrated MSO coordinates. During portions of MES-SENGER's first Mercury year in orbit (MSO1), the Magnetometer was off when the spacecraft experienced long eclipses or was close to the planet, resulting in the gap in data coverage (between ~ 10 h and ~ 12 h in local time) for the descending tracks. Complete coverage was obtained during MESSEN-GER's second Mercury year in orbit (MSO2). Projections span local times from 6.67 h to 17.3 h and latitudes 55°N to the pole. The color bar is saturated so that observed, but localized, pressure deficits greater in magnitude than -3 nPa are shown in red.

3.4 Discussion

The observations indicate that Mercury's northern cusp region is a persistent but dynamic feature. Not only is the cusp pressure deficit variable on a given pass (Figure 3.1), but the cusp extent and plasma pressure can vary markedly from one orbit to the next (Figure 3.3). This variability likely

3.4. Discussion

results from the influence of different IMF and solar wind conditions and the corresponding interactions with, and dynamics of, Mercury's magnetosphere. Here we focus on establishing the mean cusp pressure and particle fluence to the surface since the plasma pressure may have important consequences for exospheric processes and space weathering. We use MESSEN-GER averages of IMF B_X and predictions of the solar wind ram pressure from the WSA-Enlil solar wind model [42]. Statistics of these quantities for the cusp transits are given in Table 3.1.

MESSENGER's 12-hour eccentric orbit presents challenges to analyzing the effects of the solar wind on the cusp. First, the local time extent of the cusp is sampled only twice each Mercury year (Figure 3.3). A study of variations in cusp local time extent with solar wind conditions will require considerably more observations than are presently available. Second, IMF conditions for a given orbit are estimated from averages of MAG observations upstream of the bow shock. The 1-h time spans for these averages are comparable to the typical time between MESSENGER cusp transits and residence in the solar wind. Only the IMF x-component is generally larger in magnitude than its variability (Table 3.1) implying that the IMF averages are appropriate for investigating the effects of IMF sector structure on the cusp, but not for assessing the influence of magnetopause reconnection. Reconnection depends strongly on the sign of B_z and the magnetosphere responds to changes in the sign of B_z within minutes [24]. Because the mean value of B_z is generally smaller than its variability over the averaging intervals, the average IMF is not a good indicator of reconnection dynamics during our cusp transits. We also compared the average IMF before and after each magnetosphere transit to assess IMF stability. Only 18%of orbits that pass through the cusp exhibit an average IMF B_z that is of the same sign before and after MESSENGER's magnetosphere transit and that is at least one standard deviation different from zero. In contrast, the corresponding percentages for B_x and B_y are 79% and 42%, respectively. The data set therefore allows reliable assessment of the cusp dependence on IMF B_x , which at Mercury's orbit is the dominant IMF component and is predicted to have a strong influence on pressures in the cusp (e.g., [107]).

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Ascending tracks					
	IMF $B_{\rm x} > 0$	IMF $B_{\rm x} < 0$			
$ \mathbf{B} $	193 nT	179 nT			
P_{Plasma}	1.5 ± 0.1 nPa	2.1 ± 0.1 nPa			
IMF $B_{\rm x}$	$16.4 \pm 1.1 \text{ nT} (\sigma = 5.0 \text{ nT})$	$-9.1\pm0.7~\mathrm{nT}~(\sigma=4.7~\mathrm{nT})$			
IMF $B_{\rm y}$	-4.6 ± 1.2 nT (σ = 6.0 nT)	$5.9 \pm 1.1 \ \mathrm{nT} \ (\sigma = 5.9 \ \mathrm{nT})$			
IMF $B_{\rm z}$	$1.9 \pm 1.1 ~\mathrm{nT} ~(\sigma = 6.3 ~\mathrm{nT})$	$1.3 \pm 0.9 \text{ nT} (\sigma = 6.8 \text{ nT})$			
P_{Ram}	11.9 ± 0.4 n Pa	11.3 ± 0.4 nPa			
# of orbits	45	43			
Latitude	70.9 ° N	$71.4 \degree$ N			
Altitude	$684 \mathrm{~km}$	$706 \mathrm{km}$			
Descending tracks					
	IMF $B_{\rm x} > 0$	IMF $B_{\rm x} < 0$			
$ \mathbf{B} $	305 nT	318 nT			
P_{Plasma}	1.8 ± 0.1 nPa	1.9 ± 0.1 nPa			
$B_{\mathbf{x}}$	17.8 ± 0.9 nT (σ = 5.2 nT)	$-10.7 \pm 2.0 \text{ nT} (\sigma = 5.0 \text{ nT})$			
IMF $B_{\rm y}$	-5.1 ± 1.1 nT (σ = 7.0 nT)	$2.5 \pm 1.4 \text{ nT} (\sigma = 6.7 \text{ nT})$			
IMF $B_{\rm z}$	$1.5\pm0.8~\mathrm{nT}~(\sigma=7.6~\mathrm{nT})$	-0.02± 1.35 nT (σ = 7.0 nT)			
P_{Ram}	$14.5\pm0.5~\mathrm{nPa}$	11.4 ± 0.3 n Pa			
# of orbits	55	26			
Latitude	72.5° N	75.2° N			
Altitude	$426 \mathrm{~km}$	404 km			

Table 3.1: Average cusp properties and ambient conditions separated by ascending/descending tracks and by the sign of IMF B_x . The mean $|\mathbf{B}|$ and P_{Plasma} inferred from decreased magnetic field strength are evaluated in the cusp. The average IMF B_x , B_y , and B_z are calculated from observations before and after each magnetospheric transit, and the solar wind P_{Ram} is from the WSA-Enlil model evaluated during times of passage in the cusp. Mean cusp latitude and altitude are weighted by P_{Plasma} ; the rest of the values are unweighted. Uncertainties are 1 standard error of the mean. For IMF averages, the mean standard deviation (σ) is given in parentheses; only the average B_x is consistently greater in magnitude than its variability.

Determining the variability in cusp location and pressure due to dynamics associated with magnetopause reconnection, as indicated by the IMF B_y and B_z components, is left for future analyses.

We assessed the influence of the IMF B_x and the solar wind ram pressure

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on the cusp plasma pressure as follows. Statistics were evaluated separately for ascending and descending passes and for positive and negative IMF B_x , because an anti-sunward IMF (negative B_x) is expected to facilitate plasma transport into the northern cusp (e.g., [107]). The ascending tracks were divided approximately equally between positive and negative B_x , but the magnitude of B_x was ~ 1.7 times higher for sunward than for anti-sunward IMF conditions. The results indicate a larger plasma pressure in the cusp for negative B_x . The high-altitude datasets for $B_x > 0$ and $B_x < 0$ have similar mean altitudes and mean ram pressure (P_{Ram}) values, so the 40% deeper magnetic pressure deficit for negative than positive B_x can be attributed to the IMF orientation. Variation in the cusp mean position at Mercury for the different signs of B_x and comparable average P_{Ram} , as on the ascending tracks, is at most ~ 0.5° compared to the average cusp extent of ~ 11°.

We performed SEA on $P_{\rm B-deficit}$ for each of the two ascending-track populations following the procedure described in Section 3.3, except that outside the depression interval we padded the $P_{\rm B-deficit}$ values with zeroes to fill out the time to 8 min for each event (Figure 3.1, bottom panel). The SEA profiles of $P_{\rm B-deficit}$ (Figure 3.4) confirm that the magnetic pressure deficit in the cusp for the transits with $B_{\rm x} < 0$ is, on average, larger in magnitude than for $B_{\rm x} > 0$ and show that the cusp is present regardless of the sign of $B_{\rm x}$. The sunward/anti-sunward direction of the IMF thus modulates plasma pressures but is not the dominant factor determining pressure in the cusp.

The plasma pressure in the cusp appears to increase with increasing solar wind ram pressure. For $B_x > 0$ the descending tracks exhibit a higher P_{Ram} and also a lower $P_{\text{B-deficit}}$ than the ascending tracks, indicating that the cusp pressures increase with increasing P_{Ram} . The influence of P_{Ram} may account for the smaller difference in $P_{\text{B-deficit}}$ between the descending track observations for $B_x > 0$ and $B_x < 0$, as the mean P_{Ram} is substantially higher for the $B_x > 0$ events. Presumably, $P_{\text{B-deficit}}$ for the $B_x > 0$ descending track cases would have been smaller in magnitude had P_{Ram} for these tracks been comparable to that for the $B_x < 0$ descending track observations.



Figure 3.4: SEA of the magnetic pressure deficit for ascending tracks grouped by IMF $B_x > 0$ (red) and IMF $B_x < 0$ (black). A larger-amplitude magnetic pressure deficit is observed for orbits when the IMF has a negative, or antisunward, B_x component.

The estimate of cusp plasma pressures allows us to calculate the average particle flux bombarding the surface. We describe this calculation and the results in Section 3.4.1. Obtaining an average particle flux to the surface has significant implications for studies relating to the generation of Mercury's exosphere as well as space weathering of the surface.

3.4.1 Surface flux calculation

To calculate the surface flux of particles, we assume an isotropic particle distribution entering from the well-mixed magnetosheath plasma and that the dominant ions in the solar wind are protons. We know the plasma pressure at the spacecraft altitude, so our aim is to derive an expression for the plasma pressure ratio between the spacecraft altitude and the surface, and from there to estimate the pressure at the surface. Once we know the pressure at the surface, we can calculate the flux at the surface by using an equation that relates the flux of isotropic gas particles to their pressure.

We consider a particle distribution function $f_0(\mathbf{v})$ at the spacecraft altitude. The pressure at this altitude is given by the second velocity moment of the distribution function

$$P_0 = \int v^2 f_0(\mathbf{v}) d^3 v.$$
 (3.1)

In spherical coordinates, $f_0(\mathbf{v}) = f_0(v, \alpha, \phi)$ and $d^3v = v^2 \sin(\alpha)d\phi d\alpha dv$, where ϕ is the velocity azimuth about the magnetic field, and α is the velocity pitch angle (the angle between \mathbf{v} and \mathbf{B}). We assume that the distribution is gyrotropic, so that there is no dependence on ϕ , and we also assume that it is isotropic, so that the speed distribution is independent of α , to get $f_0(\mathbf{v}) = 2\pi V(v)A(\alpha)$. In the absence of electric fields parallel to \mathbf{B} , the particle speed distribution (V(v)) is independent of B, so that it is the same everywhere along the field line. We can then rewrite equation (3.1) in spherical coordinates as

$$P_0 = 2\pi \int v^4 V_0(v) dv \int A_0(\alpha) \sin(\alpha) d\alpha.$$
(3.2)

The pressure can be decomposed into the pressure due to particle motions parallel to and perpendicular to the magnetic field, **B**, since $v^2 = v_{\parallel}^2 + v_{\perp}^2$, where $v_{\parallel} = v \cos(\alpha)$ and $v_{\perp} = v \sin(\alpha)$. Then, using equation (3.1) for simplicity,

$$P_0 = \int v_{\parallel}^2 f_0(\mathbf{v}) d^3 v + \int v_{\perp}^2 f_0(\mathbf{v}) d^3 v, \qquad (3.3)$$

yielding $P_0 = P_{\parallel} + P_{\perp}$.

Since the magnetic pressure deficit reflects the perpendicular particle pressure, at the spacecraft altitude we only need P_{\perp} (as this is all we measure), which can be written as

$$P_{\perp 0} = 2\pi \int v^4 V_0(v) dv \int A_0(\alpha) \sin^3(\alpha) d\alpha, \qquad (3.4)$$

after substituting in for v_{\perp} . The total particle pressure at the surface, P_s is

given by

$$P_s = 2\pi \int v^4 V_s(v) dv \int A_s(\alpha) \sin(\alpha) d\alpha.$$
(3.5)

If we can estimate the ratio of $P_s/P_{\perp 0}$, we can calculate P_s because from our measurements of the plasma pressure at the spacecraft altitude we know $P_{\perp 0}$. Thus, because we assumed that there are no parallel electric fields, $V_s(v) = V_0(v)$, the pressure ratio between the two altitudes is

$$\frac{P_s}{P_{\perp 0}} = \frac{\int A_s(\alpha) \sin(\alpha) d\alpha}{\int A_0(\alpha) \sin^3(\alpha) d\alpha}.$$
(3.6)

From Liouville's theorem, in the absence of collisions and loss of particles, df/dt = 0, so that $f_0(\mathbf{v}_0) = f_1(\mathbf{v}_1)$. However, the planet's surface represents a loss for the particles, and we can estimate an upper limit for this loss. For a charged particle, the local magnetic field and instantaneous pitch angle are related by the first adiabatic invariant, the magnetic moment, which yields that

$$\frac{\sin^2(\alpha)}{B} = \text{constant}, \tag{3.7}$$

so that the pitch angle of particles that mirror at the surface is given by

$$\alpha_{m,i} = \sin^{-1}(\sqrt{B_i/B_s}),\tag{3.8}$$

where B_i is the magnetic field magnitude at the altitude of interest, which in this case can be either at the spacecraft or the surface. For more details on magnetic mirroring, see Section 1.2.4. This means that all particles heading toward the surface with pitch angles smaller than $\alpha_{m,i}$ will encounter the surface, while particles with pitch angles larger than $\alpha_{m,i}$ will mirror above the surface and return moving away from the surface. All particles that mirror at or below the surface are missing from the upward flowing portion of $f(\mathbf{v})$. Thus we can rewrite $A_i(\alpha)$, where *i* can represent either '0' or 's' for spacecraft altitude or surface, in terms of the function in the absence of surface losses, A, which is independent of α and altitude. We get that

$$\alpha = 0 \text{ to } \pi - \alpha_{m,i} : \quad A_i(\alpha) = A$$

$$\alpha = \pi - \alpha_{m,i} \text{ to } \pi : \quad A_i(\alpha) = 0, \qquad (3.9)$$

where $\alpha_{m,i}$ represents the mirroring pitch angle at a particular altitude. So this basically means that due to Liouville's theorem, we can assume that $A(\alpha)$ is constant for the part of the pitch angle distribution that is not lost, and for the part that is lost, $A(\alpha)$ is zero by definition. Since the parts of the particle distribution that are lost due to absorption by the surface are not going to contribute to the pressure, the pressure integrals should only go from 0 to $\pi - \alpha_{m,i}$. We can then rewrite equation (3.6) as

$$\frac{P_s}{P_{\perp 0}} = \frac{\int_0^{\pi - \alpha_{m,s}} \sin(\alpha) d\alpha}{\int_0^{\pi - \alpha_{m,0}} \sin^3(\alpha) d\alpha}.$$
(3.10)

For the top integral, $\alpha_{m,s}$ is simple, since $B_i = B_s$ in that case, and this yields from equation (3.8) that $\alpha_{m,s} = \pi/2$. The integral is then straightforward, and a value of 1 is obtained for the top integral. For the bottom integral, $\alpha_{m,0}$ can be calculated from equation (3.8) and an average field magnitude at the spacecraft altitude of the ascending tracks. With a value for $B_0 = 186$ nT, the pressure ratio becomes $P_s/P_{\perp 0} = 0.77$.

We can estimate a lower limit on this pressure ratio by assuming that there is no surface loss, that is, particles are perfectly reflected. In this case, the integral for P_s goes again from $0 < \alpha < \pi/2$, while the integral for $P_{\perp 0}$ goes from $0 < \alpha < \pi$, yielding a pressure ratio of 0.75. From the mean magnetic pressure deficit along the high-altitude ascending tracks, we have that $P_{\perp 0} = 1.79$ nPa. We can use this, along with the two estimates for the pressure ratio, to estimate an upper limit for P_s of 1.38 and a lower limit of 1.34 nPa, with an average of $P_s = 1.36 \pm 0.23$ nPa, where the uncertainty corresponds to the average pressure ratio times half the difference in mean pressures for positive and negative IMF B_x . For an isotropic gas, the flux of particles through a surface is given by

$$\Phi = P/\sqrt{2\pi mkT},\tag{3.11}$$

where m is the particle mass, T is the temperature, and k is the Boltzmann's constant. To estimate the flux of particles to the surface, we double the surface pressure that we estimated above because that only accounts for the downgoing half of the distribution (as the upgoing part was absorbed by the surface). Thus the flux is given by

$$\Phi = 2P_s / \sqrt{2\pi m k T} \tag{3.12}$$

We assume that the plasma is dominated by protons, so $m = m_p$, and we use a characteristic energy of particles in the cusp of 1.05 ± 0.95 keV from FIPS data (since FIPS observed particles with energies between 0.1 to 2 keV in the cusp) [22]. These assumptions yield an average surface flux of $\Phi_s = (2.1 \pm 1.0) \times 10^{12}$ particles m⁻²s⁻¹. This value agrees with the average flux over perihelion and aphelion conditions of $\sim 3 \times 10^{12}$ m⁻²s⁻¹ predicted by Sarantos et al. [110].

From here, we estimate the total number of particles that hit the surface in the cusp region by projecting the area of the cusp at the altitude of the descending tracks down to the surface. We calculate an upper and a lower limit for the area. Our lower limit is estimated from the area A of a trapezoid that encompasses the minimum region over which the cusp is observed at the altitude of the descending tracks in Figure 3.3, projected down to the surface, using $A \propto 1/B$. Our upper limit is estimated by taking a circular area of radius equal to half the maximum latitudinal extent of the cusp at the lower altitudes, and again projecting this down to the surface. The mean cusp area is then calculated to be $(5.2 \pm 1.6) \times 10^{11}$ m², centered at 74.7° MSO latitude on the surface, and we find that $(1.1 \pm 0.6) \times 10^{24}$ particles bombard the northern cusp region every second. The uncertainty in our total flux is primarily from the uncertainty in the proton temperature and the area of the cusp.

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As a check on the validity of Liouville's theorem we compared the pressures estimated for the ascending and descending tracks. For complete surface absorption, the ratio of the perpendicular pressures at the two altitudes should be $P_{\perp 1}/P_{\perp 0} = 0.93$, so we expect the pressure ratio to be between 0.93 and 1. For IMF $B_x < 0$, under which solar wind P_{Ram} values are comparable between ascending and descending tracks (Table 3.1), the pressures are the same to within the uncertainties.

Because of the northward offset of Mercury's dipole, and the resulting weaker surface magnetic field at high southern than northern latitudes, we expect the flux of precipitating particles to occur over a larger area in the southern cusp region than in the north. In the absence of observations of the southern cusp, we use the offset of the dipole magnetic field to estimate the total number of particles reaching the surface in the south. We calculate the central MSO latitude of the southern cusp to be about 64° S. The magnetospheric model [17] predicts a surface field strength at this latitude of 158 nT. From the ratio of the model surface field strength in the south to that in the north we estimate that the cusp area in the south is 2×10^{12} m², and the number of particles reaching the surface in the southern cusp region is correspondingly higher, 4×10^{24} particles s⁻¹. Over a Mercury solar day, the planetary surface rotates under the cusp, so the cusp precipitation reaches all planetary longitudes in a band extending $\sim 1600 \text{ km} (\sim 38^{\circ})$ in latitude. The IMF B_x effect we observe here, corresponding to 40% higher pressures in the northern cusp for negative than for positive IMF B_x , implies that the flux to the southern cusp should dominate regardless of the IMF direction. These hemispheric flux differences would lead to a persistently greater exospheric source from the south, as sputtering is most likely a contributing factor in populating the exosphere of Mercury [110]. In addition, if solar wind ion sputtering is a dominant source of space weathering at Mercury, this signature may be observed in surface reflectance spectra. It is, however, possible that Mercury's current magnetic field configuration has not been in place sufficiently long compared with space weathering timescales for this hemispheric asymmetry to be evident in surface reflectance and color differences. Alternatively, the surface may already have reached saturation, and

hemispheric differences will be muted.

3.5 Conclusions

From six months of MESSENGER MAG observations we have characterized Mercury's northern cusp region and found that it is persistently present but variable in extent and in the depth of its magnetic field depression. We focused on the role of the IMF $B_{\rm x}$ direction and the solar wind ram pressure in modulating the average plasma pressure in the cusp because of possible observable consequences for exospheric processes and space weathering. The northern cusp is clearly evident even during sunward IMF conditions but exhibits 40% higher plasma pressures on average during anti-sunward conditions, indicating that the effect of IMF $B_{\rm x}$ direction is present. Rapid variability in cusp pressures and orbit-to-orbit variations in the latitudinal extent of the cusp may be related to magnetospheric dynamics associated with southward IMF conditions. We estimate that on average $(1.1 \pm 0.6) \times 10^{24}$ particles per second reach Mercury's surface in the northern hemisphere cusp region, thus (via sputtering) contributing a source for the exosphere. Because of the northward offset of the planetary dipole, the flux of particles bombarding the southern cusp should be a factor of 4 higher, yielding a greater exospheric source in the south. Similarly, space weathering in the south due to cusp precipitation should occur over an area 4 times larger than in the north (or equivalently, over a latitudinal extent that is a factor of two larger). The implications of the north-south magnetic asymmetry for exospheric dynamics are therefore substantial and warrant efforts to confirm the estimated difference in surface magnetic field intensities. Whether a north-south asymmetry is evident in surface reflectance differences depends on the length of time that the present north-south asymmetry in the magnetic field has been maintained.
Chapter 4

Mercury's surface magnetic field determined from proton-reflection magnetometry¹

Solar wind protons observed by MESSENGER exhibit signatures of precipitation loss to Mercury's surface. In this chapter, we apply proton reflection magnetometry to sense Mercury's surface magnetic field intensity in the planet's northern and southern hemispheres [117]. The results are consistent with a dipole field offset to the north and show that the technique may be used to resolve regional-scale fields at the surface. The proton loss cones indicate persistent ion precipitation to the surface in the northern magnetospheric cusp region and to the southern hemisphere at low nightside latitudes. The latter observation implies that most of the surface in Mercury's southern hemisphere is continuously bombarded by plasma, in contrast with the premise that the global magnetic field largely protects the planetary surface from the solar wind.

4.1 Introduction

A remarkable feature of Mercury's weak, internal, magnetic field, indicated by orbital observations, is a ~ 480 km northward offset of the magnetic equator from the planetary equator [17–19]. The low magnetic field

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4.1. Introduction

strength and the northward offset provide constraints on Mercury's enigmatic dynamo mechanism [118] and lead to direct interactions between the solar wind and the planet's surface. The weak magnetic field allows precipitation of solar wind plasma to the surface in the northern magnetospheric cusp region. Since the surface magnetic field strength at any southern latitude is predicted to be weaker than at the corresponding northern latitude, enhanced ion sputtering and space weathering in the southern hemisphere is possible. In this study we apply proton-reflection magnetometry, adapted from electron reflectometry [119–121], to determine Mercury's surface magnetic field strength in both hemispheres and measure particle precipitation to the surface.

Electron reflectometry (ER) has been used extensively at the Moon [122], Mars [121], and Ganymede [123] to sense remotely the magnetic field strength at the surface. ER depends on the magnetic mirroring effect, that is, the reflection of electrons by convergent magnetic fields. Electrons that would mirror below the surface are lost and the flux of reflected electrons exhibits a sharp flux drop at the pitch angle (the angle between the particle velocity and the local magnetic field direction) corresponding to mirroring at the surface. The in-situ magnetic field together with the pitch angle of the last reflected electrons, the cut-off pitch angle, indicates the surface magnetic field strength. This technique has not yet been applied using protons.

At Mercury, protons with energies of 0.3 - 10 keV are regularly detected inside the magnetosphere [22] by the Fast Imaging Plasma Spectrometer (FIPS) [50] on the MErcury Surface, Space ENvironment, GEochemistry, and Ranging (MESSENGER) spacecraft. The Energetic Particle Spectrometer [50] on MESSENGER detects electrons with energies only above 35 keV and observes significant fluxes too infrequently to use ER. We therefore use FIPS observations of protons within Mercury's magnetosphere [53], together with magnetic field observations from MESSENGER's Magnetometer (MAG) [49], to estimate the magnetic field strength at Mercury's surface.

4.2 The technique of proton reflection magnetometry

In this section, we apply the electron reflectometry technique to protons. Similarly to electron reflectometry, proton-reflection magnetometry relies on magnetic mirroring, whereby protons travelling in a helical path along magnetic field lines reflect back along these field lines when the parallel component of their velocity vector (v_{\parallel}) becomes zero and all the velocity is in the perpendicular component (v_{\perp}) . Such a configuration occurs in convergent magnetic fields. The proton mass is much greater than the electron mass, and so proton-reflection magnetometry can only be conducted in regions where the ambient magnetic field is high enough for the protons to be "magnetized", i.e. to be directed by the magnetic field. It can be seen from equation (1.7) that the proton gyroradius is significantly higher than the electron gyroradius in the same ambient magnetic field due the higher mass, signifying that in weak magnetic fields, at the Moon for example, the dominant force affecting the motion of the protons will not be the Lorentz force. This is not the case at Mercury however, where the planetary magnetic field is high enough to "magnetize" both electrons and protons.

Protons reflect at varying altitudes depending on their initial pitch angles, and those that are initially aligned very closely to the ambient magnetic field do not reflect before they reach the surface and are absorbed by the surface. Thus a spacecraft flying above convergent magnetic fields will detect particles travelling in the direction of the magnetic field toward the planet's surface, as well as particles reflected back along field lines after having undergone mirroring. The spacecraft will also detect a loss cone in the reflected particle distribution that corresponds to particles absorbed by the surface. In what follows, we will show how it is possible to use the pitch angle information of the protons absorbed by the surface to infer the strength of the surface magnetic field.

In general, when conducting proton-reflection magnetometry at each integration period of a plasma spectrometer it is important to ensure that the same downgoing and upgoing plasma populations are sampled. This in turn requires that the distance travelled by the spacecraft during the proton's round-trip time has to be less than or equal to the gyroradius of the proton. This requirement however is not used in our analyses, because the limited FOV of the FIPS sensor means that we typically see only the incident or the reflected population at any one integration time, and so we average the results from many integration times to obtain the full incident and reflected populations (see Section 4.3).

4.2.1 How to obtain surface field strengths

In this section we describe how to use the magnetic mirroring formalism outlined in Section 1.2.4 to determine the magnetic field magnitude at the surface of the planet. The key factor in determining the surface magnetic field strength relies on measuring the pitch angle distribution of the protons (i.e., the flux of protons as a function of pitch angle). If there was no loss mechanism in the system (i.e. the surface), all particles incident along the magnetic field direction would reflect back along the field lines. Thus if the initial population entering from the magnetosheath along open field lines was originally isotropic (i.e. equal fluxes at all pitch angles), the observed particle pitch angles would be in the range of 0° to 90° for the incident population (in Mercury's northern hemisphere), and if these all reflected, particles with pitch angles in the range of 90° to 180° would be detected returning along the field lines. However, particles that are initially closely aligned with the magnetic field do not reflect before reaching the surface and are absorbed. There will thus be protons with a range of initial pitch angles which will be missing from the returning population. This missing part of the reflected pitch angle distribution, the loss cone (Figure 4.1), is determined by the ratio between the magnetic field magnitude at the spacecraft and that at the surface.

The cut-off pitch angle, α_c , that is the pitch angle in the reflected part of the distribution beyond which no more protons are observed, is the pitch angle that corresponds to protons mirroring exactly at the surface. These are the "last" protons which still make it back to the spacecraft, as all protons



Figure 4.1: Schematic pitch angle distribution showing the cut-off in reflected charged particle flux at angle α_c due to absorption by the surface.

that would have a reflected pitch angle larger than the cut-off pitch angle are absorbed. Thus protons with a cut-off pitch angle are remote-sensing the surface magnetic field because they mirror due to the magnetic field strength at the surface. We can therefore replace α_0 with α_c and B_m with $B_{\rm S}$, where $B_{\rm S}$ is the surface field strength in equation (1.16) to get

$$B_{\rm S} = \frac{B_0}{\sin^2 \alpha_c}.\tag{4.1}$$

Thus if we can measure the magnetic field strength at the spacecraft altitude, B_0 , and obtain α_c , it is possible to use equation (4.1) to obtain the surface magnetic field strength. Therefore the key to estimating B_S is determining the cut-off pitch angle in the pitch angle distribution.

4.3 Application of the technique to MESSENGER observations

For our analyses, we used one Earth-year of observations from 7 June 2011 to 7 June 2012. Combining observations from MAG and FIPS al-

lowed the calculation of proton pitch angle and the derivation of pitch-angle distributions (PADs) within the magnetosphere.

FIPS measures energy per charge (E/q), time of flight, and arrival incidence angle for ions, and completes one scan over the full range of E/qvalues every 8 s [124]. It has a conical field of view (FOV) of approximately 1.4π sr, with two symmetric cutouts of 15° near the instrument's symmetry axis and also near the plane perpendicular to the symmetry axis. Due to the limited field of view of FIPS, the full proton PAD (from 0° to 180°) is not visible at any one integration time. To build PADs spanning $\alpha = 0^{\circ}$ to 180° at times when the spacecraft was in a given region, we combined pitch angle distributions from instrument integration times when the incident proton population was observed with PADs when the reflected proton population was visible. This approach required averaging PADs over many time intervals and normalizing each contributing PAD to account for obstructions in the FIPS FOV. Here we focus on regions where the highest proton counts were detected by FIPS. The analysis thus determines Mercury's average long-wavelength (e.g., dipole) field but does not resolve shorter-wavelength structure.

4.3.1 Deriving individual pitch angle distributions

FIPS Pulse Height Analysis (PHA) data are used to create the proton pitch angle distributions. These distributions cannot be constructed simply from the raw data files because the following instrument limitations have to be taken into account: the size of the solid angle of the pixels on the detector, the efficiency of the micro-channel plates (MCP) detector, and the FIPS field of view (FOV) obstructions at each integration time. Because only angular structure is of interest for these accumulations, proton events from all E/q steps were added together to improve the signal-to-noise ratio.

To account for the solid angle size variation with zenith angle, θ , of the pixels on the detector, we divide each proton count obtained from the PHA data (Figure 4.2A) by the solid angle size given by

$$\sin\theta d\theta d\phi = 4.248 \times 10^{-5} r^2 - 1.114 \times 10^{-3} r + 1.15 \times 10^{-2} \tag{4.2}$$

where r is given by

 $r = 3.60348 - 0.045364 \ \theta + 0.00648116 \ \theta^2 - 0.0000451434 \ \theta^3. \tag{4.3}$

This normalization does not have a large effect on the proton counts because the variation in solid angle size with zenith angle is small. Next, because the FIPS MCP detector efficiency varies with location on the detector [125], our counts are also divided by an MCP correction factor to ensure that counts that fell on the lower efficiency part of the detector carry a higher weight. After rotating the vector magnetic field data measured by MAG into the FIPS reference frame, the pitch angles for all proton counts are computed. The solid-angle-weighted and MCP-normalized proton counts that fall within 10° pitch angle contours are then summed into 10° pitch angle bins, centered on pitch angles of 5°, $15^{\circ},...,165^{\circ}, 175^{\circ}$.

The summed proton counts are then also weighted by a factor that takes into account the FIPS visibility of each magnetic pitch angle bin [125]. This normalization is a function of the FIPS orientation with respect to the magnetic field direction and FOV obstructions at each integration time (Figure 4.2B). The counts are divided by this factor, ensuring that pitch angle bin centers which are not fully in the FOV are weighted higher to account for unobserved proton counts at those pitch angles. We thus finally arrive at a pitch angle distribution at every integration time of the instrument, where the proton counts have been properly weighted in order to take into account all the instrument limitations (Figure 4.2C). Errors assigned to each pitch angle bin center incorporate counting statistics and the fraction of proton gyrophase angles that was visible in the FOV for each pitch-angle bin.

4.3.2 Averaging pitch angle distributions

Because of the limited FOV of FIPS, there are no integration periods when both incident and reflected particle distributions are fully observed. In order to increase the signal-to-noise ratio in our observations, as well as to represent the entire PAD, we averaged PADs from different integration periods when the spacecraft was over the same MSO location. To connect



Figure 4.2: Example integration period of FIPS. (A) Zenith angle versus azimuth angle in the FIPS reference frame. The green dots show the proton count locations in angle-angle space, and pitch angle contours are for that integration period (the magenta curve is a pitch angle of 90° , and the green curve is 150°). The spacecraft altitude, latitude, and longitude are also given. (B) Weighting factor as a function of pitch angle. The limited FOV of FIPS as well as the viewing geometry given the local magnetic field direction are taken into account in the weighting. (C) Derived pitch angle distribution at the given integration period. Error bars reflect counting statistics as well as the fraction of the proton gyrophase angle that was visible in the FOV for each pitch angle bin.

incident and reflected particle populations from different integration periods, we normalized each individual PAD by the weighted proton count observed at a pitch angle of either 85° or 95° , depending on which of the two pitch angles was visible in the distribution at that time. If both were observed, the proton count at 85° was used. This methodology was applied to electron reflectometry measurements at the Moon [126], where the average proton counts in the incident population were used to normalize the PAD. Since FIPS does not observe the full incident population at the same time as the reflected population, we used the proton count at $\alpha \sim 90^{\circ}$ to approximate the average incident population. We then computed weighted averages of the individual PADs, where the error at each pitch-angle bin center on the PADs was assigned as the weight. The error assigned to the final averaged distribution is one standard error in all the proton counts that were averaged at each pitch angle bin center. Averaged pitch angle distributions were derived in the cusp region in the northern hemisphere, as well as in regions of high proton flux on the nightside at low latitudes in the southern hemisphere.

4.3.3 Surface field estimates

Averaged pitch angle distributions were derived in the cusp region in the northern hemisphere (geographic latitude $\lambda > 60^{\circ}$ N), as well as in regions of high proton flux on the nightside at low latitudes ($0^{\circ} < \lambda < 30^{\circ}$ S) in the southern hemisphere. Although high proton fluxes to the dayside southern hemisphere, especially in the southern cusp region, are expected, MESSEN-GER is unable to measure these because its eccentric orbit allows it to be inside the magnetosphere only over northern latitudes on the dayside. The averaged PAD for the northern cusp region, which includes 485 reflected and 185 incident population scans, is shown in Figure 4.3A. The incident population is approximately independent of pitch angle (i.e., isotropic), consistent with protons entering along open field lines from the magnetosheath. However, there is a void, or loss cone, in the fluxes of reflected protons, from $\alpha \sim 120^{\circ}$ to 180° , evidence for incident protons having been 'lost' to Mercury's surface. Relative to the sharp cut-off pitch angle observed in ER

at the Moon [127], the edge of the proton loss cone is smoothed over ~ 30° in α . We attribute this to the combined effects of the FIPS angular resolution of ~ 15° and pitch angle diffusion from wave-particle scattering in the cusp. Broadband magnetic field fluctuations between 0.001 and 10 Hz, are consistently observed in this region [128], and would scatter protons in α .

We estimated the loss cone angle consistent with diffusive scattering by fitting solutions to the diffusion equation to the loss cones of our pitch angle distributions. We solved the one-dimensional diffusion equation

$$\frac{\partial u(\alpha, t)}{\partial t} = D_{\alpha} \frac{\partial^2 u(\alpha, t)}{\partial a^2}$$
(4.4)

where u is the proton count and D_{α} is the diffusion coefficient, with a step-function initial condition:

$$u(\alpha, 0) = \begin{cases} c_1 & \text{for } \alpha \le \alpha_c, \\ c_2, & \text{for } \alpha > \alpha_c. \end{cases}$$
(4.5)

where c_1 and c_2 are constants set by the average maximum and minimum weighted proton counts in the PAD. The boundary conditions were given by:

$$\frac{\partial u(0,t)}{\partial \alpha} = \frac{\partial u(\pi,t)}{\partial \alpha} = 0.$$
(4.6)

The use of reflective boundary conditions (i.e. zero gradient in the flux at the boundaries), as opposed to one with a negative gradient, is justified for this diffusion process as there is no loss of particles in the field aligned (or anti-aligned) direction from pitch angle diffusion. This is due to the fact that the diffusion process occurs after particles have mirrored and are travelling up along magnetic field lines away from the surface. As such, particles that are scattered into the loss cone will be scattered back and forth across $\alpha = 180^{\circ}$, but will not be lost to the surface. There should thus be no loss of protons at the boundaries.



Figure 4.3: Pitch angle distributions in the northern and southern hemispheres. (A) Results for Mercury's northern cusp. (B) Results for the lowlatitude southern hemisphere nightside. Average proton counts are in red with standard errors. The black curve shows the diffusion model fit to the reflected portion of the distribution; the fit uncertainty is in gray. Yellow shading indicates the loss cone; the black error bar shows the uncertainty in α_c . (C) Comparison of the southern hemisphere PAD (red curve) from panel (B) to a model single-sided loss cone distribution (black curve) and an observed double-sided loss cone distribution (blue curve). The character of the southern hemisphere PAD is in-between that of a single-sided and a fully formed double-sided loss cone (see text below and in Chapter 5).

The solution to equation (4.4) is

$$u(\alpha, t) = \sum_{n=1}^{\infty} B_n \cos(n\alpha) \, \mathrm{e}^{-n^2 D_\alpha t}, \qquad (4.7)$$

where

$$B_n = \frac{2}{\pi} \left[\int_0^{\alpha_c} c_1 \cos(n\alpha) d\alpha + \int_{\alpha_c}^{\pi} c_2 \cos(n\alpha) d\alpha \right].$$
(4.8)

We fit equations equations (4.7) and (4.8) to our loss cones and allowed the cut-off pitch angle, α_c , and $D_{\alpha}t$ to vary freely. We used a grid search method that minimized the median absolute deviation (MAD) between the model and the observations. Figure 4.4 shows the contour plots of the absolute value of the residuals for the models best fit to the averaged PAD for the northern cusp region as well as for the southern hemisphere. We established upper and lower bounds on the parameters α_c and $D_{\alpha}t$ by identifying an allowable upper bound on the misfit, corresponding to a 95% confidence limit (bold contours in Figure 4.4). The upper and lower bounds on α_c were identified as the locations of the intersection of a horizontal cut (passing through the minimum misfit) with the bold contour. A corresponding vertical cut yielded the limits on the $D_{\alpha}t$ parameter. The bounds on the best-fit model, shown by the grey shaded regions of Figure 4.3, were determined from the diffusion curves corresponding to the upper and lower limits for α_c and $D_{\alpha}t$.

From the fit to the northern cusp PAD, we obtain a cut-off pitch angle, α_c , of $121^\circ \pm 3^\circ$, which together with the measured average magnetic field strength at the spacecraft altitudes (< 550 km), $B_0 = 302.4 \pm 53.0$ nT, implies a surface field strength of $B_{\rm S} = 412 \pm 98$ nT where we have used $B_{\rm S} = B_0/\sin^2(\alpha_c)$. The uncertainty accounts for the standard error in the fit value of α_c and also the standard deviation of B_0 , computed from all intervals in the average PAD.

We mapped the average observation location in the cusp down to the surface by tracing the magnetic field lines to the surface using Mercury's offset dipole magnetic field. The observation altitudes ranged from 282 km



Figure 4.4: (A) Residuals of the diffusion equation curve best fit to the loss cone of the averaged PAD in the northern cusp region, as a function of the cut-off pitch angle, α_c , and the product of the diffusion coefficient and time, $D_{\alpha}t$. The bold contour marks the residual level from which the errors on the fit parameters were obtained. (B) Same as A but for the southern hemisphere averaged PAD. The ratio of the bold contour to the minimum misfit is ~ 7 in (A) and ~ 3 in (B).





Figure 4.5: (A) Stereographic projection plot (looking down from above the north pole) showing the surface magnetic foot-point locations of the reflected proton observations in the northern cusp versus local time and latitude. The Sun is to the right. Latitudes north of $45^{\circ}N$ are shown. (B) Corresponding plot for the southern hemisphere (looking through the planet from above the north pole). Latitudes from $10^{\circ}N$ to $90^{\circ}S$ are shown.

to 549 km, with a mean value of 414 km. In latitude, the cusp observations at the surface extended 15.6° degrees in latitude and 7.5 h in local time and were centered on noon at 76.4°N latitude on the surface (Figure 4.5A).

We also find high proton fluxes in the latitudinal band $0^{\circ} < \lambda < 30^{\circ}$ S on the nightside, with a clear loss-cone signature in the derived PAD (Figure 4.3B), although with larger uncertainties than for the northern cusp region (Figure 4.3A). In the southern hemisphere, observations as far south as possible are desirable for observing the long-wavelength structure in the magnetic field. However, due to MESSENGER's eccentric orbit and high altitudes in the southern hemisphere, we are restricted to observations north of approximately 30°S latitude. In this averaged PAD, we included 128 scans in the reflected population and 315 scans in the incident population. The similar error bars on most of the incident and reflected population fluxes in Figure 4.3B, despite the higher number of observations being included in the incident side, is due to the significantly larger standard deviations in the fluxes of the incoming protons. The best-fit diffusion model to the loss cone gives $\alpha_c = 43^{\circ+7}_{-13}$, and this together with $B_0 = 52.5 \pm 14.8$ nT at the spacecraft altitude corresponds to a surface field strength of $B_{\rm S} = 113^{+87}_{-61}$ nT. The observation altitudes were higher than those in the northern cusp region, ranging between 1160 and 1980 km, with a mean of 1535 km. The mapped surface locations span $23^{\circ}{\rm S} < \lambda < 34^{\circ}{\rm S}$, with a mean of 27.8°S, and local times spanning the nightside from 16 h to 5.3 h, centered on 23.5 h (Figure 4.5B).

The apparent secondary loss cone in the incident population for this averaged PAD (Figure 4.3B) implies that these observations may correspond to closed field lines on the nightside. However, as Figure 4.3C shows, even though these observations are likely from a closed field line region, the shape of the southern hemisphere PAD is significantly different from that a fully formed double-sided loss cone found on closed field lines on the dayside (see Chapter 5), and is in between the character of an idealized single-sided and a fully double-sided loss cone distribution. This is owing to fresh proton populations drifting onto closed field lines on the nightside from the plasma sheet, thereby continuously replenishing particles in the loss cone. We can thus approximate this nightside southern hemisphere PAD as a single-sided distribution, which suggests that the inferred surface magnetic field strength for this region may be a lower limit.

4.4 Consistency checks

To ensure that the FIPS proton data and our averaging method are viable for conducting consistent proton reflection measurements we performed a number of consistency checks. We binned the observations in the northern cusp region (the region where we have the most observations) in altitude as well as in latitude and local time, and tested that the derived averaged PADs yielded expected behavior with respect to the size of the loss cones and the estimated magnetic field strengths.



Figure 4.6: Pitch angle distributions derived from observations binned in altitude: red denotes low-, green mid-, and blue high-altitude observations.

4.4.1 Altitude binning

It is expected that over regions of approximately constant magnetic field on the surface, if the altitude of the spacecraft observations increases (and thus the measured magnetic field strength decreases) then the cut-off pitch angle will increase (or equivalently, the loss cone size will decrease). We tested that this is the case in our observations in the northern cusp region. Three altitude averages of the cusp observations were obtained (Figure 4.6). The first was at low altitude with a mean of 325 km and included 155 scans in the reflected and 85 in the incident population. The second was at mid altitude with a mean of 418 km and included 98 scans in the reflected and 96 in the incident side. And the third was at high altitude with a mean of 510 km and consisted of 173 scans in the reflected and 4 in the incident population.

Despite having approximately equal number of observations in the inci-

4.4. Consistency checks

dent and reflected populations for the mid altitude bin, the errors on the incident population are much larger because of the large standard deviations in the measured proton fluxes at those pitch angles. We fit diffusion equation curves to these averaged PADs to derive a cut-off pitch angle for each. We find that the cut-off pitch angle increases with increasing altitude, as expected. The derived cut-off pitch angles were: $116^{\circ+6}_{-4}$, $123^{\circ+3}_{-4}$, $126^{\circ+2}_{-5}$ for the low-, mid-, and high-altitude bins, respectively. The smaller difference in loss cone cut-off angle between the mid- and high-altitude bins is attributed to the difference in the average latitude of observations between these two bins (the mean latitude at the spacecraft altitude of the mid altitude bin was 76.1°N, whereas it was 73.7°N for the high altitude bin). This latitude bin than for the high-altitude bin, offsetting the expected altitude-dependent difference in cut-off pitch angle.

4.4.2 Latitude binning

We also obtained averages of PADs in three latitude and two local-time bins in the cusp region to test whether we observe the latitudinal increase in the magnetic field strength expected for an intrinsic dipole field. The observations were binned not just in latitude but also in local time to minimize the range of altitudes included in each average. Phasing of MESSENGER orbit-correction maneuvers with local time resulted in dusk observations that were systematically taken at lower altitudes than the dawn observations. By splitting the data into dawn and dusk sections, and binning each section separately into three latitude bins (the latitude at the spacecraft altitude and the altitude of the bin centers is given in Figure 4.7), we obtained six PAD averages in the northern cusp region.

After fitting diffusion equation curves to these PAD averages and obtaining the cut-off pitch angle for each averaged PAD, we find that the magnetic field strength increases with increasing latitude, as expected. The mean measured magnetic field strengths at the spacecraft altitudes on the dawn side are 250.2 ± 33.8 nT, 276.2 ± 39.8 nT, and 311 ± 46 nT for low-,



Figure 4.7: Pitch angle distributions binned in latitude and local time in the northern cusp. (A) Averaged PADs for latitude bins on the dawn side of the northern cusp region. Colour coding is as follows: red is low-, blue mid-, and green high-latitude. The mean latitudes and altitudes of the observations in each bin are also given. (B) Same as A but for dusk-side observations.

mid-, and high-latitude data bins, and the corresponding surface magnetic field strengths are 334^{+65}_{-72} nT, 376 ± 70 nT, and 432 ± 92 nT, respectively. The loss cones are not as well defined on the dusk side, especially in the case of the low-latitude bin, which has very high errors reflecting the small number of observations that were included in the average. The existence of a loss

cone in the averaged PAD for the dusk-side low-latitude bin is questionable, and so we ignore this bin in our calculations. The mid- and high-latitude bins on the dusk side do exhibit loss cones, albeit more smeared out than those seen on the dawn side. We estimate that the magnetic field strength at the surface for these two bins of data (using $B_0 = 327.0 \pm 46.9$ nT and $B_0 = 343.0 \pm 37.1$ nT for the mid- and high-latitude bins) to be 412^{+140}_{-74} nT and 432 ± 62 nT for the mid- and high-latitude bins, respectively, consistent with an increase in field strength with increasing latitude. The estimates for the surface field agree within 10% between the dawn and dusk side for the mid- and high-latitude bins.

4.5 Discussion and conclusions

Estimates of the surface magnetic field strength in the northern cusp and low latitude southern hemisphere are compared with predictions from the best-fit time averaged magnetospheric model [19] in Figure 4.8 and Table 4.1. The results from proton-reflection magnetometry are significantly lower than the model magnetic field. Such a difference is expected however because the model is a vacuum magnetic field model, whereas our PADs demonstrate the presence of plasma extending to the surface of the planet. The plasma generates a diamagnetic field, which will reduce the surface field below the vacuum model prediction [89, 116, 129].

Fortunately, the proton data provide the information required to estimate the diamagnetic effect (see Appendix A for a detailed derivation of the diamagnetic field and flux to the surface). The flux of particles at the surface in the northern cusp region can be determined from the loss cone size and mean proton temperature, T_p , and density, n_p , in the cusp, yielding the proton pressure $P_p = n_p k T_p$, where k is Boltzmann's constant. The typical proton density ($n_p \approx 30 \text{ cm}^{-3}$) and temperature ($T_p \approx 12 \text{ MK}$) derived from FIPS observations in the cusp [125] are for an isotropic particle distribution; the anisotropy associated with the loss cone produces an underestimate of plasma density [130]. Taking this anisotropy into account yields a surface flux of 3.7×10^{12} particles m⁻²s⁻¹, which is approximately in agreement with

Latitude	Local time	$B_{\rm PR}$	B_{Plasma}	$B_{ m Plasma}$ - $B_{ m PR}$	$B_{\rm Model}$
(MSO	(MSO)	(nT)	(nT)	(nT)	(nT)
surface)					
76.4°N	12 h	412 ± 98	-86 ± 11	498 ± 99	647
$27.8^{\circ}\mathrm{S}$	23.5 h	113^{+87}_{-61}	-44 ± 6	157^{+87}_{-61}	224

4.5. Discussion and conclusions

Table 4.1: Surface magnetic field strength from proton-reflection magnetometry compared with magnetospheric model predictions. The center latitudes and local times of the northern cusp and southern hemisphere PAD averages are given, as well as our surface field estimates with and without the diamagnetic effect of the plasma near the surface. The vacuum magnetic field model predictions for comparison with the corrected proton-reflection magnetometry estimates are also listed.

the flux determined by Winslow et al. [89] for the cusp region. In the region of our southern hemisphere average, we have $n_p \approx 5 \text{ cm}^{-3}$ and $T_p \approx 20 \text{ MK}$, which give a flux of 4.4×10^{11} particles m⁻²s⁻¹. To take into account the fact that the loss cone in the incident population for this PAD is only $\sim 70\%$ filled in (i.e. the PAD is not fully single-sided), we multiply this flux value by 0.7 to obtain the flux to the surface in this region of 3.1×10^{11} particles $m^{-2}s^{-1}$. The diamagnetic effect reflects particle motions only in the direction perpendicular to the local magnetic field, so we used the perpendicular particle pressure to calculate the diamagnetic field at the surface. We find an average diamagnetic effect, ΔB_{Plasma} , in the northern cusp of 86 ± 11 nT, and $\Delta B_{\text{Plasma}} = 44 \pm 6 \text{ nT}$ in the southern hemisphere at low latitudes. The uncertainties in these values represent the limits on the diamagnetic field derived from the pressure at the spacecraft altitude and that at the surface. After accounting for this effect, our magnetic field estimates agree with the model predictions in the southern hemisphere and are within 8% (our upper bound) of the model prediction in the cusp (Table 4.1).

The validity of the offset dipole can be tested by estimating the ratio between the surface magnetic field strength in the northern cusp and that in the southern hemisphere. From the mean observation locations, we find that the ratio from proton reflection magnetometry, corrected for the diamagnetic effect, is $3.2^{+3.0}_{-1.6}$ in agreement with the result of 3.3 for an offset dipole field



Figure 4.8: Mollweide projection of surface magnetic field strength from the best-fit time-averaged magnetospheric model [19] in MSO coordinates. The equator offset and the magnetopause and magnetotail fields cause departures from the field of a dipole alone. Black contours mark the approximate outline of the regions sampled by proton-reflection magnetometry observations.

alone. The full magnetospheric model [19] yields a surface field ratio of 2.6 to 3.6 for the latitudinal and longitudinal extent of our observation locations, with a mean value of 2.9. A centered dipole field alone gives a ratio of 1.5 for the latitudes sampled and is not in agreement with our inferred surface field ratio. Even if we consider that the southern hemisphere surface magnetic field strength may be a lower limit, we can still conclude that the offset dipole is confirmed as long as $B_{\rm S}$ in the southern hemisphere is not more than a factor of ~ 2.5 greater than our derived value. Such a large value for the low latitude southern hemisphere field is incompatible with vector magnetic field measurements by MESSENGER. Lastly, it is important to note that the large uncertainty assigned to our southern hemisphere loss cone size is approximately equal to the difference in loss cone angle size between our PAD in the southern hemisphere and a fully formed doublesided PAD from the same region (see difference between red and blue curves for the reflected population of Figure 4.3C). Thus our large uncertainty on α_c for the southern hemisphere population (and the corresponding large uncertainty on $B_{\rm S}$) likely captures any possible pitch angle diffusion related

erosion of the loss cone size.

Proton-reflection magnetometry thus provides independent confirmation of the offset dipole nature of Mercury's magnetic field, although weaker terms of higher degree and order in a multipole expansion of the field are not ruled out by these results. Significantly, proton-reflection magnetometry demonstrates that persistent proton precipitation to the surface occurs on the nightside of Mercury's low-latitude southern hemisphere. This result implies that such precipitation may also be occurring at middle latitudes on the nightside, where closed field lines reach the surface, because the northward offset of the magnetic equator results in weak surface field strengths (not more than 25% above our measured value) everywhere in the southern hemisphere. Together with the proximity of the magnetopause to the surface and the large cusp region on the dayside southern hemisphere [17, 18], this suggests that most of Mercury's southern hemisphere surface may be continuously bombarded by plasma. Such continuous precipitation in the southern hemisphere implies that space weathering is not confined to the cusp region and may thus show limited latitudinal variation. Because of MESSENGER's eccentric orbit, proton reflection magnetometry at higher southern latitudes is not feasible, but observations from the ESA-JAXA BepiColombo mission [131], with a less eccentric orbit that will provide lowaltitude observations in both hemispheres, may provide an opportunity to probe the surface field more extensively and with greater spatial resolution.

Chapter 5

Regional-scale surface magnetic fields and proton fluxes to Mercury's surface from proton-reflection magnetometry

Following the successful implementation of the proton-reflection magnetometry technique to measure the surface magnetic field strength in two regions on Mercury's surface, we expand on this work [132]. In this section, we extend our study with 1.5 years of additional observations (including data from 2012 and 2013) by MAG and FIPS with the aim of resolving regional-scale structure in the surface magnetic field strength and deriving a plasma precipitation map of the surface.

Regional-scale structure in the magnetic field, if present, may be due to either core or crustal sources. If they are of core origin, surface field strength measurements could provide further constraints for dynamo models. Conversely, if the short-wavelength structure is due to crustal remanent magnetization, correlations between crustal fields and geologic features can help establish constraints on the temporal evolution of Mercury's dynamo field - an important constraint for geophysical evolution models.

Determining the average proton flux to various regions of Mercury's surface is important from a space weathering and exosphere production perspective because it identifies regions on the surface that are exposed to direct bombardment by the solar wind. The collision between energetic solar wind or magnetospheric ions with surface grains can release atoms via momentum transfer [133], a process known as sputtering. Sputtering can release any atomic species present in the near-surface mineralogy, so particles in the exosphere reflect the composition of the surface on an atomic level [134]. Such direct ion sputtering is only one of several mechanisms that space weathers planetary surfaces, and studies using MESSENGER exospheric measurements are currently under way to determine the role that sputtering plays in the generation of Mercury's exosphere. Thus it is highly desirable to distinguish regions on the surface that are exposed to relatively higher or lower fluxes of solar wind particles to inform exospheric studies.

Differential particle fluxes to various regions on the surface can also yield clues about magnetospheric processes. For example, particle fluxes to the surface outside of the cusp regions (where they are mostly expected), signify that drift and scattering processes are both occurring in the magnetosphere, allowing particles access to closed field line regions of the magnetosphere. In this work we improve on our previous estimates of proton precipitation to Mercury's surface.

5.1 Methods

In this section, we are interested in planetary processes, and conduct most of our analyses in a planetocentric Mercury body fixed (MBF) coordinate system. The MBF coordinate system governs the geometry of internally generated fields at Mercury. The MBF and MSO +z axes are nearly identical, and hence so are the MSO and MBF latitudes due to Mercury's small obliquity [4, 135].

In order to resolve regional-scale structure in either the surface magnetic field or in proton fluxes to the surface, we systematically derive averaged pitch angle distributions over all locations around the planet where high proton fluxes are observed. We grid the surface of the planet into MBF latitude and longitude bins, with 10° spacing in latitude and 20° spacing in longitude and record all individual pitch angle distributions that fall into

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each grid cell. This grid cell size is chosen to reveal variations in the surface field and fluxes on length scales as small as possible while still ensuring a high signal-to-noise ratio for the average PADs in each grid cell. Because our observations are taken at spacecraft altitudes, we map the observation locations along magnetic field lines to the planet's surface using the offset dipole magnetic field, and determine the surface grid cell corresponding to each FIPS integration period. We then evaluate an averaged pitch angle distribution from all the individual PADs in each grid cell, where the averaging is done as described in Section 4.3.2.

Observations were restricted to times when MESSENGER was inside the magnetosphere and spacecraft altitudes were below 3000 km. Bow shock and magnetopause boundary crossing times were not available past 15 December 2012, thus we used a magnetopause model shape with the best-fit time-averaged parameters defined in Chapter 2 to determine whether observations were within the magnetosphere. The altitude restriction also helps in confining observations to within the magnetosphere, except on the dayside where the magnetopause is not more than ~ 2500 km from the surface on average. We later examine more restrictive altitude bins of observations.

These analyses reveal loss cones in pitch angle distributions mapping to the surface in the northern hemisphere north of 40°N, as well as in the southern hemisphere between 20°S and 40°S. In the northern hemisphere, the observations between 40° and 60°N are not part of the cusp region and thus represent a population of high proton fluxes, that we were not able to observe previously with fewer observations. This region of high proton fluxes is consistent with higher plasma pressures in the same latitude range observed at spacecraft altitudes by Korth et al. [136]. Figures 5.1 and 5.2 show example observations of PADs from the northern and southern hemispheres, respectively. Figure 5.1 shows a clear transition in PAD characteristics with latitude: north of ~ 70°N latitude the PADs have single-sided loss cones, i.e. isotropic incident population and a clear loss cone in the reflected population, while south of ~ 60°N loss cones are observed in both the incident and reflected populations. These are termed double-sided loss cones. By comparing Figures 5.1 and 5.2, it can be seen that the character



Figure 5.1: Example pitch angle distributions (averaged over measurements taken at altitudes < 3000 km) in the northern hemisphere shown for 10° latitude bands. The latitude, longitude grid cell for each PAD is given in the top right corner of each panel. The change in character of the PADs can be seen with increasing latitude. Data is more sparse in the $80^{\circ} - 90^{\circ}N$ latitude band, which is signified by the large error bars on the normalized fluxes for the incident population.

of the PADs in the northern hemisphere between $40^{\circ} - 60^{\circ}$ N is very similar to those observed in the southern hemisphere. We attribute the doublesided loss cones observed in these regions to particles bouncing back and forth along closed magnetic field lines, with particles being lost to the surface on the side of the magnetic field line with the weaker magnetic field





Figure 5.2: Example PAD (averaged over measurements taken at altitudes < 3000 km) in the southern hemisphere. The latitude, longitude grid cell for the PAD is given in the top right corner of the figure. Most PADs observed in the southern hemisphere exhibit similar double-sided loss cones as observed at low latitudes in the northern hemisphere.

strength at the planetary surface. At Mercury, the large northward offset of the magnetic equator yields much weaker surface fields in the southern hemisphere than the northern hemisphere. Thus, particles on closed field lines will predominantly be lost to the southern hemisphere. This transition in PAD characteristics allows us to determine that on average, the boundary between open and closed field lines in the northern hemisphere maps to latitudes of $60^{\circ} - 70^{\circ}$ N on the surface. It is also important to note that the average loss cone size appears to be $10^{\circ} - 20^{\circ}$ larger for the double-sided loss cones than for the single-sided loss cones, a separate indication that particles are being lost to the southern hemisphere where the surface field is weaker.

To obtain surface magnetic field strengths in the latitude-longitude grid cells where a loss cone is observed, the cut-off pitch angle (or equivalently, the loss cone size) has to be established. In each grid cell we can also calculate the proton flux to the surface using the loss cone angle, the proton number

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density and temperature. Calculating the cut-off pitch angle by fitting the solution to the diffusion equation to the loss cone as described in Section 4.3.3 is straightforward in the region of single-sided loss cones. Here, near-isotropic proton populations enter the cusp region along open magnetic field lines; some are lost to the surface, while the returning population undergoes pitch-angle scattering from the plasma waves present in the region. In this process the PAD diffuses from a step-function-like drop-off in the fluxes, to the observed smooth loss cone. We thus use the solution to the diffusion equation with a step-function initial condition to fit to the single-sided PADs to obtain the cut-off pitch angle.

In the case of double-sided loss cones the use of the diffusion equation is no longer physically well motivated because the PADs are in an equilibrium state. After just one bounce period the PADs will have loss cones on both sides corresponding to the southern hemisphere, and will thus not have an initially step-function-like PAD that undergoes diffusion. In this regime, due to persistent pitch angle scattering as the particles bounce back and forth along the close field line, the loss cone size (which corresponds to the southern hemisphere) will slowly get larger as particles from the non-losscone part of the distribution get scattered into the loss cone. The particles scattered into the loss cone over time will make it down to the surface, so this larger loss cone angle of the double-sided distributions (which we term "large void" so as not to confuse terminology with the original loss cone) does still reflect particles that made it to the surface, albeit largely due to scattering. Thus we can still obtain the flux of particles to the surface from this large void, however the flux does not correspond to a per second flux to the surface, but to a flux over the length of time the particles have spent bouncing back and forth along the field line. This time cannot be easily determined as we have no information as to when the particles entered onto the closed field lines. However, by estimating how large the void is in the PADs we can obtain a flux to the surface over some average time that particles remain trapped on closed field lines at Mercury prior to being lost to the surface, the magnetopause, or down the magnetotail through a combination of drift motions and pitch angle diffusion.

5.1. Methods

The void size is obtained by fitting a curve that matches the shape of these PADs, and since the diffusion curve is still a fairly good approximation of the shape of the void, we implement it here again knowing that the parameter $D_{\alpha}t$ derived from these fits is no longer physically meaningful. Deriving a physically motivated shape model for the double-sided loss cones would involve properly accounting for a number of different factors, including: pitch-angle diffusion along the particle's bounce motion, various particle drift processes, and loss to the surface. It would also require knowledge of the length of time (number of bounce periods) taken to reach the observed PAD shape. Thus modeling of the voids is not pursued here.

Calculating the surface field strength from the large voids of the doublesided PADs is not warranted because the size of the large voids does not correspond to the actual loss cone size that would be there in the absence of pitch angle diffusion and at present it is not known how much they have been altered. It is important to note that the estimated surface field strength that was obtained in Chapter 4 for the southern hemisphere is still justified, because the PAD in Figure 4.3 exhibits a much more isotropic incident population (and can thus be approximated as a single-sided loss cone) than the very clearly anisotropic incoming populations in the newly derived PADs in the southern hemisphere latitude/longitude grid cells. The discrepancy between the previously derived PAD and the new PADs in the same latitude range is due to the binning in body-fixed longitude, in which we average pitch angle distributions from all local times. Previously, our observations in the southern hemisphere were all confined to the nightside, where fresh proton populations, which are mostly isotropic, drift in from the plasma sheet onto closed field lines. There are no sources of fresh proton populations on the davside at low-to-mid latitudes, however, and because there are a factor of ~ 2 more dayside observations than nightside observations, the dayside observations dominate the signal in our southern hemisphere bins here. We have verified that if we confine our new observations to the nightside, we still detect nearly isotropic incident populations; however, the reduced number of integration periods means that binning in both local time and body fixed longitude is not vet possible. Thus here we retain the PADs binned in

body-fixed longitude but averaged over all local times.

The two different regimes of the PADs (single-sided and double-sided) allow us to calculate magnetic field strengths and per-second fluxes from single-sided loss cones at high latitudes in the northern hemisphere and to obtain bounce-averaged fluxes from double-sided loss cones at lower northern hemisphere latitudes and in the southern hemisphere. However, as the voids in the double-sided PADs observed in the north still correspond to the southern hemisphere, we do not gain more information from these about the fluxes to the north, only to the south. In the following two sections we describe our results for the surface field strengths and fluxes separately.

5.2 Resolving regional scale surface magnetic field strengths

Due to the constraints mentioned in the previous Section, deriving surface magnetic field strengths from the latitude/longitude binned observations is only feasible in the northernmost latitude range, north of the boundary between open and closed field lines. Not all latitude-longitude grid cells exhibit loss cones in the PADs, we thus visually inspect each averaged PAD and decide whether there is a clear loss cone to which we can fit a diffusion curve. Although high proton fluxes from the cusp are expected in the latitude band between 60°N and 70°N at the surface, we do not observe any well defined loss cones in this region. This may be because the transition between open and closed field lines occurs in this area.

High proton fluxes and clear loss cones are evident in the latitude band between 70° and 80° N. Some PADs with loss cones are also observed up to 90°N, although there are fewer observations in this region and thus the error bars are on average higher on the normalized PAD proton counts. To derive surface magnetic field strengths in these regions we binned the observations in different altitude ranges to reduce the error on $B_{\rm S}$. However, even for a 10 km altitude range the standard deviation in the measured magnetic field strengths at spacecraft altitudes was greater than 50 nT in some cases.



Figure 5.3: Example PAD corresponding to the $70^{\circ} - 80^{\circ}N$ latitude and $60^{\circ} - 80^{\circ}E$ longitude grid cell with the solution to the diffusion equation fit to the loss cone.

This is likely due to the large variation in plasma pressures inside the cusp at the integration times averaged in the PAD. We are not currently able to properly account for the variations in the diamagnetic field with time in different regions. Thus variations in B_0 translate directly into uncertainties in the estimated magnetic field strengths at the surface. To minimize this problem, instead of binning the observations in altitude, we bin them by the measured magnetic field strength at the spacecraft altitude, ensuring that the diamagnetic field values are more similar at the selected times.

Clear loss cones were obtained for nearly all longitude bins between 70° and 80°N and one longitude bin at 90°N when data with $B_0 > 300$ nT were binned in the different regions on the surface. The average standard deviation in B_0 in these bins was ~ 25 nT, about a factor of two smaller than if the observations had been binned in altitude. In all the surface grid cells where PADs with clear loss cones are detected, we fit diffusion curves to the reflected population to estimate the cut-off pitch angle. Figure 5.3 shows an example averaged PAD with a diffusion curve fit to the loss cone. This PAD





Figure 5.4: Regional-scale surface magnetic field strengths estimated by proton-reflection magnetometry in the northern hemisphere, a) uncorrected for the diamagnetic field and b) corrected for the diamagnetic field. The maps are in a stereographic projection, looking down from above the north pole.

maps to between $70^{\circ} - 80^{\circ}$ N latitude and between $60^{\circ} - 80^{\circ}$ E in longitude. The magnetic field strength at the spacecraft for this particular example was $B_0 = 327 \pm 22$ nT and the derived α_c was $120^{\circ+6}_{-4}$, yielding a surface magnetic field strength of 436^{+82}_{-64} nT, not including the diamagnetic field correction. The uncertainties were derived using the approach described in Chapter 4. From similar fits to all suitable PADs, we establish a map of the regional-scale magnetic field strength in the northernmost hemisphere, shown in Figure 5.4.

The method of estimating the diamagnetic field described in Section 4.5 has been improved upon from that described in Section 4.5. Instead of taking an average proton number density, temperature, and thermal pressure over the entire cusp region, we take derived values of n_p , T_p , and P_p (the proton thermal pressure) from proton observations at specific integration times from 2011 that match the times averaged in our PADs. These derived products were made available by the FIPS instrumentation team on MESSENGER, and are described in [52]. Unfortunately the times that these products are derived do not extend into 2012 or 2013, but we can still use them to establish





Figure 5.5: a) Diamagnetic field corrected surface field strengths with error bars as a function of longitude in the northern cusp. The number of individual PADs averaged in each grid cell was between 7 and 32. b) Upper limits on the surface magnetic fields strengths (from the addition of the upper error bars in a) to the corrected surface field values) in the cusp.

average proton conditions in the latitude bins used here, although not in the longitude bins due to the limited number of observations. We set the errors on the derived FIPS products to be 1 standard error in the estimated n_p , T_p , and P_p values within each latitude bin.

The FIPS estimated average P_p at spacecraft altitude is 1.62 ± 0.17 nPa at 70° – 80° N latitude. Using this total proton pressure at the spacecraft altitude, we calculate the perpendicular pressure at the spacecraft using equation (A.7), correcting for the fact that P_p was derived assuming an isotropic distribution function. We also calculate the perpendicular pressure at the surface using equation (A.9). From equation (A.1), we obtain the corresponding diamagnetic field for these perpendicular pressures, and take the average of the two to get an approximate range for the diamagnetic field at the surface. The errors on the corrected surface magnetic field strengths incorporate the average error in the diamagnetic field, and are shown in Figure 5.5a, while the upper limits on our proton-reflection magnetometry derived surface magnetic field strengths are shown in Figure 5.5b.

On average the errors on $B_{\rm S}$ are higher than the estimated variation in $B_{\rm S}$

from one longitudinal bin to the next. However, we can draw two important conclusions from these surface field estimates. One is that we do observe an increased magnetic field strength with increasing latitude, $B_{\rm S}$ for the $80^{\circ} - 90^{\circ}$ N bin is significantly higher than the $B_{\rm S}$ for the latitude bin below it. More importantly, the observed fields at high latitudes combined with our southern hemisphere surface field result in Chapter 4 are somewhat weaker but consistent within the uncertainty with an offset dipole (see Figure 4.8). This result establishes that higher degree and order core or crustal fields on Mercury must either be very weak, or must be on length scales much smaller than our longitudinal bin of ~ 300 km.

5.3 Particle fluxes to Mercury's surface

We expand our calculation of the proton flux to the surface in the northern cusp and southern hemisphere low latitude region described in Section 4.5 to incorporate the regional scale proton loss cones observed in Section 5.2 above, and build a proton precipitation map of Mercury's surface. We can obtain fluxes to the surface from both single-sided and double-sided PADs as described in Section 5.1, with the caveat that the fluxes determined from double-sided loss cones are over an average bounce life-time of protons in Mercury's magnetosphere. It is important to note that we also assume for these double-sided loss cones that the original incident population was isotropic, prior to becoming trapped on closed field lines, i.e. these flux estimates are upper limits. The double-sided loss cones measured in the northern hemisphere correspond to particles being lost to the southern hemisphere surface. Therefore, our measurements allow for the resolution of regional-scale fluxes to the surface in the northernmost region of the northern hemisphere, where the single-sided loss cones are observed, and to regions of the southern hemisphere mapped by the double-sided loss cones observed in the north and the PADs observed in the south.

To determine the proton fluxes to the surface, we average individual PADs from all spacecraft altitudes (up to 3000 km) for the southern hemisphere measurements but only average over observations at < 550 km al-



Figure 5.6: Example pitch angle distribution exhibiting symmetry in the double-sided loss cone. The same double-sided PAD is shown in both panels, from 50° to 60°N latitude, 0° – 20°E longitude. Diffusion equation fits to the $\alpha = 0^{\circ} - 90^{\circ}$ side yield a void size of $72^{\circ+6}_{-27}$, while fits to the $\alpha = 90^{\circ} - 180^{\circ}$ yield a void size of $80^{\circ+1}_{-10}$, and are thus in agreement within the uncertainty.

titude in the $70^{\circ} - 80^{\circ}$ N latitude band where low spacecraft altitudes are available. We do not apply a similar B_0 binning to these averaged PADs as for the surface field calculations. This approach is justified for determining the fluxes because by restricting the range of magnetic field measurements at the spacecraft we may exclude times of high proton fluxes into the magnetosphere and thus would only obtain a lower limit of the fluxes to the surface. We also significantly increase the signal to noise ratio in our averaged PADs with this approach. The proton number density and temperature needed for the derivation of the flux to the surface (as described in detail in Appendix A.2) was calculated in 10° latitudinal bins from FIPS moment estimates as described in Section 5.2. To establish where the double-sided PADs observed in the northern hemisphere map to in the south, we traced the magnetic field line from the northern hemisphere observation point at spacecraft altitudes to the southern hemisphere foot-point location on the surface of the planet.

Observations from the northern hemisphere double-sided loss cones map

to latitudes $20^{\circ} - 30^{\circ}$ S, while observations from the southern hemisphere map to latitudes $20^{\circ} - 40^{\circ}$ S on the surface. MESSENGER cannot observe more equatorward latitudes on the surface in either the northern or southern hemisphere because the high spacecraft altitudes over the equatorial region do not allow it to cross closed field lines that map to the surface near the equator.

We determine the size of the large voids in the double-sided PADs by fitting the solution to the diffusion equation to the PAD. Error bars on the normalized proton flux in the PADs are on average lower for the reflected population, i.e. $\alpha = 0^{\circ} - 90^{\circ}$ for the southern hemisphere observations and $\alpha = 90^{\circ} - 180^{\circ}$ for the northern hemisphere observed PADs. We therefore fit the diffusion curve to the reflected population in each hemisphere. This carries the assumption that the voids on both sides of the PAD essentially equilibrate (i.e. the PAD is symmetric about $\alpha = 90^{\circ}$) after one full bounce period, amounting to the same void size on both sides (as they both correspond to the southern hemisphere surface). We test this assumption on a few example cases for which the normalized fluxes on both sides of the PAD have similar error bars, and find that they are in agreement within the error on the fit α_c (see Figure 5.6).

Figure 5.7 shows the proton flux map to the surface in the southern hemisphere. The fluxes are nearly an order of magnitude higher than the value in the southern hemisphere low latitude band derived in Section 4.5. This discrepancy is due to the nearly single-sided loss cone nature of the PAD in Figure 4.3b, which yields a per second flux to the surface, while the smeared out double-sided loss cones correspond to proton populations that have been trapped for numerous bounce periods and thus yield a bounce-averaged flux to the surface. Thus this is not a valid comparison, as the two estimates are established over very different time-scales. There were two longitudinal/latitudinal grid cells (in the latitude range of $20^{\circ} - 30^{\circ}$ S) that were observed both from northern and southern hemisphere PADs. The flux values estimated from these two methods were in good agreement in the two grid cells; for example, for the longitude range of $80^{\circ} - 100^{\circ}$ E a flux of $1.9^{+0.7}_{-1.3} \times 10^{12}$ particles m⁻² (bounce life-time)⁻¹ was obtained from the


Figure 5.7: Regional-scale resolution of proton flux to the southern hemisphere. a) Proton flux to the southern hemisphere surface on a steregraphic projection plot (looking through the planet from above the north pole). b) Ratio of the error on the flux to the flux estimate. The number of individual PADs averaged in each grid cell was between 50 and 1100.

southern hemisphere measurements and a flux of $(1.9 \pm 0.8) \times 10^{12}$ particles $m^{-2}/(bounce life-time)^{-1}$ was derived from the northern hemisphere PAD.

In the northern hemisphere, we also establish a flux map in 20° longitudinal bins between 70° and 80° N latitudes from observations at less than 550 km altitudes, shown in Figure 5.8. The results for the northern and southern hemispheres are also shown together in a global flux map in Figure 5.9. The global flux map shows particle fluxes in regions where proton-reflection magnetometry directly confirms proton precipitation all the way to the surface (i.e., where loss cones are observed). From Figure 5.9 it can be seen that the fluxes in the north are on average a factor of 2 higher than in the south, as the open field lines provide a direct path for particles entering from the solar wind.

An intriguing result that can be seen from the northern hemisphere flux map in Figure 5.8 is the apparent flux increase in longitude bins near 0° and 180° longitudes compared to the average fluxes observed near $90^{\circ}E$ and $90^{\circ}W$. We attribute this flux increase to Mercury's 3:2 spin orbit resonance,



Figure 5.8: Regional-scale resolution of proton flux to the northern hemisphere high latitude surface. a) Proton flux to the northern cusp surface (looking down from above the north pole). b) Ratio of the error on the flux to the flux estimate. Uncertainties on the fluxes are lowest in the region of interest, near 0° and 180° longitudes, although even in these regions the ratio is still only slightly lower than the signal detected. The number of individual PADs averaged in each grid cell was between 11 and 172.



Figure 5.9: Global map of proton flux to the surface.

which causes the planet's 0° and 180° longitudes to always face the Sun during perihelion or aphelion. These longitudes experience local noon at perihelion where the solar wind density is highest, and thus they receive higher plasma fluxes from the solar wind through open field lines in the cusp (which is on the dayside). A map of the average solar wind density relative to the maximum expected at Mercury's distances from the Sun is shown at different longitudes on the planet in Figure 5.10, with maximum solar wind densities always occurring near 0° and 180° longitudes. As our observations have been averaged over several Mercury years in the 2.5 Earth years included, it is possible for this orbit-averaged signal to be present in our measurements. Such a flux difference would have been present since Mercury entered its 3:2 spin orbit resonance. Higher space weathering rates and greater exospheric generation from these regions is therefore expected. Any differences in the surface elemental composition at these longitudes may introduce different exospheric species compared to the background average. If differences in exospheric species are correlated with surface elemental compositions at these longitudes, the origin of such particles may be traced back to this region due to increased solar wind sputtering in this area.

We expect the plasma flux to be 57% higher in the longitude bins near 0° and 180°, then near longitudes of 90°E and 90°W (Figure 5.10), which face the Sun at approximately the mean orbital distance of Mercury from the Sun. This is because the solar wind density decreases as $1/r^2$ with distance from the Sun, yielding a 57% difference in solar wind density between perihelion and the mean orbital distance of 0.39 AU. Our map of the fluxes shows a signal of ~ 40%, slightly smaller than the expected value. Figure 5.8b shows the ratio of the uncertainty on the flux to the flux estimates for each longitude bin cell for the northern hemisphere (similar plot in Figure 5.7b shows the ratio for the southern hemisphere), where the errors take into account the error on α_c as well as one standard error on the derived n_p and T_p values in each 10° latitude range. The error to flux ratio is lowest in the area of interest and is ~ 30% there, signifying that the flux difference is likely a genuine signal, however, more observations are needed to fully resolve its amplitude and longitudinal extent. A less pronounced, but similar



5.4. Conclusions

Figure 5.10: Predicted relative solar wind density at local noon as a function of longitude, normalized to the maximum value.

signal is also observed in the southern hemisphere fluxes (Figure 5.7). This signal is not significant above the uncertainty level, but may hint at the fact that higher proton fluxes to the southern magnetospheric cusp occurring at perihelion can reach to lower latitudes (where our observations are made) due to particles drifting from open to closed field lines.

5.4 Conclusions

In this chapter, we have extended our original work on proton-reflection magnetometry at Mercury with an additional 1.5 years of observations to resolve regional-scale variation in the surface magnetic field strength as well as in the proton flux to the planet's surface. Our most significant findings can be summarized as follows:

• A new region of proton fluxes in the northern hemisphere is detected, observed at all local times in the latitude range $40^{\circ} - 60^{\circ}$ N, in agree-

ment with plasma pressure measurements at spacecraft altitudes conducted by Korth et al. [136]. The pitch angle distributions observed in this region exhibit double-sided loss cones, similar to those in our southern hemisphere observations.

- The character change in the PADs from the southern hemisphere low latitude region to the northern hemisphere polar region, indicates the boundary between closed field lines at mid to low latitudes and open field lines at high latitudes. We find that this boundary in the northern hemisphere is between $60^{\circ} 70^{\circ}$ N.
- Upper limits for the surface magnetic field strength estimates in the northernmost region are in agreement with the best-fit time-averaged magnetospheric model predictions, with a maximum estimated surface field value from proton-reflection magnetometry of ~ 700 nT and a corresponding model prediction of ~ 750 nT. There are no regions where the surface field strength is stronger than that expected from an offset dipole field with a moment of 190 nT $R_{\rm M}^3$ [19]. This suggests either that higher degree and order core or crustal fields are very weak, or that they occur on length scales much smaller than the resolution of our measurements (~ 300 km).
- Proton fluxes to the surface are estimated for the southern hemisphere in the latitude range 20° 40°S and in the northern hemisphere at latitudes 70° 90°N. Despite the fairly large uncertainties in these measurements, two significant conclusions can be drawn: 1) the per second fluxes everywhere in the northern cusp region are approximately a factor of 2 higher than the bounce-averaged fluxes in the south (which occur over much longer time-scales); 2) increased fluxes are detected in the north near 0° and 180° longitudes compared with fluxes near longitudes of 90°E and 90°W. This may reflect the increased incident solar wind density at these longitudes at local noon that results from the 3:2 spin-orbit resonance of Mercury. Although the signal that we observe is just above the level of the uncertainty, with further obser-

vations it may be possible to better resolve the amplitude of the signal as well as its longitudinal extent.

Finally, new low-altitude observations soon to be acquired by MESSEN-GER in the months leading up to its planned impact into Mercury on 28 March 2015, will allow for the surface magnetic field strength estimates to be tested by orbital MAG observations. These measurements, in conjunction with low altitude proton-reflection magnetometry, may be able to significantly improve the resolution of our current proton-reflection magnetometry estimates at the surface and may reveal higher degree and order short-wavelength structure in the internal field.

Chapter 6

Conclusions

In this final chapter, we summarize key results obtained in this thesis and also list future directions.

In Chapter 2 we explored the nature of the initial interaction between Mercury's global magnetic field and the solar wind and IMF by studying Mercury's bow shock and magnetopause boundaries. The results from this work exemplify that the boundaries of Mercury's magnetosphere, and therefore the magnetosphere itself, are highly variable and greatly affected by the interplanetary medium. Although we do not have continuous solar wind observations upstream of the bow shock, variations in the location of these boundaries can still be perceived on the timescale of a few minutes, as the boundaries are observed to move back and forth across the spacecraft as the solar wind and IMF conditions change. Solar wind conditions are however available on an orbit-by-orbit basis, allowing us to quantify variations in the boundaries on the timescale of hours. The large observed spread in the boundary locations shows that the magnetosphere can expand and contract significantly from one orbit to the next, and we found that this expansion/contraction is controlled dominantly by the solar wind parameters. By building probability density maps of the boundaries, we determined the highest likelihood shape and region for these boundaries on average, and also under various solar wind conditions. We then parametrized the model shape of the magnetopause as a function of the solar wind ram pressure, and the bow shock as a function of the Alfvén Mach number, the two dominant solar wind influences on these boundaries. These parametrized empirical shape models for the magnetopause and bow shock, which are the most important results from this work, can be used to define the boundary conditions for Mercury's magnetosphere and are thus highly relevant for any

magnetospheric study.

A fruitful area of further inquiry that may be addressed with more observations is how reconnection affects the magnetopause location at Mercury. At Earth it has been shown that during times of southward directed IMF. reconnection at the dayside magnetopause erodes magnetic flux on the dayside and transfers it into the tail [57, 58]. This moves the magnetopause planetward and increases the flaring of the magnetotail. We have shown in Chapter 2 that on average Mercury's magnetotail is highly cylindrical; this low flaring implies that magnetic flux has a short residence time in the tail and is transferred back to the dayside much more quickly than our observations can resolve. This observation supports the very short Dungey cycle time at Mercury. However, further observations may be able to inform our understanding of the reconnection related planetward erosion of the dayside magnetopause. It is possible however, that no significant IMF $B_{\rm z}$ dependence will be detected, since at Mercury reconnection has been observed to occur under varying magnetic shear angles, i.e. not just under southward pointing magnetic fields [34]. Resolving this issue would be possible by analyzing many more Mercury years of MESSENGER observations to extract any underlying magnetic shear angle dependence on the magnetopause shape and location. With more observations it would also be possible to derive a more accurate shape model for the magnetopause, which includes asymmetries near the poles due to the cusp regions as well as any other potential asymmetries in the shape causing departures from a figure of revolution, which could not be resolved from the data used in Chapter 2 alone.

In Chapter 3, using magnetic field observations we investigated the region of space where solar wind particles are expected to reach Mercury's magnetosphere, the magnetic cusp region. Due to MESSENGER's eccentric orbit, observations of only the northern hemisphere cusp region are available at this time. Through observations of diamagnetic depressions in the magnetic field on the dayside at high northern latitudes, that were accompanied by high frequency variations in the field indicative of the presence of plasma, we mapped out Mercury's northern cusp region and the associated plasma pressure within it. Prior to MESSENGER, numerous models of Mercury's magnetosphere had made predictions of the latitudinal and local time extent of Mercury's magnetic cusps and therefore the boundary between open and closed field lines (see review in [137]). Since the models did not incorporate the north-south asymmetry of the magnetic field (as this was not known prior to MESSENGER), most models predicted the northern cusp to have the same latitudinal extent as the southern cusp, and also predicted the northern cusp to extend much further south (some as far south as 30° N), than what we observe. Our observations from 3 Mercury years of data under varying solar wind conditions indicate that the limits of the northern cusp are $56^{\circ} - 84^{\circ}$ MSO latitude at spacecraft altitudes and 7 - 16 h in local time.

One of the major results from the work presented in Chapter 3 is that the observed plasma pressure in the northern cusp region is highly affected by the solar wind ram pressure and the IMF B_x direction. We find that the plasma pressure is significantly increased for high ram pressures and for an anti-sunward IMF configuration, which facilitates plasma transport into the northern cusp [107]. An equally significant result is that from the measured plasma pressures in the cusp at spacecraft altitudes we establish the plasma flux down to the surface of the planet. This is found to be in good agreement with the predicted average flux over perihelion and aphelion conditions to the northern hemisphere by Sarantos et al. [110]. Due to the north-south asymmetry of the internal field, we also predict that the particle flux to the surface near the southern hemisphere cusp is a factor of 4 greater than in the north, implying a greater exospheric source from this region and higher space weathering of the surface.

Investigating the affects of magnetopause reconnection on the magnetic cusps as indicated by the IMF B_z and B_y components is an interesting area for further study. By the end of the MESSENGER mission there will be altogether 17 Mercury years of observations of the northern cusp region, which can be used to investigate IMF B_y effects on the local time extent of the cusp that have been observed to occur at the Earth [102, 104]. Zhou et al. [103] showed at the Earth that during times of southward IMF, there is a clear local time shift in the cusp location depending on the sign of IMF $B_{\rm y}$. This is attributed to a shift in the center of the reconnection site by the IMF $B_{\rm y}$ component; positive $B_{\rm y}$ is expected to shift the reconnection site duskward in the northern hemisphere and thus also shift the region of open flux, i.e. the cusp, duskward; the opposite is expected for negative $B_{\rm y}$. As was shown in Chapter 3, the IMF $B_{\rm y}$ component is more stable at Mercury than the $B_{\rm z}$ component, thus such a study focusing on the $B_{\rm y}$ effects on the cusp from 17 Mercury years of observations by MESSENGER is feasible and would improve our understanding of reconnection related affects on Mercury's magnetosphere.

In Chapter 4 of this thesis, we present a new method to quantify the intensity of solar wind proton precipitation to Mercury's surface and make the first measurements of magnetic field strength at the planet's surface using magnetic field and plasma spectrometer observations by MESSENGER. The two most significant findings from this work are: 1) Loss cone observations directly confirm particle precipitation to the surface and show that solar wind plasma persistently bombards Mercury's surface not only in the magnetic cusp regions as expected but over a large percentage of the entire southern hemisphere. (2) The north-south asymmetry in Mercury's long-wavelength magnetic field structure is confirmed at the surface, not just at spacecraft altitudes, independently confirming this unusual and challenging (from a dynamo modeling perspective) feature of the global planetary field and providing key constraints on Mercury's internal dynamo.

The asymmetry in Mercury's magnetic field strength causes hemispheric disparities in the efficiency of solar wind particle penetration into the magnetosphere, and subsequently the surface. The persistent particle precipitation we observe to low southern latitudes on the planet's nightside implies that most of the southern hemisphere is continuously bombarded by plasma, a result that stands in contrast with the canonical view that a global magnetic field protects the surface of an airless body from bombardment by the solar wind. Such precipitation plays a major role in the generation of Mercury's exosphere and space weathering of the surface. This widespread plasma bombardment means that Mercury's surface might show limited latitudinal variation in spectral signatures of space weathering, a result that is in agreement with recent MESSENGER findings [138] but not with predictions from magnetospheric models prior to MESSENGER observations (e.g., [85]). Our results demonstrate for the first time that proton reflection magnetometry can be applied successfully to measure the planetary surface field (providing the only means to do so at Mercury currently) and to quantify plasma precipitation directly to the surface. This technique may prove advantageous alongside electron reflectometry, at other planetary bodies and could be applied at Ganymede, with the upcoming JUICE mission.

Significant wave activity has been observed inside Mercury's magnetosphere [128, 139], which can cause pitch angle diffusion of protons. In future work it would be fruitful to investigate the nature of this scattering process by determining the associated diffusion coefficient, since the diffusion coefficient is dependent on the wave modes involved in the scattering. One possible way to determine the diffusion coefficient is from the power spectrum of the magnetic fluctuation levels [36]; as there are many different wave modes causing the diffusion, the diffusion coefficient will likely be a superposition of different wave power spectra. Anderson et al. [139] have documented varying levels of magnetic fluctuations inside the magnetosphere depending on external conditions. It would thus be interesting to investigate how the diffusion coefficient varies with external conditions, and if any changes in pitch angle diffusion are observed with proton-reflection magnetometry. Once the diffusion coefficient is established, it would be possible to estimate the diffusion timescale associated with the pitch angle diffusion process in Mercury's magnetosphere.

In Chapter 5, we extend our proton-reflection magnetometry work at Mercury with 1.5 years of additional observations to further probe Mercury's surface magnetic field structure and better resolve proton flux precipitation to the surface. The observed transition from double-sided to single-sided loss cones in the pitch angle distributions marks the boundary between open and closed field lines and is shown to occur between 60° and 70° N on the surface, in agreement with the lower latitudinal boundary of the cusp we detected at spacecraft altitudes in Chapter 3. We map all the regions on the surface of the planet in $10^{\circ} \times 20^{\circ}$ latitude/longitude grid cells where proton loss cones are observed; these indicate the regions where proton precipitation directly impacts the surface. Our observations allow for the estimation of surface magnetic field strengths in the northern hemisphere and for the calculation of proton fluxes both to the northern and southern hemisphere.

Most significantly, we find that in the northernmost region, regional-scale variations in the surface magnetic field strength must be either very weak or be on length scales much smaller than the resolution of our observations $(\sim 300 \text{ km})$. The observed increase in magnetic field strength with latitude is consistent with the latitudinal magnetic field variation predicted by an offset dipole field. We also find that the bounce-averaged fluxes observed to the southern hemisphere low latitude region are approximately a factor of two smaller than the instantaneous fluxes estimated to the northern hemisphere cusp region, although these are likely to be an order of magnitude lower instantaneously as was shown in Chapter 4. Mercury's 3:2 spin-orbit resonance is expected to cause a variation in proton fluxes to the surface with body-fixed longitude due to solar wind density changes along Mercury's eccentric orbit. We detect an increase in proton fluxes near 0° and 180° longitudes, consistent with the expected signal; however, since the measured flux increase is only slightly above the uncertainty, further observations are needed to better resolve its amplitude and longitudinal extent. Such a longitudinal signature in proton fluxes to the surface may affect the exospheric species observed at Mercury if there are large-scale longitudinal variations in the surface composition, and is also expected to be accompanied by differential space weathering of the surface in these regions.

Finally, in the last part of its mission phase, MESSENGER will investigate a completely new and unexplored region of Mercury's magnetosphere, that at low altitudes. The low altitude campaign of MESSENGER, which will last approximately 10 months of orbital observations prior to impact into Mercury, will allow measurements to be taken in the northern cusp in the altitude range of 25 - 150 km. These observations are not only crucial from an internal magnetic field perspective, since at these low altitudes the internal field is certainly expected to dominate, but are also vital in terms of understanding the dynamics of the low altitude magnetosphere and particle fluence to the surface. Proton reflection magnetometry alongside vector magnetic field measurements at spacecraft altitudes will be able to resolve regional-scale magnetic field structure that may be present in the northern hemisphere where these low altitudes will be reached. Particle flux measurements at these low altitudes will also validate our estimates of surface fluxes from high altitude observations. These final low altitude observations from MESSENGER will complete the picture of Mercury's northern hemispheric magnetosphere, albeit leaving behind many unanswered questions about the planet's southern hemisphere for future missions, such as BepiColombo.

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Appendix A

Diamagnetic field and particle flux calculation for proton reflection magnetometry

A.1 Diamagnetic magnetic deficit calculation

In this section we discuss the derivation of the diamagnetic field depression at the surface presented in Section 4.5 for the northern cusp and southern hemisphere low latitude band. The diamagnetic field is due to particle motions perpendicular to the magnetic field, and thus the diamagnetic depression can be calculated from the perpendicular particle pressure given by

$$\Delta B = \sqrt{2\mu_0 P_\perp}.\tag{A.1}$$

Thus we need to calculate the perpendicular pressure at the surface to obtain the diamagnetic field at the surface. The pressure is the second moment of the phase space distribution function, or phase space density, and the perpendicular component can be calculated from this as was done in Section 3.4.1. The total pressure can also simply be obtained from the ideal gas law $P = n_p k T_p$ where the proton number density, n_p and temperature, T_p at the spacecraft are needed, and this value can be obtained from FIPS measurements. Thus our aim is to obtain the perpendicular pressure at the surface from the total pressure at the spacecraft altitude, taking into account



Figure A.1: (A) Cartoon diagram showing a loss cone in an idealized pitch angle distribution in the southern hemisphere (A) and northern hemisphere (B).

the observed loss cone from the proton reflectometry measurements.

From Liouville's theorem, the phase space density of particles is preserved along the field line in the absence of particle collisions and loss; here it is also assumed that the particles' total velocity does not change along the field line (an approximation which is needed for proton reflectometry). It is then apparent from equation (3.1) that the total pressure stays constant along the field line in the absence of collisions and loss. As we know from our proton-reflection magnetometry measurements however, particles are being lost to the surface, and so the pressure that we measure at the spacecraft altitude needs to be corrected for this loss to obtain the pressure at the surface.

The proton number density, n_p , obtained from FIPS measurements (see Section 4.5 and [130]) was derived assuming an isotropic proton distribution at the spacecraft altitude, which our observations have shown to not be a good assumption. In reality, the distribution has a loss cone from 0 to α_c in the southern hemisphere, and from α_c to π in the northern cusp region (see Figure A.1 A and B). Having such a void in the pitch angle distribution leads to an underestimate in the phase space density, f_p , because under the isotropy assumption the observed proton number counts by FIPS are spread evenly over 0 to π in pitch angle space, when in fact those particles are only coming from the non-void region of the angle space. Thus, the corrected phase space density should be higher, and we need to account for this enhancement in order to obtain accurate estimates of the proton pressure and flux (as they both depend on the phase space distribution).

The number density is given by the zero-order velocity moment of the distribution function

$$n = \int f(\mathbf{v}) d^3 v \tag{A.2}$$

and in spherical coordinates we have $d^3v = v^2 \sin \alpha d\phi d\alpha dv$. The proton number density, n_p is equal to the integral over the true phase space density, f_{p0} , from α_c to π in the southern hemisphere and 0 to α_c in the northern hemisphere, whereas in the derivation of n_p from FIPS measurements, the integral was implicitly assumed to run from 0 to π over a phase space distribution that is isotropic. Let f_{p0} be the true phase space density, while f_p is the phase space density of an isotropic distribution with density n_p :

$$n_p = 2\pi \int v^2 dv \int_0^\pi \sin(\alpha) d\alpha f_p(v,\alpha) = 2\pi \int v^2 dv \int_{\alpha_c}^\pi \sin(\alpha) d\alpha f_{p0}(v,\alpha),$$
(A.3)

shown here for the southern hemisphere example. If we write $f = FV(v)A(\alpha)$ where V(v) is the speed distribution and $A(\alpha)$ is the pitch angle distribution, and F is a constant, then $V_p(v) = V_{p0}(v)$ holds if the temperatures are the same. The pitch angle dependence can be written as $A_p(\alpha) = 1$ for all α , and $A_{p0}(\alpha) = 0$ for $\alpha = 0$ to α_c and $A_{p0}(\alpha) = 1$ for $\alpha = \alpha_c$ to π in the southern hemisphere (Figure A.1A), while the reverse is true for A_{p0} in the northern hemisphere (Figure A.1B).

From equation (A.3), the ratio F_p to F_{p0} is given by the ratios of the solid angles with non-zero $A(\alpha)$:

$$\frac{F_{p0}}{F_p} = \frac{\int_0^\pi \sin(\alpha) d\alpha}{\int_{\alpha_c}^\pi \sin(\alpha) d\alpha} = \frac{2}{1 + |\cos(\alpha_c)|},\tag{A.4}$$

where the limits of integration are again for the southern hemisphere

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case, but the result with the absolute value sign accounts for both the northern and southern hemisphere cases (since $\alpha_c < \pi/2$ in the southern hemisphere but $> \pi/2$ in the northern hemisphere).

The perpendicular pressure is given by

$$P_{\perp} = \int v_{\perp}^2 f(v, \alpha) dv d\alpha.$$
 (A.5)

For an isotropic distribution it can be shown that $P_{\perp}/P = 2/3$ so that $P_{\perp} = 2/3n_pkT_p$. However, since we don't have an isotropic distribution, the ratio of P_{\perp}/P will be different as we have to consider equation (A.5) integrated over only the range of non-zero $A(\alpha)$. Thus we need to take into account the difference between the phase space distributions, and so we need the factor F_{p0}/F_p .

Because we can obtain the total pressure for the isotropic distribution from FIPS measurements, we take the ratio of the perpendicular pressure for the loss cone distribution to the total pressure for the isotropic distribution

$$\frac{P_{\perp p0}}{P_p} = \frac{\int v_{\perp}^2 f_{p0} dv d\alpha}{\int v^2 f_p dv d\alpha} = \frac{2\pi F_{p0} \int v^4 V(v) dv \int_{\alpha_c}^{\pi} \sin^3 \alpha d\alpha}{2\pi F_p \int v^4 V(v) dv \int_0^{\pi} \sin \alpha d\alpha},$$
(A.6)

where the limits of the integration in the numerator are for the southern hemisphere case, but the results below apply to both cases. Using equation (A.4), we arrive at

$$P_{\perp p0} = \left[1 - \frac{1}{3} \left(\frac{1 + |\cos^3 \alpha_c|}{1 + |\cos \alpha_c|}\right)\right] P_p.$$
 (A.7)

This result is the ratio at the spacecraft altitude, but the actual magnetic field due to the plasma at the surface must consider the perpendicular plasma pressure at the surface. To calculate the perpendicular pressure at the surface, we need to map the true phase space density at the spacecraft f_{p0} to the surface. By Liouville's theorem, f_{p0} remains constant along phase space trajectories, so we use the same f_{p0} at the surface to evaluate the pressure. Note however that if we assume that the surface is perfectly absorbing, then there will be no upward going particles and the only integrate over the downgoing particles. That is, at the surface we have

$$\frac{P_{\perp p0,surf}}{P_p} = \frac{2\pi F_{p0} \int v^4 V(v) dv \int_{\pi/2}^{\pi} \sin^3 \alpha d\alpha}{2\pi F_p \int v^4 V(v) dv \int_0^{\pi} \sin \alpha d\alpha},\tag{A.8}$$

which yields

$$P_{\perp p0,surf} = \frac{2}{3(1+|\cos\alpha_c|)}P_p \tag{A.9}$$

where the absolute value sign again takes care of both the northern and southern hemispheres.

Because the system is not infinite, but has gradients in the magnetic field between the spacecraft and the surface, the simple expression for ΔB in equation (A.1) is not precise and the actual ΔB should be somewhere between that indicated at the spacecraft altitude and that at the surface. Substituting in values for α_c , n_p , and T_p in equation (A.9) and taking ΔB to be between $\sqrt{2\mu_0 P_{\perp p0}}$ and $\sqrt{2\mu_0 P_{\perp p0,surf}}$, we arrive at the results listed in Section 4.5 for the diamagnetic field.

A.2 Particle flux calculation

As we have more information at our disposal from the loss cone size than what we had available in Section 3.4.1, we can do a more accurate calculation of the proton surface flux in the northern hemisphere cusp region as well as calculate the surface flux in the southern hemisphere low latitude band. The flux of particles per unit area per unit time with velocity parallel to the magnetic field line v_{\parallel} through a surface normal to the magnetic field is given by

$$\Phi = \int f_{p0}(\mathbf{v}) v_{\parallel} d^3 \mathbf{v}, \qquad (A.10)$$

where we are using the parallel velocity $(v_{\parallel} = v \cos(\alpha))$ because we want the flux through the surface, thus we need the velocity that is parallel to the surface normal. This assumes that the field lines are perpendicular to the surface right at the surface. We assume gyrotropy again, such that the integral over $d\phi = 2\pi$ and use the same decomposition for the phase space density as in Section A.1. We also need to ensure that the flux is obtained from the corrected, or true, phase space density f_{p0} as we established in Section A.1. The flux can then be rewritten as

$$\Phi = 2\pi \int v^3 V_{p0}(v) dv \int F_{p0} A(\alpha) \cos(\alpha) \sin(\alpha) d\alpha, \qquad (A.11)$$

where the integral over α runs from 0 to α_c in the northern hemisphere, and α_c to π in the southern hemisphere. This integral will yield an upper limit for the flux because by integrating over these values in pitch angle we are assuming that the loss cone remains completely filled in for the downgoing (incident) particles and that all of the particles within the loss cone are absorbed at the surface.

 F_{p0} is simply a constant and can be written in terms of F_p as from equation (A.4) above. We assume as before that $V_{p0} = V_p$, and we can then rewrite equation (A.11) as

$$\Phi = 2\pi \int v^3 V_p(v) dv \int F_p \frac{2}{1 + |\cos(\alpha_c)|} A(\alpha) \cos(\alpha) \sin(\alpha) d\alpha, \quad (A.12)$$

which due to $f_p(v) = F_p V_p(v)$ (where we have assumed that the phase space density has no angular dependence as before since $A(\alpha) = 0$ or 1) can be simplified to be

$$\Phi = 2\pi \frac{2}{1 + |\cos(\alpha_c)|} \int v^3 f_p(v) dv \int_{\alpha_c}^{\pi} \cos(\alpha) \sin(\alpha) d\alpha$$
(A.13)

for the southern hemisphere integration limits. In order to solve equation (A.13), we need to assume a Maxwellian velocity distribution:

$$f(v) = n_p \left(\frac{m_p}{2\pi kT_p}\right)^{3/2} \exp\left(-\frac{m_p v^2}{2kT_p}\right), \qquad (A.14)$$

where n_p is the number density quoted in Section 4.5. Substituting all

this into equation (A.13), we find

$$\Phi = 2\pi n_p \left(\frac{m_p}{2\pi kT_p}\right)^{3/2} \frac{2}{1+|\cos(\alpha_c)|} \int_0^\infty v^3 \exp\left(-\frac{mv^2}{2kT_p}\right) dv \int_{\alpha_c}^\pi \cos(\alpha) \sin(\alpha) d\alpha,$$
(A.15)

which yields

$$\Phi = n_p \frac{\sin^2 \alpha_c}{(1 + |\cos \alpha_c|)} \sqrt{\frac{2kT_p}{\pi m_p}}.$$
(A.16)

Due to the absolute value sign, this equation is again valid for both the southern and northern hemispheres. Substituting in values for the parameters yields the results discussed in Section 4.5.