

**Imperial College
London**

The Dynamics of the Outer Boundaries in Global Simulations of Planetary Magnetospheres

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Submitted in part fulfillment of the requirements for the degree of
Doctor of Philosophy

30th September 2018

Abstract

The solar wind — magnetosphere interaction consists of a set of coupled regions with different dominating physical processes. Understanding this interaction is important in both furthering our understanding of space plasma physics, and mitigating the effects of space weather. The outer boundaries of the magnetosphere are important as they mediate variations in the solar wind and their impact on magnetospheric dynamics.

The Gorgon magnetohydrodynamic code is developed to simulate planetary magnetospheres. It is benchmarked in the context of the Earth's magnetosphere, where it performs well compared to other global MHD codes. The code is used to study the behaviour of the outer boundaries of the magnetosphere, the bow shock and magnetopause, focussing on their response to different magnetospheric drivers.

Gorgon is first adapted to Neptune's magnetosphere, whose dipole is significantly tilted from its rotation axis. The outer boundaries undergo daily variation, resulting in reconnection being heavily modulated by the precessing internal field. The results found have implications on the global circulation of plasma and energy in inclined magnetospheres, emphasising the need for three-dimensional, time-dependent simulations.

The Earth's magnetosphere is studied using real varying solar wind, where the outer boundaries' position and velocity undergo rapid changes. The variable solar wind and curved bow shock causes different parts of the bow shock to respond differently which not captured in empirical models that are widely used to place spacecraft observations in context. In situ observations of the shock may be biased toward observing high velocity shocks, as they rarely cross a stationary shock. A complex mesh of flux ropes are also formed. Their generation creates complex topology, which affects their subsequent evolution, allowing some flux ropes far into the tail. Their complex topology has interesting consequences on the impact these flux ropes have on the transfer of magnetic flux and plasma into magnetosphere.

Declarations

I declare that the work within this thesis is my own, and where other work has been used, it has been appropriately referenced.

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Lars Mejnertsen, 11th October 2018

The work presented in this thesis has contributed to two peer-reviewed published papers:

- L. Mejnertsen et al. (2016). “Global MHD simulations of Neptune’s magnetosphere”. In: *Journal of Geophysical Research: Space Physics* 121.8, pp. 7497–7513. ISSN: 21699380. DOI: **10.1002/2015JA022272**
- L. Mejnertsen et al. (2018). “Global MHD Simulations of the Earth’s Bow Shock Shape and Motion Under Variable Solar Wind Conditions”. In: *Journal of Geophysical Research: Space Physics* 123.1, pp. 259–271. ISSN: 21699380. DOI: **10.1002/2017JA024690**

Acknowledgements

I would first like to thank my supervisors, Jonathan Eastwood and Jerry Chittenden. Having worked with Jonathan and Jerry since my MSci project almost 5 years ago, I can say it has been a pleasure to work with them, and I would like to thank them for their guidance, support and patience. I would also like to thank my collaborators, Adam Masters, Heli Hietala and Steven Schwartz for their invaluable contributions into the work in this thesis and the published papers.

These four years at Imperial have been made a lot more enjoyable thanks to the wonderful students and staff in the Space and Atmospheric Physics Group, who provided many interesting discussions and entertaining times.

Many thanks to my friends and family for their continued support and enjoyable distractions. I would like to thank my parents especially for their support and encouragement throughout my life: I would not have made it this far without them. Finally, I would like to thank Paula, for sticking by me through all the ups and downs of doing a PhD, supporting me all the way, and always keeping me smiling.

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Nomenclature

ACE	Advanced Composition Explorer
AU	Astronomical unit, 1.50×10^8 km
BS	Bow Shock
k_B	Boltzmann's Constant, 1.38×10^{-23} J/K
q	Charge
ρ_c	Charge Density
$\vec{J}$	Current density
DSCOVR	Deep Space Climate Observatory
R_E	Earth Radius, 6.38×10^3 km
m_e	Electron mass, 9.11×10^{-31} kg
e	Elementary charge, 1.6×10^{-19} C
ϵ	Energy density
FTE	Flux Transfer Event
GSE	Geocentric Solar Ecliptic coordinate system
GMHD	Global magnetosphere simulations using MHD
γ	Ideal gas constant, (5/3 for an ideal gas)
IMF	Interplanetary Magnetic Field
R_M	Magnetic Reynolds number
$\vec{B}$	Magnetic field strength
$\vec{A}$	Magnetic vector potential, $\vec{B} = \nabla \times \vec{A}$
P_B	Magnetic pressure
MHD	Magnetohydrodynamics
MP	Magnetopause
ρ	Mass density
$\vec{p}$	Momentum density

R_N	Neptune radius, 24.6×10^3 km
NSO	Neptune Solar Orbital coordinate system
n	Number density
μ_0	Permeability of free space, $4\pi \times 10^{-7}$ H/m
ϵ_0	Permittivity of free space, 8.85×10^{-12} F/m
β	Plasma beta
m_p	Proton mass, 1.67×10^{-27} kg/m ³
PE	Prediction Efficiency
sw	Solar wind
T	Temperature
P	Thermal pressure
$\vec{v}$	Velocity

1 Introduction

A magnetosphere is the extension of a planet's magnetic field into interplanetary space, where it encounters the solar wind — an ionized plasma streaming from the Sun at supersonic speeds — along with the embedded interplanetary magnetic field. This interaction is complex and spans over many length and time scales, and falls within the realm of collisionless plasma physics. The study of planetary magnetospheres is interesting for two reasons. From a plasma physics perspective, they provide a parameter regime unique to interplanetary space. The collisionless nature of the plasmas within the solar wind-magnetosphere coupling system provides interesting phenomena which are not fully understood, such as collisionless shocks, magnetic reconnection, and plasma turbulence. These conditions cannot be easily re-created in the laboratory. Secondly, the planetary magnetic field largely protects the planet from the solar wind, by deflecting it around and past the planet. Under certain circumstances, the solar wind can enter into the magnetosphere, potentially damaging infrastructure and communication networks. This is the manifestation of space weather. Understanding the process by which space weather affects the Earth is important to mitigate its adverse effects. In order to answer such questions, a multitude of spacecraft have been sent to orbits around the Earth, the Sun and the other planets, as well as beyond the solar system. They observe both the local plasma and magnetic field parameters and remotely image phenomena of interest such as the Sun and planetary aurora.

Although spacecraft observations are an integral part of space physics today, the field much pre-dates the space age. Sightings of the aurora, with their bright, colourful displays in the sky, have long fascinated people over the centuries, naturally producing speculation as to their origin. It was only with the invention of compasses in the 1700s, and later magnetometers and telescopes, that more scientific studies could be performed. These suggested that deviations in the Earth's magnetic field, activity on the Sun, and the presence of aurora are all correlated, leading to speculation that the aurora were due to energetic particles from the Sun (Birkeland,

1901). This is now known as the solar wind, whose existence has been postulated as early as 1943, where it was noted that the curved tail of a comet could be due to the presence of the solar wind (Biermann, 1951). It was not fully observed until the Mariner 2 spacecraft first crossed the bow shock in 1963 (Neugebauer and Snyder, 1966). The link between the solar wind and the Earth's magnetic field led to the idea of the 'magnetosphere' (Gold, 1959), the region of interplanetary space where the planet's intrinsic field dominates.

Space weather is the generic term for the effect solar activity has on health and infrastructure at Earth. In an ever-increasing digital world, it has the potential to cause significant damage to a whole host of systems, both space and terrestrial (Cannon et al., 2013). While adverse space weather conditions, such as the arrival of a large coronal mass ejection (Tuija I. Pulkkinen, 2007), can create beautiful auroral displays, the currents generated in the ionosphere have the capability to do lasting damage to power grids by overloading transformers (Boteler, Pirjola, and Nevanlinna, 1998; Juusola et al., 2015; Kappenman, 1996) and underground pipelines (Boteler, 2000), and disrupt communications vital for aircraft operations and the Global Positioning Service (Cannon et al., 2013; Kintner, Ledvina, and Paula, 2007). Orbiting satellites can also suffer effects ranging from temporary disruption to permanent damage, as the magnetosphere which usually protects them can be compressed to within geostationary orbit or energise radiation belt particles (J. P. Eastwood et al., 2017). These potential impacts have already been realised: the most high profile case occurring in Quebec, 1989, where adverse space weather caused a significant power outage (Bolduc, 2002). This has led governments to state an interest in being able to predict such events, with the UK government placing a likelihood of 5% to 50% of a severe event occurring in the next five years, with a medium impact on infrastructure (*National Risk Register of Civil Emergencies* 2015; Cannon et al., 2013).

Global simulations of the magnetosphere offer a complementary approach to spacecraft observations for improving both our understanding of space plasmas and space weather prediction capabilities. As the name implies, they are able to simulate the entire magnetosphere, capturing the global processes required to understand magnetospheric dynamics as a whole. Global simulations allow us to capture the entire magnetospheric system at an instant in time, and thus reveal the interconnection between different plasma regions and dynamics on large scales.

In this chapter, the basic concepts of magnetohydrodynamics and the magnetospheric system will be described. Magnetohydrodynamics is a well developed theory and can explain many large-scale magnetospheric phenomena. However, significant simplifications in plasma physics

are made in the derivation of magnetohydrodynamics; part of this chapter will discuss the implications of this when applied to the magnetosphere, and for accurate reproduction of observational results. The solar wind-magnetosphere system is described too, from the origin of the solar wind at the Sun, to its effect on planetary magnetospheres. Finally, the simulating of the magnetosphere globally using the magnetohydrodynamic equations will be detailed. There are a variety of competing codes which have been used for a number of different studies; these are also critically reviewed. The chapter concludes with a brief overview of some of the open questions relating to the physics of the outer magnetospheric boundaries.

1.1 Space Plasma Physics

A plasma is a collection of positively and negatively charged particles, which flow freely and respond to electric and magnetic forces. Their overall charge distribution is zero, acting in some ways as a neutral gas. However, each individual particle interacts with the electromagnetic fields generated by the particles, and thus plasmas exhibit complex collective behaviour. This material is covered in many textbooks, for example Baumjohann and Treumann (2012a) and Margaret G Kivelson and Christopher T Russell (1995) but is included here for completeness.

To define a more quantitative estimate of the parameter regime which plasmas inhabit, an important length scale is the Debye length: at scales smaller than this length, the plasma can no longer be considered neutral. This defines the concept of quasi-neutrality of a plasma. Thermal motions of individual particles can cause local charge imbalances, which are generally unstable and create waves within the plasma in order to restore neutrality. In a plasma of temperature, T , and electron number density, n_e , the Debye length, λ_D , is given as,

$$\lambda_D = \sqrt{\frac{\epsilon_0 k_B T}{n_e e^2}}, \quad (1.1)$$

where $\epsilon_0 = 8.85 \times 10^{-12}$ F/m is the electric permittivity of free space, $k_B = 1.38 \times 10^{-23}$ J/K is Boltzmann's constant, and $e = 1.6 \times 10^{-19}$ C is the elementary charge. It is derived from the fact that each particle is subject to the electrostatic attraction/repulsion of its neighbours, whose magnitude is proportional to the inverse of the distance squared. However, since there are an equal number of positive and negative charges, the surrounding particles will also provide a shielding of this force, reducing it drastically over large distances. This is known as Debye

shielding, and attributes the characteristic size of this shielding to the Debye length, shown in equation (1.1).

Derived from the Debye length is the plasma parameter, Λ , which quantifies how effective Debye shielding is. It is a measure of how many particles are within a Debye sphere (a sphere of radius λ_D), quantified in equation (1.2). In order for the shielding to be effective, there must be enough particles in the sphere to shield the Coulomb interaction.

$$\Lambda = n_e \lambda_D^3 \quad (1.2)$$

The final important scale is the plasma frequency, ω_{pe} . It is the timescale over which quasi-neutrality is restored if the charge distribution is displaced, allowing a non-zero net charge to exist within the plasma. Due to the charge imbalance, an electric field is formed which creates a restoring force proportional to the initial displacement. This is a classic case of simple harmonic motion, with an associated frequency: in this case, the plasma frequency, specified in equation (1.3), showing specifically the proton-electron plasma frequency. The restoring force affects the electrons to a greater extent than the ions, since they have a much lower mass, and hence inertia, than ions.

$$\omega_{pe} = \sqrt{\frac{n_e e^2}{m_e \epsilon_0}} \quad (1.3)$$

The plasma approximation is summarised by equations (1.4) to (1.7). The first equation (1.4) states that there must be an equal number density of positive and negative charges, enforcing quasi-neutrality on a large scale. Within the plasma, typical length scales, L , must be larger than the Debye Length (equation (1.5)), and there must be a high enough particle density (equation (1.6)) in order to maintain quasi-neutrality and free flowing particles. Lastly, the plasma frequency must be much larger than the neutral collision timescale, τ_n (equation (1.7)). Collisions between neutral and charged particles forces the two species into equilibrium, causing the plasma to behave more like a gas.

$$n_i \approx n_e \quad (1.4)$$

$$\lambda_D \ll L \quad (1.5)$$

$$\Lambda \gg 1 \quad (1.6)$$

$$\omega_{pe} \tau_n \gg 1 \quad (1.7)$$

Collections of particles with these properties are considered to be a plasma. They evolve according to the individual motions of the particles due to their interaction with the electromagnetic field. A charged particle in an electromagnetic field moves according to the Lorentz force,

$$m \frac{d\vec{v}}{dt} = q \left(\vec{E} + \vec{v} \times \vec{B} \right), \quad (1.8)$$

where m , $\vec{v}$ and q are the particle's mass, velocity, and charge, and $\vec{E}$ and $\vec{B}$ are the electric and magnetic fields. However, the motion and position of these charges generates and modifies the electromagnetic fields. This is described by Maxwell's equations,

$$\nabla \cdot \vec{E} = \frac{\rho_c}{\epsilon_0} \quad (1.9)$$

$$\nabla \times \vec{E} = -\frac{\partial \vec{B}}{\partial t} \quad (1.10)$$

$$\nabla \cdot \vec{B} = 0 \quad (1.11)$$

$$\nabla \times \vec{B} = \mu_0 \vec{J} + \mu_0 \epsilon_0 \frac{\partial \vec{E}}{\partial t} \quad (1.12)$$

where the plasma generates changes in the electromagnetic field by its charge density, ρ_c , and current density, $\vec{J}$; $\mu_0 = 4\pi \times 10^{-7} \text{H/m}$ is the permeability of free space. It consists of four vector equations: Gauss' Law in equation (1.9), Faraday's Law in equation (1.10), Gauss' Law for magnetism (otherwise known as divergence free condition) in equation (1.11) and the Ampère-Maxwell Law in equation (1.12).

In order to capture the full dynamics of a plasma, both Newton's second law (equation (1.8)) for every particle and Maxwell's equations (equations (1.9) to (1.12)) must be solved simultaneously. This is not feasible considering the number of particles present in space plasmas: for example, the solar wind near Earth, a relatively rarefied space plasma, has a particle number density typically $\sim 10^6 / \text{m}^3$.

Fortunately, a number of simplifications can be made to render the modelling more feasible. In what is called Single Particle Motion, insight can be gained by considering just one particle and its movement due to electromagnetic fields. More self-consistent approaches include kinetic theory and magnetohydrodynamics, which solve both solutions for the plasma and the electromagnetic fields simultaneously.

1.1.1 Single Particle Motion

To gain insight into the micro-scale behaviour of plasma, it is useful to consider the dynamics of individual particles in electromagnetic fields, otherwise known as Single Particle Motion. Each individual particle is subject to the Lorentz force (equation (1.8)), but their feedback into the electromagnetic fields is neglected.

Moving charged particles in a uniform magnetic field are subject to a centripetal force perpendicular to the magnetic field, causing circular motion or gyration called cyclotron motion. They do not experience a force along the magnetic field direction, meaning the particles can stream relatively freely up and down field lines, whilst gyrating.

The frequency and radius of these cyclotron orbits, called the cyclotron frequency (Ω_c) and Larmor radius (r_L) respectively, are dependent on the mass, charge and energy of the particle, as given in equations (1.13) and (1.14),

$$\Omega_c = \frac{qB}{m} \quad (1.13)$$

$$r_L = \frac{v_\perp}{\Omega_c} = \frac{mv_\perp}{qB} \quad (1.14)$$

It provides a distinction between the different charges which constitute the plasma. The light, negatively charged electrons will move in small circular paths at high frequency, and in the opposite direction to the relative large positive ions, whose orbits are large and low frequency.

In the presence of electric and non-uniform magnetic fields, particles undergoing cyclotron motion will tend to drift. Drifting behaviour can be explained by considering the change in the particle's Larmor radius due to changes in the background magnetic field strength (e.g. gradient drift), or in the particle's energy (e.g. $\vec{E} \times \vec{B}$ drift).

1.1.2 Kinetic Theory

In kinetic theory, the plasma is described by a six dimensional distribution function, $f(\vec{x}, \vec{v})$, three dimensions in space and three in velocity. Macro-scale quantities, such as number density, are found by integrating f over different volume elements. The evolution of the distribution function, in the absence of collisions, is described by the Vlasov equation,

$$\frac{\partial f}{\partial t} + \vec{v} \cdot \nabla_{\vec{x}} f + \frac{q}{m} \left(\vec{E} + \vec{v} \times \vec{B} \right) \cdot \nabla_{\vec{v}} f = 0 \quad (1.15)$$

where $\nabla_{\vec{x}}$ describes the gradients in space, $\vec{x}$, and $\nabla_{\vec{v}}$ describes the gradients in velocity space, $\vec{v}$.

Kinetic theory describes accurately some of the interesting features that are not found in neutral gasses. In a neutral gas, particle collisions create Maxwellian velocity distributions. Since space plasmas are collisionless, this does not occur, allowing non-Maxwellian distributions. Furthermore, due to the presence of magnetic fields, there may be an anisotropy between velocities perpendicular and parallel to the magnetic field.

Solving the Vlasov equation is not a trivial task. Most known analytical solutions make a number of assumptions, and solving them numerically is computationally prohibitive. Kinetic theory includes physics which occurs over small length scales (inertial lengths) and timescales (plasma frequencies), which need to be resolved in any numerical simulation (Büchner, 2005). Scaling this up to simulating the magnetosphere, in three spatial and three velocity dimensions, requires a vast amount of computing resources. A number of simplifications can be made to lessen computational cost: the Vlasiator, a hybrid-Vlasov code of the Earth's magnetosphere, solves the Vlasov equation for ions and treats electrons as a charge neutralising fluid (Alfthan et al., 2014; Pfau-Kempf et al., 2016). However, this code still requires a large amount of computing resources, and has yet to be run in the full six dimensions (currently only running with two spatial and three velocity dimensions).

1.2 Magnetohydrodynamics

Magnetohydrodynamics (MHD) takes a fluid approach to modelling plasmas. The plasma is treated as an electrically conductive fluid, describing the evolution of average fluid parameters such as density, velocity, and temperature and the embedded electromagnetic fields.

Many assumptions have been made to simplify the physics governing plasma motions to generate the MHD equations. MHD treats the plasma population as a single fluid; it does not distinguish between the different components of the plasma, such as the ions and electrons. It also neglects 'fast' timescales associated with electromagnetic wave propagation and the restoration of quasi-neutrality. By taking a fluid approach, it also assumes all particles follow a Gaussian distribution, which is not necessarily true in plasma that is collisionless and the directionality imposed by the presence of magnetic field. Fortunately, within the magnetosphere,

the Gaussian approximation holds more or less true, except in areas such as the ring current and the plasma sheet (Joachim Raeder, 2003).

In general, the time evolution of an electrically charged fluid and the embedded magnetic field can be described by the resistive MHD equations (1.16) to (1.20),

$$\frac{\partial \rho}{\partial t} + \nabla \cdot \rho \vec{v} = 0, \quad (1.16)$$

$$\rho \left(\frac{\partial \vec{v}}{\partial t} + (\vec{v} \cdot \nabla) \vec{v} \right) = -\nabla P + \vec{J} \times \vec{B}, \quad (1.17)$$

$$\frac{d}{dt} \left(\frac{P}{\rho^\gamma} \right) = 0 \quad (1.18)$$

$$\frac{\partial \vec{B}}{\partial t} = -\nabla \times \vec{E}, \quad (1.19)$$

$$\text{where, } \vec{E} + \vec{v} \times \vec{B} = \eta \vec{J} \quad (1.20)$$

The quantities ρ , $\vec{v}$ and P are the mass density, velocity and thermal pressure of the fluid, with an ideal gas constant, γ . $\vec{J}$ and $\vec{B}$ are the current density and the magnetic field strength, and η is the magnetic resistivity.

The first three equations (equations (1.16) to (1.18)) are the inviscid Navier-Stokes equations, which describe how a fluid evolves when subjected to internal pressures and external forces. Equation (1.16) describes the property of mass conservation: for a given fixed position in the fluid, the rate of change of mass is equal to the sum of the mass flowing in and out of that position. Equation (1.17) describes Newton's second law, otherwise known as the momentum equation. Here, the momentum density is changed in the presence of forces, such as pressure gradients, ∇P , and the Lorentz force, $\vec{J} \times \vec{B}$, and also due to the amount of momentum advected with the fluid, $\rho (\vec{v} \cdot \nabla) \vec{v}$. Finally, equation (1.18) is the adiabatic energy equation for an ideal gas, where $P = (n_e + n_i)k_B T = 2nk_B T$ (as $n_e \simeq n_i \simeq n$).

Since MHD averages over the ion and electron motion, the current density, $\vec{J}$, is not calculated by the definition of current, $\vec{J} = e (n_i \vec{v}_i - n_e \vec{v}_e)$. Instead, Ampère's law is used, as given in equation (1.21). The displacement current is neglected as it is negligible in most space plasmas — the plasma varies on time-scales much slower than electromagnetic wave propagation — and the usual MHD equations do not consider the propagation of electromagnetic waves.

$$\mu_0 \vec{J} = \nabla \times \vec{B} - \frac{1}{c^2} \frac{\partial \vec{E}}{\partial t} \simeq \nabla \times \vec{B} \quad (1.21)$$

Table 1.1: Typical values of plasma beta in the magnetosphere system. ¹From Margaret G Kivelson and Christopher T Russell (1995); ²From P. J. Cargill et al. (2005).

Region	n (/cm ³)	T (K)	$ \vec{B} $ (nT)	β
Solar Wind at 1 AU ¹	6.6	1.2×10^5	7	~ 1
Magnetosheath ²	26.4	2.4×10^6	28	~ 10
Plasmasphere at $3R_E$ ¹	10^3	1.2×10^5	1100	$\sim 10^{-3}$
Cusps ²	10	2.5×10^6	240	$\sim 10^{-2}$

In equation (1.22), Ampère's law (equation (1.21)) is substituted into the $\vec{J} \times \vec{B}$ force. The force due to the magnetic field is revealed to have two components,

$$\vec{J} \times \vec{B} = -\nabla \frac{B^2}{2\mu_0} + \frac{1}{\mu_0} (\vec{B} \cdot \nabla) \vec{B} \quad (1.22)$$

The first term on the right hand side has the same functional form as the force due to thermal pressure gradients, and is hence known as the magnetic pressure. It represents the tendency for compressed magnetic fields to expand. The second term represents the magnetic tension. Analogous to a string under tension, this force acts to reduce the curvature of the magnetic field.

The ratio between the magnetic pressure and the thermal pressure is an important dimensionless parameter, useful in understanding the plasma regime. Known as the plasma β (as in equation (1.23)), it gives a sense of what dominates the motion of the plasma, whether it is largely driven by thermal pressures, or whether the $\vec{J} \times \vec{B}$ force dominates.

$$\beta = \frac{P}{P_B} = \frac{2nk_B T}{B^2/2\mu_0} \quad (1.23)$$

β can vary drastically depending on different regions. In the solar wind, for example, the β is typically 1, though in the inner magnetosphere where the magnetic field strength increases with a $1/r^3$ relationship, β decreases such that the plasma is dominated by magnetic forces. Table 1.1 shows some typical plasma beta values found in the magnetosphere.

1.2.1 The Magnetic Induction Equation

The final equation (1.19) is Faraday's Law and describes the evolution of the magnetic field. This requires determination of the electric field using Ohm's law, given in equation (1.20), where $\eta \vec{J}$ represents the electric field due to the current, with resistivity η , and $\vec{v} \times \vec{B}$ is the motional electric field. Using Ohm's law (equation (1.20)) to substitute for $\vec{E}$ in Faraday's law (equation (1.19)), and the simplified Ampère's law (equation (1.21)) for $\vec{J}$, the magnetic induction equation is obtained,

$$\frac{\partial \vec{B}}{\partial t} = \nabla \times (\vec{v} \times \vec{B}) - \eta \nabla^2 \vec{B} \quad (1.24)$$

The change in the magnetic field is shown to have two components: the advection of magnetic field with the plasma through $\nabla \times (\vec{v} \times \vec{B})$, and the diffusion of the magnetic field through $\eta \nabla^2 \vec{B}$.

The ratio between these two terms is another important dimensionless parameter in MHD: the magnetic Reynolds number,

$$R_M = \frac{|\nabla \times (\vec{v} \times \vec{B})|}{\left| \frac{\eta}{\mu_0} \nabla^2 \vec{B} \right|} \sim \frac{\mu_0 L v}{\eta}. \quad (1.25)$$

The magnetic Reynolds number, R_M , is a measure of the amount of advection compared to the amount of diffusion over a typical length scale, L . This gives insight into how magnetic fields evolve with time, whether they move with the plasma flow ($\nabla \times (\vec{v} \times \vec{B})$ dominates) or diffuse through the plasma ($(\eta/\mu_0) \nabla^2 \vec{B}$ dominates).

The resistivity, η , is usually determined by the rate of collisions between the particles constituting the plasma. In collisional plasmas, this is usually calculated through the Spitzer resistivity, which quantifies the Coulomb interaction. However, in most space plasmas, the Coulomb collision frequency is negligibly small (Baumjohann and Treumann, 2012a). Instead, complex wave-particle interactions can cause electrons to scatter, creating an anomalous resistivity (Manfred Scholer, 1988)

1.2.2 Ideal Magnetohydrodynamics and the Frozen-in Condition

Most space plasmas are so rarefied that the average collision timescale is much longer than the relevant MHD timescales. This means that the effect of collisions can be neglected, leading to

zero resistivity. This simplifies the equations to Ideal Magnetohydrodynamics.

With no resistivity, the magnetic field does not diffuse through the plasma: it can only advect with the fluid. This gives an interesting consequence; in a perfectly conducting plasma, with $\eta = 0$, the magnetic flux through any open surface moving with the plasma remains unchanged. It ‘associates’ the magnetic field with the plasma: the motion of plasma can drag the magnetic field, and vice versa. This is known as the frozen-in approximation.

This concept of frozen-in flow explains many aspects of space plasma physics. For example, it allows the solar wind to carry the Sun’s magnetic field with it into interplanetary space. It also explains why the magnetosphere presents an obstacle for the solar wind flow: the plasma is frozen-in and tied to magnetic field lines, which do not merge with the planetary field under normal circumstances. It also gives rise to the notion of magnetic field lines being tied to the plasma it is embedded within. Plasma on a particular field line are said to be topologically linked, due to the aforementioned association with the magnetic field.

1.2.3 Magnetic Reconnection

An important process in space plasmas is magnetic reconnection, which plays a significant role in energising plasmas and changing their topology within the magnetosphere. Macroscopically, reconnection restructures the magnetic field to a lower energy state, converting the stored magnetic energy into high velocity and temperature plasma jets (Yamada, Kulsrud, and Ji, 2010).

Reconnection occurs when two separate frozen-in regions are compressed by plasma flow, as shown in figure 1.1. Embedded within these two regions are anti-aligned magnetic fields, which are also compressed together, forming a current layer in between. This current layer is very thin, on length scales where the frozen-in flux condition no longer applies. This allows the magnetic field to ‘reconnect’ through the plasma, as seen in figure 1.1.

The microscopic mechanism of reconnection is not captured in MHD, as it occurs on spatial scales where the individual motions of the plasma species must be taken into account, and the frozen-in flux theorem no longer applies. For magnetic reconnection to occur in standard MHD, there must be a non-zero resistivity. In MHD the energy transfer in reconnection is captured in the $\vec{J} \times \vec{B}$ force, accelerating the plasma by the large current in between the two regions. The

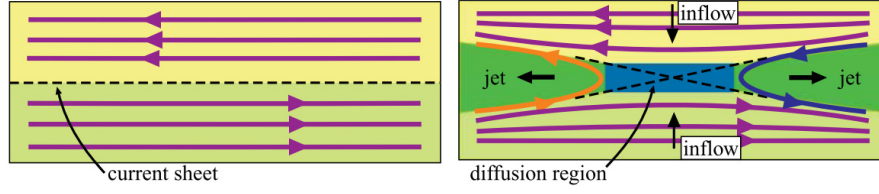


Figure 1.1: A two-dimensional picture of magnetic reconnection. Plasma compresses the anti-parallel magnetic fields, creating a diffusion region. The field reconnects, and the energy release heats and transports the plasma away from the diffusion region. Image from: Jonathan P. Eastwood (2008)

reconnection region is typically very small, allowing diffusion to overcome advection, $R_M \ll 1$. The field then disassociates itself with the plasma, allowing it to reconnect.

In general, any extra term added to the $\vec{v} \times \vec{B}$ term on the right hand side of equation (1.24) can cause the frozen-in approximation to break down. Taking into account a more complex form of Ohm's law, for example, adds extra terms which are not dependent resistivity but will still break the frozen-in approximation. This will be discussed in a later section on extended MHD. In global MHD simulations, numerical resistivity is often the cause of reconnection. Numerical resistivity is due to the fact that the equations are discretised both spatially and temporally. This produces error in the calculation of the MHD variables, and hence adds extra terms in the magnetic induction equation. A simplified example of this can be shown by taking a simplified discretisation of the induction equation (1.24) spatially. Taking only the $\hat{x}$ component of $\vec{B}$ and applying a central difference scheme, Faraday's law is given by,

$$\frac{\partial B_x}{\partial t} = -(\nabla \times \vec{E})_x = \frac{\partial E_y}{\partial z} - \frac{\partial E_z}{\partial y} \quad (1.26)$$

$$= \frac{E_y^{i,j,k+1} - E_y^{i,j,k-1}}{2\Delta z} - \frac{E_z^{i,j+1,k} - E_z^{i,j-1,k}}{2\Delta y} + \mathcal{O}(\Delta y^2, \Delta z^2) \quad (1.27)$$

where $\vec{E} = -\vec{v} \times \vec{B}$, Δy and Δz are the spatial resolution in the y and z directions, i, j, k denote the grid indices and $\mathcal{O}(a)$ denotes terms of order a or higher. In this case, the error in spatially discretising the system, $\mathcal{O}(\Delta y^2, \Delta z^2)$, essentially creates a numerical electric field which breaks the frozen-in approximation.

1.2.4 Magnetohydrodynamic Waves and Shocks

Similar to hydrodynamic fluids, MHD allows for the existence of waves. Waves allow changes in the local fluid properties to propagate through the fluid smoothly. Hence waves are an

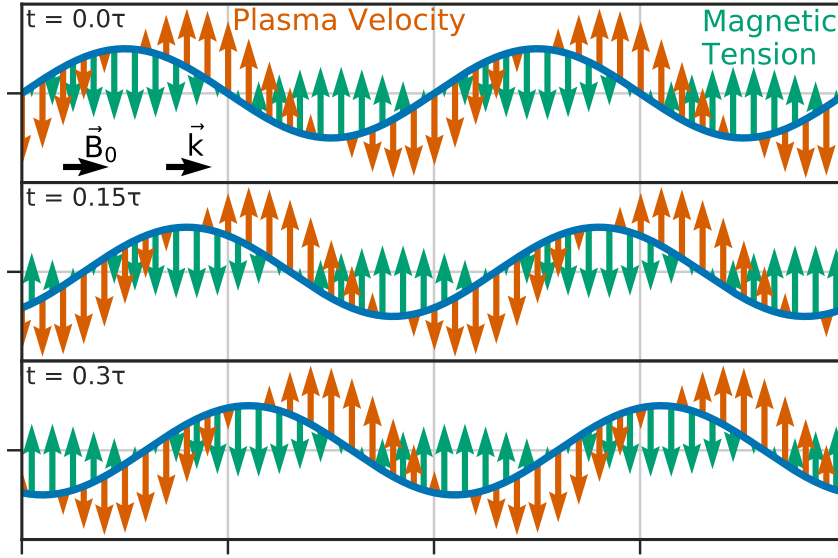


Figure 1.2: A diagram showing an Alfvén wave propagating to the right, with time increasing with every panel (τ is the period of the wave). The magnetic tension force (green) acts to reduce magnetic field curvature, hence points down the peaks, and up on the troughs. This force accelerates the plasma, which is stationary at the peaks and troughs and fastest at equilibrium. The inertia of the plasma carries the field line past equilibrium, where it is decelerated by the tension force, then back towards equilibrium. This process repeats, creating an Alfvén wave.

important mechanism for transmitting information around the system. In the hydrodynamic case these are caused by variations in pressure. In MHD the embedded magnetic field provides two wave modes: the Alfvén wave and the magnetosonic wave.

The Alfvén wave is caused by the magnetic tension force, which acts to reduce the curvature in the magnetic field, accelerating the plasma. The plasma inertia causes the field to curve in the opposite direction, providing a magnetic tension in the opposite direction. Hence, the magnetic force acts as restoring force, causing the magnetic field and plasma velocity to oscillate. A diagram of this is shown in figure 1.2. This produces a non-dispersive, non-compressional, transverse wave propagating at the Alfvén speed, $v_A = B/\sqrt{\mu_0\rho}$, as shown by its dispersion relation,

$$\omega_A^2 = k_{\parallel}^2 v_A^2. \quad (1.28)$$

As the name suggests, the magnetosonic wave is the MHD extension of a sound wave in a hydrodynamic fluid. It is caused by both pressure and magnetic field variations. Equation (1.29)

shows the dispersion relation of a magnetosonic wave, which is also dependent on the sound speed, $c_s = \sqrt{\gamma P/\rho}$. As with the Alfvén wave, the wave speed has a preferred direction. The magnetosonic phase speed is fastest perpendicular to the magnetic field, and slowest parallel to it.

$$\omega_{ms}^2 = \frac{k^2}{2} \left(v_A^2 + c_s^2 \pm \sqrt{(v_A^2 + c_s^2)^2 - 4v_A^2 c_s^2 \hat{k}_\parallel^2} \right) \quad (1.29)$$

The dispersion relation has two distinct roots, relating to the fast (positive root) and slow (negative root) magnetosonic wave. This distinction is due to magnetosonic waves being driven by two main variations, the pressure and the magnetic field. When these are correlated, their effect adds together causing a fast wave. The slow wave occurs when they are anti-correlated.

All three of the MHD wave speeds are dependent on the sound speed, Alfvén speed, and the direction of propagation with respect to the magnetic field, $\hat{k}_\parallel$. This dependence is shown in figure 1.3, with the wave phase speed plotted as a function of the angle between the magnetic field and the wave vector ($\vec{k}$), θ_{Bk} . The fast magnetosonic wave (green curve) propagates in all directions, but is fastest perpendicular to the magnetic field direction, reaching the magnetosonic speed $v = \sqrt{c_s^2 + v_A^2}$. With decreasing θ_{Bk} , the wave phase speed decreases: parallel to the field ($\theta_{Bk} = 0$), it travels at either the sound or the Alfvén speed, depending on which is larger. The Alfvén (blue curve) and slow magnetosonic wave (orange) show the opposite behaviour. They travel fastest parallel to the field, peaking at v_A for the Alfvén wave. Their speed decreases with increasing θ_{Bk} , and are slowest perpendicular to the field. In fact, both wave modes do not propagate at all in the perpendicular direction, as the phase speed reaches zero.

In hydrodynamic fluids, sound waves travel at the same speed in every direction. This is because the force due to the pressure, $-\nabla P$, is isotropic, and can accelerate the gas in any direction. The asymmetry in direction in MHD comes from the fact that the magnetic force, $\vec{J} \times \vec{B}$, always acts perpendicular to the magnetic field, and thus strongly affects variations in this direction.

In the regime where the plasma flow is larger than the wave speed, the flow is said to be supersonic¹. For a plasma, this wave speed is typically the magnetosonic speed, $c_{ms} = \sqrt{c_s^2 + v_A^2}$. In this region, waves are not fast enough to propagate the change through the fluid smoothly; this causes shocks to form. Shocks are sharp transitions, localised to a small region, where the

¹While supersonic technically refers to flows faster than the sound speed, here and throughout the thesis, it is used more generally to mean flows faster than the magnetosonic speed.

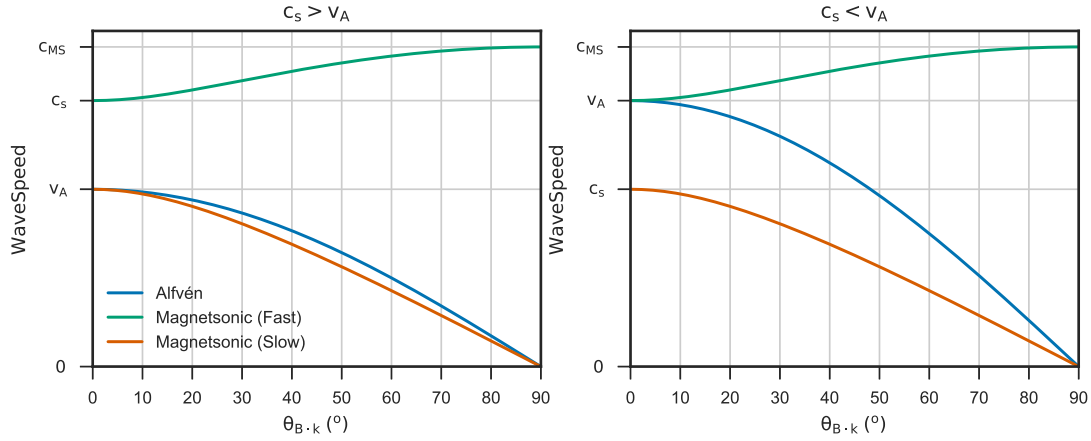


Figure 1.3: The Alfvén (blue), fast (green) and slow (orange) magnetosonic wave phase speeds as a function of the angle between the wave vector $\vec{k}$ and the mean magnetic field.

fluid is heated irreversibly, rapidly compressed and decelerated to sub-sonic speeds. As there are two compressional wave modes in MHD, there are also two shock modes: fast and slow. Shocks require a large amount of dissipation in the vicinity of the shock, which provides the irreversible heating. In hydrodynamic fluids, this dissipation is provided by collisions within the fluid. However, typical space plasmas are collisionless, which makes the formation and properties of these shocks much more complex. This dissipation is a phenomenon which occurs over small length scales (Burgess and M. Scholer, 2013): a regime where MHD is not applicable.

Fortunately, an understanding of the full small-scale physics of the shock is not needed to gain insight into such shocks, since the conservation of mass, energy and momentum always apply. Using the MHD equations, a set of ‘jump’ conditions known as the Rankine-Hugoniot conditions, can be obtained, which shed light on the large-scale effect of the shock. These are given in equations (1.30) to (1.35), they relate the plasma parameters upstream (supersonic plasma) to their parameters downstream (shocked plasma) by invoking the conservation equations. This assumes the shock width is much smaller than the region where the Rankine-Hugoniot

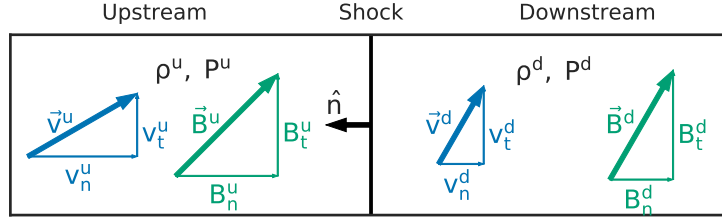


Figure 1.4: The shock geometry used in the Rankine-Hugoniot relations. The upstream values are on the left, denoted by the superscript u . Downstream is to the right, denoted by the superscript d . Vector quantities are set in the shock coordinate system, with normal the normal and tangential components, v_n and v_t .

conditions are applied, and that the shock is planar and stationary.

$$[\rho v_n] = 0 \quad (1.30)$$

$$\rho v_n [v_n] + \left[P + \frac{B^2}{2\mu_0} \right] = 0 \quad (1.31)$$

$$\rho v_n [\vec{v}_t] - \frac{B_n}{\mu_0} [\vec{B}_t] = 0 \quad (1.32)$$

$$\left[v_n \left(\frac{\rho v^2}{2} + \frac{\gamma}{\gamma - 1} P + \frac{B_t^2}{\mu_0} \right) - \frac{B_n}{\mu_0} \vec{v} \cdot \vec{B} \right] = 0 \quad (1.33)$$

$$[B_n] = 0 \quad (1.34)$$

$$B_n [\vec{v}_t] - [v_n \vec{B}_t] = 0 \quad (1.35)$$

The Rankine-Hugoniot equations use the coordinate system aligned to the shock normal, $\hat{n}$, where the subscript n corresponds to the vector component parallel to the shock normal (i.e. $B_n = \vec{B} \cdot \hat{n}$, the normal component). The subscript t corresponds to the tangential component, which is perpendicular to the shock normal. The square brackets, $[\]$, are shorthand for the difference between the upstream and downstream values (i.e. $[f] = f^u - f^d$). The geometry of the shock and the upstream and downstream parameters are shown in figure 1.4.

The first four Rankine-Hugoniot relations describe the conservation of plasma properties across the shock. The first, equation (1.30), states the amount of mass into the shock must equal the amount out of the shock: the conservation of mass. Equations (1.31) and (1.32) describe the conservation momentum in both the normal and tangential directions. Finally, equation (1.33) describes the conservation of energy, where a simple adiabatic energy equation is used: $P\rho^{-\gamma} = \text{const}$. These four are equivalent to the hydrodynamic Rankine-Hugoniot

equations, with the addition of the magnetic force and energy. The addition of the magnetic field imposes extra conditions. Equation (1.34) states the magnetic field component normal to the shock remains unchanged by the shock, and is due to the $\nabla \cdot \vec{B} = 0$ constraint. Equation (1.35) imposes Faraday's law (equation (1.19)), using only the motional electric field $\vec{E} = -\vec{v} \times \vec{B}$.

A useful property of MHD shocks is magnetic coplanarity. This states that the upstream and downstream magnetic field directions and the shock normal lie in the same plane. It is used in spacecraft analysis of shocks, where the shock normal is determined from the upstream and downstream magnetic field directions (Schwartz, 1998).

The Rankine-Hugoniot conditions are difficult to solve generally: they are highly non-linear with five independent variables. However, by considering specific orientations of the magnetic field direction with respect to the shock normal, some useful insights can be gained. When the upstream magnetic field is parallel to the shock normal ($B_t = 0$), it is said to be a parallel shock. The magnetic field strength remains unchanged through the shock, due to the $\nabla \cdot \vec{B} = 0$ constraint in equation (1.34), though the plasma still undergoes heating, compression and deceleration. Hence in a parallel shock, the MHD fluid behaves most like a hydrodynamic fluid, with the magnetic field playing a minimal role. The magnetic field plays more of a role in perpendicular shocks, when the magnetic field and the shock normal are perpendicular to one another ($B_n = 0$). In this case, both the magnetic field strength increases and the fluid is compressed through the shock. These are always fast mode shocks, as slow magnetosonic waves do not propagate perpendicular to the field.

In reality, perfectly parallel and perpendicular shocks are unlikely to occur. Variation in the interplanetary magnetic field and the curved nature of the structures such as a planetary bow shock means the shock normal can be at an arbitrary angle to both the magnetic field and the plasma flow. These are more generally called oblique shocks, which are further split into two categories: fast and slow mode shocks. Similar to slow magnetosonic waves, where the thermal pressure and the magnetic field strength are anti-correlated, slow mode shocks decrease the magnetic field strength after the shock, deflecting the magnetic field direction toward the shock normal. Fast mode shocks increase the magnetic field strength, and hence the shock deflects the field direction away from the normal. This is illustrated in figure 1.5, which shows a diagram of fast, slow shocks, and the intermediate 'shock'. Fast and slow mode shocks are strongly related to their wave counterparts, the fast and slow magnetosonic wave. One may expect a similar such shock for the third wave, the Alfvén wave. This is the intermediate mode, which

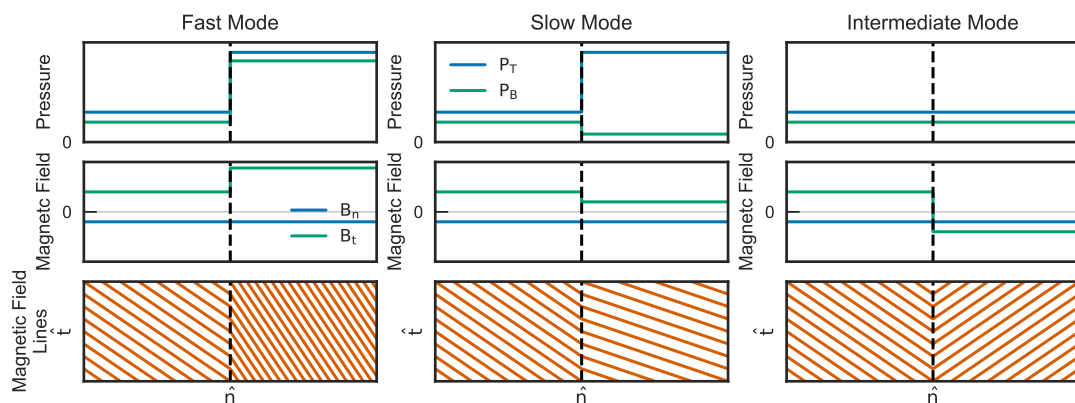


Figure 1.5: The three shock modes in MHD: fast (left), slow (middle), intermediate (right). The shock surface is the black dashed line through the middle, with upstream to the left, and downstream to the right. The velocity and the shock normal are parallel, with an arbitrary magnetic field. The top row shows the change in thermal and magnetic pressure across the shock. The middle row shows how the magnetic field rotates, through the shock. The bottom row shows magnetic field lines and how they kink through the shock.

causes 180° rotation in the tangential magnetic field. However, similar to the Alfvén wave, the intermediate ‘shock’ does not compress nor heat the plasma at the discontinuity, and hence is not technically a shock. It fits better as a specific case in the group of rotational discontinuities, where the magnetic field rotates across the boundary. In general, MHD shocks and rotational discontinuities are subsets of general MHD discontinuities.

When taking into account the effect of individual particles, the geometry of the magnetic field upstream and downstream of the shock becomes especially important. Large-scale particle processes observed at the shock, such as reflection of particles off the shock, are dependent on this geometry (J. P. Eastwood et al., 2005). Furthermore, geometry also has a role in determining the dissipation that forms the shock (Burgess and M. Scholer, 2013). This will be further explored in the section on the bow shock.

1.2.5 Extended Magnetohydrodynamics

A number of assumptions and simplifications have been made in order to arrive at the MHD equations (1.16) to (1.20). MHD can be extended to include more complex physics, taking into

account some of the anisotropies in proton and electron motions or asymmetries with respect to the magnetic field.

For example, rather than using a scalar pressure, an anisotropic pressure tensor can more accurately capture the dynamics and temperature of the plasma. This usually consists of splitting the tensor pressure into two components, parallel and perpendicular to the magnetic field.

Furthermore, the Ohm's law given in equation (1.20) has already been simplified, neglecting terms considered to be small compared to the motional electric field and resistive electric field. The generalized Ohm's law includes the Hall effect, $\vec{J} \times \vec{B}$, the electron pressure contribution, $\nabla \cdot \mathbf{P}_e$, and the electron inertia term, $\partial_t \vec{J}$,

$$\vec{E} + \vec{v} \times \vec{B} = \eta \vec{J} + \frac{1}{ne} \vec{J} \times \vec{B} - \frac{1}{ne} \nabla \cdot \mathbf{P}_e + \frac{m_e}{ne^2} \frac{\partial \vec{J}}{\partial t} \quad (1.36)$$

Each of these new terms in the generalised Ohm's law are due to the fact that electrons behave differently than ions: they are several orders of magnitude lighter, and hence respond to changes in the fields much faster than the ions. None of these extra terms are related to collisions in the plasma, and hence are valid in most space plasmas. Similar to the effect of resistivity, these extra terms break the frozen-in condition for the plasma. Hence they are expected to play a role in regions of large current, and significant electron pressure gradients, such as in reconnection.

However, these are not trivial to implement, adding extra complexity. For example, the extended Ohm's law requires determination of the electron pressure tensor, which is not readily available from the MHD equations. It also allows for the existence of new waves, like Whistler waves, which may add extra computational difficulties.

1.3 Solar wind interaction with magnetised planets

This section, contains a general overview of the magnetospheric system and its interaction with the solar wind. The solar wind is a significant driver in the magnetospheric system, controlling the size and shape, and hence dynamics of the outer boundaries of the magnetosphere. Understanding the generation and evolution of the solar wind explains many of the features observed in the solar wind throughout the heliosphere. Once the solar wind arrives at a magnetised planet, it interacts with its magnetic field to form complex processes which have

a global effect throughout the magnetosphere. This material is covered in a many textbooks, for example Baumjohann and Treumann (2012a) and Margaret G Kivelson and Christopher T Russell (1995).

1.3.1 Generation of the Solar Wind and Interplanetary Magnetic Field

The solar wind is a stream of hot and rarefied plasma that emanates from the Sun in all directions, travelling outward through interplanetary space. It carries the Sun's magnetic field with it, creating the interplanetary magnetic field (IMF).

Generation of the solar wind is a complex and not fully understood phenomenon. This is largely due to the complex and unpredictable motions and dynamics of the corona, the Sun's outermost plasma region, where the solar wind is generated (Owens and Forsyth, 2013). The solar corona is a hot, highly conductive and low beta plasma, attached to the solar magnetic field emanating from the solar surface, the photosphere. The large thermal pressure gradient between the corona accelerates the coronal plasma anti-Sunward, generating the ambient solar wind (Parker, 1965).

Simplified models of the solar wind can provide physical insight into some of the properties and structure observed. Parker (1965)'s solar wind model adequately describes the evolution by treating the solar wind as an unmagnetised gas which undergoes isothermal expansion. His spherically symmetric model, derived from the one-dimensional hydrodynamic fluid equations, gives a number of possible solutions. Only one corresponds to observation: the coronal plasma is initially sub-sonic, then is accelerated by the large thermal pressure gradient to supersonic speeds, forming the observed supersonic solar wind. The transition from sub- to supersonic occurs at a critical point, approximately 5.8 solar radii (0.02 AU).

Since the solar wind is highly conductive, the coronal magnetic field is frozen-in to the plasma and is carried outward with the bulk plasma flow. Whilst the magnetic field lines are carried radially anti-Sunward by the solar wind, they are also tied to locations on the rotating solar surface, creating a spiral IMF structure called the Parker spiral (Parker, 1965). This is illustrated in figure 1.7 which shows a view of the ecliptic plane of the heliosphere. The spiral angle (the angle the magnetic field makes to the flow direction) is dependent on the solar wind speed, the solar rotation rate and the distance from the Sun. The Parker spiral explains the general structure of the IMF observed throughout the heliosphere. For typical solar wind parameters

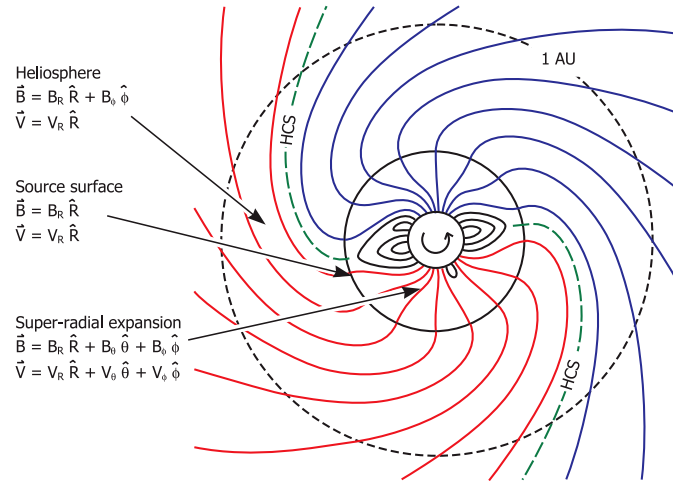


Figure 1.6: The structure of the IMF and coronal magnetic field in the ecliptic plane. The lines represent the magnetic field, with red and blue colours denoting the polarity: toward or away from the Sun. The green dashed line shows the Heliospheric Current Sheet (HCS), which separates the two polarities. Close to the Sun, within the ‘source surface’, is the coronal magnetic field, which consists of open magnetic fields from coronal holes, and closed magnetic fields from helmet streamers. Further away, the field gains the Parker spiral profile. From Owens and Forsyth (2013).

at Earth (1 AU), the spiral angle is approximately 45° , which is in agreement with observation. As the solar wind moves through the heliosphere, the spiral becomes more wound up and the angle tends to 90° , as was observed by the Voyager spacecraft at Neptune (Slavin and Holzer, 1981).

During a solar minimum, when solar activity is at its lowest, the Sun’s magnetic field can be considered dipolar close to the Sun’s surface, with two regions of magnetic polarity that are separated by the Heliospheric Current Sheet (HCS). The dipole is offset by a few degrees with respect to the Sun’s rotation axis (Owens and Forsyth, 2013), and so, as the Sun undergoes its ~ 25 day rotation, the radial IMF direction has a sector structure, reversing direction every half solar day in the ecliptic plane. This is also illustrated in figure 1.6 by the colour of the field lines (e.g. blue for inward, red for outward).

In reality, the solar corona is not uniform nor symmetric, as illustrated by the near Sun magnetic field in figure 1.6. This is due to the Sun’s magnetic field, which extends out of the Sun’s surface into the corona (Owens and Forsyth, 2013). The high conductivity in the corona keeps plasma attached to these field lines, which are either closed (black), called coronal loops, or

Table 1.2: Typical solar wind conditions at 1 AU (Hundhausen, 1995).

Density	Proton	6.6 cm^{-3}
	Electron	7.1 cm^{-3}
	He ²⁺	0.25 cm^{-3}
	Total	$8.1 m_p \text{ cm}^{-3}$
Average flow speed		450 km s^{-1}
Temperature	Proton	$1.2 \times 10^5 \text{ K} = 10.4 \text{ eV}$
	Electron	$1.4 \times 10^5 \text{ K} = 12.1 \text{ eV}$
Magnetic field	Strength	7 nT
	Spiral angle	45°

open (coloured), called coronal holes. The open field in coronal holes allow plasma to freely stream out into the heliosphere, producing the ‘fast’ solar wind, which travels at $\sim 800 \text{ km/s}$. The ‘slow’ wind is thought to be produced from helmet streamers: the boundaries between coronal loops and holes. Generated by the interaction between plasma on the open-closed field line boundary, the slow solar wind does not exhibit the same expansion as on open field lines, hence is slower than the fast solar wind, travelling at $\sim 400 \text{ km/s}$. During a solar minimum, coronal holes are commonly found at high solar latitudes, and coronal loops at low latitudes. During a solar maximum, the solar surface is much more active, leading to more unpredictable positions of coronal loops and holes.

1.3.2 The Heliospheric Solar Wind

Parker’s solar wind model can explain the presence of the solar wind throughout the heliosphere, and the general structure of the interplanetary magnetic field.

At the average radius of Earth’s orbit ($1 \text{ AU} = 1.50 \times 10^8 \text{ km}$), the typical slow solar wind parameters are shown in table 1.2. As the solar wind moves through the heliosphere, the evolution of the average parameters can be modelled by considering the conservation of mass, momentum and magnetic flux. This leads to a number of scaling laws, which are used to scale

solar wind parameters observed at 1 AU (Slavin and Holzer, 1981),

$$v = v_0, \quad (1.37)$$

$$n = n_0 \left(\frac{r_0}{r} \right)^2, \quad (1.38)$$

$$B = B_0 \frac{r_0}{r} \sqrt{2 \left(\frac{r_0}{r} \right)^2 + 1}, \quad (1.39)$$

$$T_p = T_{p0} \left(\frac{r}{r_0} \right)^{-2/3}, \quad (1.40)$$

$$T_e = T_{e0} \left(\frac{r}{r_0} \right)^{-1/3}, \quad (1.41)$$

where the subscript 0 denotes the reference values (e.g. parameters shown in table 1.2 at $r_0 = 1$ AU), and r is the radial distance from the Sun. Most of the quantities above give some radial dependence on the solar wind parameters, decreasing in size with distance from the Sun, with the exception of the velocity, which stays constant.

In reality, the solar wind is much more variable than Parker's model suggests (Owens and Forsyth, 2013). Observations of the solar wind by spacecraft reveal fluctuations in all plasma quantities over observed length and timescales, due to a combination of waves, turbulence, shocks and rotational or tangential discontinuities (Timothy S. Horbury et al., 2001; Bruno and Carbone, 2013).

The solar wind also exhibits some large-scale transient structures generated by the Sun such as corotating interaction regions (CIR) and coronal mass ejections (CME). These consist of increases in solar wind density and magnetic field strength, and often create enhanced and prolonged geomagnetic effects (Owens and Forsyth, 2013).

A corotating interaction region is an example of a transient structure in the solar wind (E. J. Smith and Wolfe, 1976; Pizzo, 1991; Gosling and Pizzo, 1999; Crooker et al., 1999). If the slow and fast streams are ejected from the same latitude, the fast solar wind catches up with the slow region in front of it. The frozen-in nature of the plasma means the two regions cannot pass through each other, creating a stream interaction region. This region mediates the transition between the slow and fast solar wind through two shocks, propagating away from the stream interface in the solar wind rest frame, as seen in figure 1.7. Similarly, the lagging of a slow stream behind a fast stream creates a large rarefaction region.

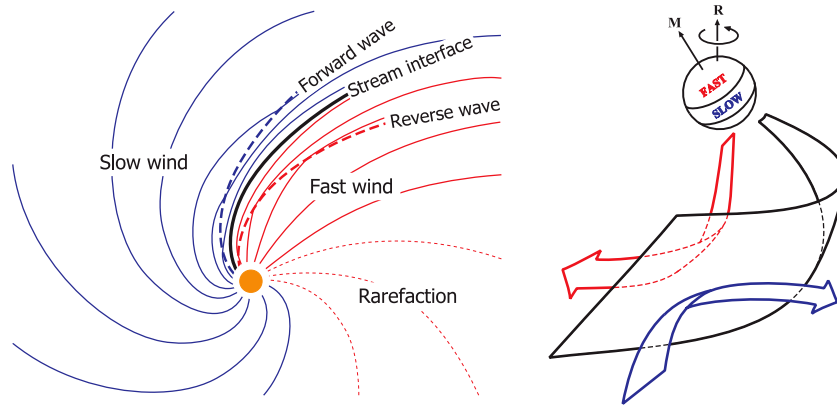


Figure 1.7: A Corotating Interaction Region as seen in the ecliptic plane on the left, and its 3D structure on the right. The lines denote magnetic field lines, coloured by polarity. Fast wind, usually generated from high solar latitudes, catches up with the slow, generated from low latitudes, creating the interaction region. From Owens and Forsyth (2013).

Coronal mass ejections are large-scale eruptions of plasma from the solar surface (Chen, 2011). Whilst they begin with their size on the scale of coronal loops (10^4 m), they can rapidly expand to fill a large proportion of the solar surface. As these CMEs move through the heliosphere, they are accelerated to the ambient solar wind speed (Gopalswamy et al., 2000; Peter J. Cargill, 2004). They exhibit a complex magnetic field topology, with strong rotational flux rope structures and enhanced plasma density (Chen, 2011); they are particularly effective in opening up magnetic flux in Earth's magnetosphere, causing strong substorms and space weather effects (Gosling, 1993; Tuija I. Pulkkinen, 2007).

1.3.3 The Magnetosphere

Many of the planets in the Solar System generate their own magnetic field through their internal dynamo. This field extends into interplanetary space where it interacts with the solar wind and interplanetary magnetic field (S. Chapman and Ferraro, 1931). Due to frozen-in flow, the solar wind cannot penetrate the planetary field under normal circumstances, which separates the magnetic environment around the planet from the heliospheric magnetic field, forming a magnetosphere. An illustration of this is shown in figure 1.8.

The magnetosphere can be roughly sub-divided into a number of interconnected regions: the dayside, the inner magnetosphere, and the magnetotail. Each of these regions contain plasma

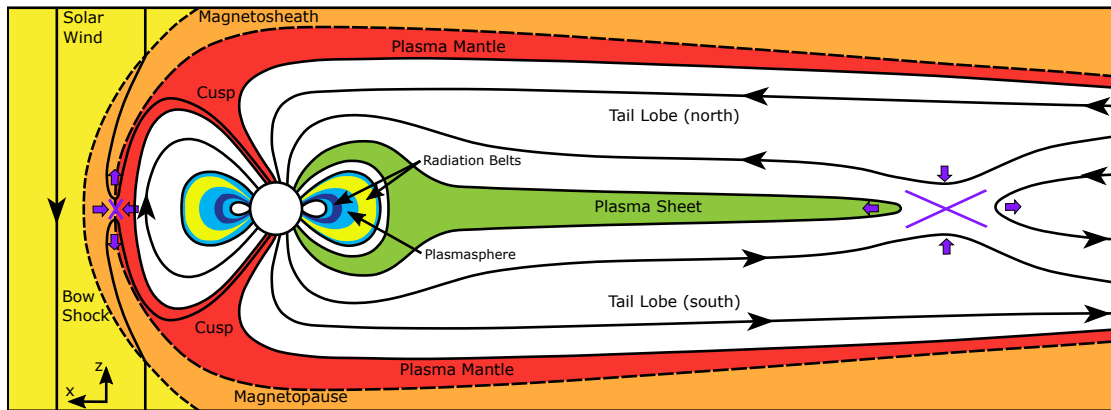


Figure 1.8: The magnetosphere in the noon-midnight plane. To the left, is the Sun and the oncoming solar wind, shown in yellow. The solar wind encounters the magnetopause, which causes the bow shock to form Sunward. Near the Earth is the inner magnetosphere, which consists of the plasmasphere and radiation belts on closed field lines. The magnetosphere is elongated on the nightside, creating the long magnetotail. On closed field lines, the plasma sheet is shown in green. The turbulent magnetosheath region between the shock and the magnetopause is shown as the orange region. Magnetic field lines are shown in black. In this case, the IMF is southward, causing dayside and nightside reconnection. The x-lines are shown in purple, with plasma in- and outflow. From J. P. Eastwood et al. (2015).

with different parameter regimes, and local processes and dynamics. They each play a part in the global dynamics of the magnetosphere, whose effect we observe as space weather at Earth.

The direct interaction between the solar wind and magnetosphere occurs on the dayside. Here, the solar wind encounters the planet's magnetic field, which acts as an obstacle to its flow due to the frozen-in condition. The magnetopause is the boundary that separates the solar wind and magnetospheric plasma (plasma tied to the planetary field). Its position is determined by the pressure balance between the solar wind dynamic pressure and the planetary magnetic field pressure (S. Chapman and Ferraro, 1931; Beard, 1960). This forms the characteristic shape of the magnetosphere, shown in figure 1.8, with the planetary magnetic field dipole compressed in the Sunward direction (to the left) and elongated in the anti-Sunward direction (to the right) (Tsyganenko, 1995). Sunward of the magnetopause is the bow shock, which mainly decelerates the supersonic solar wind flow such that it can be deflected around the magnetosphere (Spreiter, Summers, and Alksne, 1966; Donald H Fairfield, 1971). Between the bow shock and the magnetopause is a region of turbulent shocked plasma known as the magnetosheath (Ness, Scarce, and Seek, 1964; D. H. Fairfield, 1976; Lucek et al., 2005; Burgess

and M. Scholer, 2013; Krasnoselskikh et al., 2013), shown as the orange region in figure 1.8.

The region near the planet is the inner magnetosphere, which contains the plasmasphere, ionosphere and radiation belts. The plasmasphere is a cold ($\sim 1\text{eV}$) and dense ($\sim 10^3/\text{cm}^3$) plasma (Chappell, 1972; Park, Carpenter, and Wiggin, 1978), extending from $3 R_E$ to $5 R_E$ in the case of the Earth (shown by the blue and yellow colours in figure 1.8). Due the high magnetic field strength in this region, its dynamics are dominated by the approximately dipolar planetary magnetic field, creating field-aligned structures which tie plasma populations in the plasmasphere to the ionosphere. The ionosphere is the weakly ionized outermost layer of the planet's atmosphere, and is considered to be the inner boundary to this region. It is an important source of plasmaspheric plasma, populating the plasmasphere via the polar wind. It is formed by ionizing the top layer of the atmosphere through photo-ionization from solar radiation on the dayside and electron precipitation from the magnetosphere, which creates a ring of enhanced conductivity around the planet's magnetic poles called the auroral oval. The ionosphere also plays an important role in controlling global magnetospheric dynamics, linking in to the global magnetospheric current systems. Aside from the cold plasmaspheric plasma, the inner magnetosphere has a significant population of energetic particles (Daly, 1994). They are trapped in 'radiation belts', shown as the dark blue and yellow regions in the inner magnetosphere in figure 1.8. These have a strong effect on space weather, as these regions can coincide with spacecraft orbits, potentially damaging spacecraft by impact ionisation (Lucci et al., 2005; Koons and Fennell, 2006).

Further anti-Sunward, the magnetotail forms in the planet's wake. This region consists of the plasma sheet (green region in figure 1.8), in between two magnetic lobes. The plasma sheet contains relatively low temperature and low mass density plasma, and magnetic field lines which have been stretched out from their usual dipolar shape. Located approximately along the ecliptic plane, the magnetic field B_x component switches direction from pointing Sunward to anti-Sunward. This is the location of the tail current sheet, which connects to the magnetopause current sheet. It is accompanied by a thin region of plasma known as the plasma sheet. The dynamics of the magnetotail are of particular interest to space weather. During conditions where dayside reconnection occurs, magnetic flux can become stored in the magnetotail (Keiling, 2009; Sharma et al., 2008). Upon reaching a threshold not fully understood, the magnetic energy stored in the tail is transferred to the plasma usually through tail reconnection, creating a substorm. This accelerates particles into the polar ionosphere, creating adverse space weather.

The solar wind driving is not necessarily the largest driver in a magnetospheric system. Whilst it is the case for Earth, planets like Jupiter and Saturn are driven strongly by plasma ejected from their moons (Krupp et al., 2004; Mitchell et al., 2009). Some of these moons, like Ganymede, are magnetised themselves, creating a complicated magnetosphere within a magnetosphere system (M. G. Kivelson et al., 1997).

1.3.4 The Magnetopause

The magnetopause is the boundary between the solar wind and the magnetospheric plasma. It separates the hot, turbulent shocked solar wind from the relatively cold and rarefied magnetospheric plasma. The boundary is in pressure balance with the solar wind acting to compress the planet's intrinsic magnetic field.

The magnetopause can be identified by a number of interesting features. The magnetospheric plasma density is relatively low compared to the magnetosheath plasma, meaning there is a large jump in density across the magnetopause (Schwartz, 1998). There is also typically a change in magnetic field magnitude and orientation, creating a current sheet known as the Chapman-Ferraro current layer, which is one of the main current systems in the magnetosphere.

Its sub-solar location (or stand-off distance) can be approximated by considering the pressure balance between the solar wind's dynamic pressure and the magnetic pressure of the planet's intrinsic field,

$$\rho_{\text{sw}} v_{\text{sw}}^2 = 2 \frac{B_{\text{intr.}}^2}{2\mu_0} \quad (1.42)$$

where $B_{\text{intr.}}$ is the planet's intrinsic magnetic field at the magnetopause, with solar wind density (ρ_{sw}) and speed (u_{sw}). By assuming a dipole field, $B_{\text{intr.}} = B_{\text{P}}(R_{\text{P}}/r)^3$, the standoff distance of the magnetopause can be estimated with the following relation,

$$r_{\text{MP}} = R_{\text{P}} \left(\frac{2B_{\text{P}}^2}{\mu_0 \rho_{\text{sw}} u_{\text{sw}}^2} \right)^{1/6}, \quad (1.43)$$

where R_{P} is the planet radius, B_{P} the equatorial intrinsic magnetic field strength. The factor of two on the right hand side arises from the fact that there is a current sheet along the magnetopause, which acts to bolster the magnetic field planet-ward of the magnetopause, and negate the field Sunward. Using the standard solar wind parameters at Earth in table 1.2, and

an equatorial field strength of 30 nT, this estimates Earth's magnetopause stand-off position to be approximately $10 R_E$ ($R_E = 6.37 \times 10^3$ km, Earth radius), which agrees with observation.

Whilst this first order estimate gives an understanding on what supports the magnetopause, more complex models of the magnetopause have been developed using spacecraft observations. Models such as Shue et al. (1998) and Lin et al. (2010) have been developed by locating spacecraft magnetopause crossings, and fitting a set magnetopause shape under different solar wind conditions. Such magnetopause models usually assume some functional form for the magnetopause shape. Whilst the Shue et al. (1998) only fits to a simple conical, axisymmetric shape, the Lin et al. (2010) model includes cusps and asymmetries in the ecliptic and noon-midnight planes.

Under some conditions, solar wind plasma can enter through the magnetopause surface. Observations have shown that the effectiveness of plasma entry into the magnetosphere is strongly dependent on the angle between the IMF and the magnetospheric field at the magnetopause (De Keyser et al., 2005). In locations where they are anti-parallel, magnetic reconnection occurs which breaks the frozen-in condition locally, allowing plasma entry.

Dungey (1961) proposed a mechanism by which magnetic flux can be opened on the dayside through reconnection, becoming an open field line (field line 2 in figure 1.9). The solar wind carries this open field line with it, dragging its footpoint through the ionosphere around the magnetic poles of the Earth (field lines 3 to 6 in figure 1.9) towards the nightside of the planet, where it reconnects again in the tail (field line 7 in figure 1.9), becoming closed again (field line 8 in figure 1.9). Whilst open, the motional electric field of the solar wind can penetrate into the magnetosphere, causing Sunward convection of magnetospheric plasma by $\vec{E} \times \vec{B}$ drift. The plasma drifting Sunward gets channelled into the polar ionosphere, causing large geomagnetically induced currents. This cycle repeats for as long as dayside reconnection occurs. The closing of magnetic fields through reconnection on the tail side also acts to divert the cross tail current sheet along the field into the poles (Kepko et al., 2015). The ionosphere closes these field-aligned currents in areas of higher conductivity (causing the auroral oval), completing the circuit by diverting the current back into the tail current (Kepko et al., 2015). On a macroscopic scale, energy is transferred onto the ionosphere by Ohmic heating. This is a simplified view of the complex current systems that exist in the inner magnetosphere and magnetotail (Kepko et al., 2015).

The Dungey cycle is a two dimensional model, and in particular it implies the reconnection on the dayside is steady. In reality, the magnetopause is a three-dimensional, with the potential

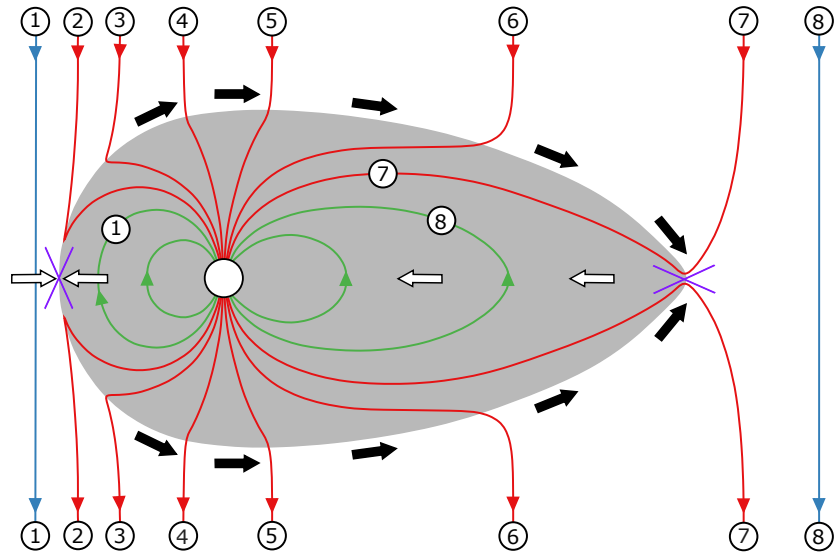


Figure 1.9: The Dungey cycle, showing the movement of field lines (grey arrows) and plasma convection direction (white arrows). Adapted from Baumjohann and Treumann (2012b)

for reconnection to occur over a region. Variations in the reconnection rate, or multiple x-line reconnection can form twisted magnetic field structures, known as flux ropes or flux transfer events (C. Russell and Elphic, 1978; Southwood, C. Farrugia, and Saunders, 1988; Manfred Scholer, 1988; Lee and Fu, 1985). These three-dimensional structures vastly complicate the Dungey cycle.

The Dungey cycle is triggered by reconnection occurring on the dayside. In the case of the Earth, its intrinsic magnetic field does not significantly precess. Its magnetic moment stays predominantly perpendicular to the solar wind flow, in the $-z_{\text{GSE}}$ direction. Hence the IMF B_z component primarily controls the position of reconnection on the magnetopause: the Earth's magnetic field on the dayside sub-solar region is positive, meaning the IMF must have a 'southward' orientation, $\text{IMF } B_z < 0$, for reconnection to occur in this region. Under 'northward' IMF, $B_z > 0$, the IMF drapes around the magnetopause, reconnecting at cusps: this is where the IMF is anti-parallel to Earth's field. In other magnetised planets, the importance of the Dungey cycle is not as clear. Jupiter and Saturn's magnetospheres have strong internal drivers, namely their active moons, which can dominate their magnetosphere's dynamics. Neptune and Uranus' both have magnetic moments significantly offset from their rotation axes, meaning the planetary field varies significantly over their respective days, further inhibiting reconnection and the Dungey cycle (Cowley, 2013; Masters, 2015).

Whilst dayside reconnection is thought to be most effective at allowing plasma to enter the magnetosphere, there are a number of other proposed models for plasma entry. It is thought that wave-particle interactions can cause plasma to diffuse across the magnetopause (Haerendel et al., 1978; T. D. Phan et al., 2005). In regions along the magnetopause flank, the fast magnetosheath plasma flows past the approximately stationary magnetospheric plasma. This causes the Kelvin Helmholtz instability to form, creating large, rolled-up vortices at the flank magnetopause (Hasegawa et al., 2004). The micro-scale eddies created in this interaction can allow plasma to diffuse across the magnetopause boundary.

1.3.5 The Bow Shock

The bow shock forms upstream of the magnetopause, due to the obstacle formed by the magnetosphere. With solar wind Mach numbers typically ~ 8 , it is usually categorised as a strong shock. Due to the blunt profile of the magnetopause, the bow shock is detached from the magnetopause. Its main role is to decelerate and deflect the supersonic solar wind around the magnetosphere, creating a turbulent magnetosheath region between the magnetopause and bow shock.

In Farris and C T Russell (1994)'s semi-empirical relation, the standoff position of the shock is given by,

$$D_{BS} = D_{OB} \left(1 + 1.1 \frac{(\gamma - 1) M_1^2 + 2}{(\gamma + 1) M_1^2} \right), \quad (1.44)$$

which relates the shock standoff position, D_{BS} , to the standoff position of the obstacle, D_{OB} , and the upstream Mach number, M_1 , and the gas constant γ . This shows the position of the bow shock is strongly dependent on the position of the obstacle, the magnetopause, which in turn is strongly dependent on the dynamic pressure of the solar wind. To a lesser extent, the shock is dependent on the solar wind Mach number (M_1), which determines width of the magnetosheath, and has been calculated from gas dynamic theory.

More recently, empirical models of the bow shock have been devised. These take a catalogue of spacecraft shock crossings, binned according to upstream parameters and fitted to a two-dimensional surface. There are a number of empirical models in the literature (e.g. Merka (2005), Iver H Cairns and J G Lyon (1995), and Jeřáb et al. (2005)), all derived using different datasets and reference surfaces. These models have shown empirically that the bow shock shape and

position is largely determined by the solar wind dynamic pressure, and to a lesser extent, the solar wind Alfvén Mach number (Peredo et al., 1995; Merka, 2005). In general, these models have been shown to agree with spacecraft observations of the shock, producing the correct positions, shock normals (e.g. T. S. Horbury et al. (2002) and K. Meziane, Alrefay, and Hamza (2014)) and velocities on the whole (Meziane et al., 2015).

Whilst the bow shock's main role is to deflect plasma around the magnetosphere, the collisionless nature of the solar wind creates interesting kinetic effects only seen at collisionless shocks. Upstream of the bow shock is a region called the foreshock, which consists of particles travelling in the opposite direction of the solar wind (J. P. Eastwood et al., 2005). These reflected particles leads to numerous wave-particle interactions and particle acceleration (J. P. Eastwood et al., 2005). Whilst MHD can not capture these effects, it does provide the large-scale shape of the shock, which is important for determining the shock normal and the location of the foreshock.

1.4 Global Simulations of the Magnetosphere

Modelling of the magnetosphere is needed to understand the global processes and dynamics of the system as a whole. Such models can loosely be categorised into empirical and physics-based numerical models.

A considerable amount of work has gone into creating empirical models of magnetospheric structures and processes. These are based on a catalogue of spacecraft observations of a certain phenomenon, and correlated to certain inputs to the system, such as the solar wind parameters. For example, a large number of bow shock and magnetopause models have been proposed (e.g. Shue et al. (1997), Lin et al. (2010), Farris and C T Russell (1994), Jeřáb et al. (2005), and Merka (2005)), which take spacecraft crossings of the shock or magnetopause, and relate them statistically to the observed solar wind parameters. Other modelling work includes the Tsyganenko magnetic field model (Tsyganenko, 1995), which models the Earth's magnetic field by taking into account the relevant current systems and planetary dipole, as well as the solar wind parameters (Tsyganenko, 2014). This is then fitted to spacecraft observations, where one obtains the strengths of each of the current systems.

These models are useful in quantifying the structure of the magnetosphere. Since many magnetospheric processes span over large scales, they are often used to provide a global context to

local measurements provided by spacecraft. However, they usually only capture the average properties of the magnetosphere, and can be limited by either the assumed model or the bias in spacecraft position. This means they do not necessarily do well in predicting extreme events, since there are relatively few of these that have been observed. Nor do they capture the time-dependent changes induced by, for example, coronal mass ejections and interplanetary shocks. In the case of magnetopause models for example, many assume a simple, axisymmetric conical shape to fit the data, which is not necessarily the best representation of the magnetopause shape. Furthermore, the majority of magnetopause crossings are at low latitudes: there are only few observations of the flank.

One way of attempting a physical understanding is to apply a physics-based model. This consists of applying a set of physical equations, such as MHD, to either a local region, or the magnetosphere as a whole. In global magnetosphere simulations, the magnetosphere is simulated as a whole, from the solar wind, through to an inner boundary near the planet, then to far down the magnetotail. In the literature, there are three main approaches to simulating space plasmas: kinetic (using mainly the collisionless Vlasov equations), Particle-In-Cell (PIC) and MHD.

Kinetic and PIC codes attempt to capture the full dynamics of space plasmas by either solving the Vlasov equation (kinetic) or simulating the behaviour of a large number of particles and how they interact the electromagnetic fields (PIC). There is a fourth method, which combines a kinetic model with a fluid model to improve on performance, usually treating the electrons as a fluid and the ions as kinetic. The global magnetosphere simulation, Vlasiator (Alfthan et al., 2014; Pfau-Kempf et al., 2016), is an example of a hybrid-Vlasov code.

Even when simulating small local regions, such codes are not easy to implement and have their own computational problems regarding initial conditions and resolution. Local simulations are also strongly affected by their choice of boundary conditions. For simulating the magnetosphere as a whole however, the main problem is their computational cost: with current technology, the computation time can take much longer than real time, which is too slow for any meaningful space weather prediction. Hence the majority of global magnetosphere codes use MHD, which is relatively computationally inexpensive compared to the other approaches. This allows MHD codes to run faster than real time with relatively little computing power.

The issue with physics-based codes is that compromises inevitably arise in the assumptions that are made to simplify the equations. With almost all physics-based models, some assumptions

have been made to simplify the equations. For example, MHD treats the motions of electrons and ions as one single fluid, and neglects ‘fast’ phenomena. By not including all the physics needed to simulate a system, differences may arise between simulation and reality. Another issue is that numerical models inject numerical error into the simulation. This could be from discretising the system either spatially or temporally. The most prominent of such an issue is concerning $\nabla \cdot \vec{B}$, where, if care is not taken in developing the correct solver, numerical errors can create magnetic monopoles.

Hence, work needs to be done in verifying numerical models and critically assessing under what circumstances they accurately reflect reality. However, in a system so large and sparsely observed as the magnetosphere, this is not a trivial process. This can be done by comparing the simulation to spacecraft data or empirical models.

1.4.1 Global Magnetohydrodynamic Codes

Due to the recent availability of high-performance computing resources, and services such as the Community Coordinated Modeling Center (CCMC), there have been numerous papers on global MHD simulations of planetary magnetospheres. In the case of the Earth, global magnetospheric codes usually span approximately $30 R_E$ Sunward, and 100 to $200 R_E$ anti-Sunward, containing the solar wind, magnetosheath, inner magnetosphere and tail regions. The majority of these codes use MHD in the main global simulation, but often couple into local models to capture non-MHD physics. An example of this is the ionosphere, whose plasma is strongly coupled to the neutral atmosphere, or the ring current in the inner magnetosphere.

The CCMC is a service provided by a number of US institutions which makes space physics models and simulations available to run on request. It includes a number of models of a variety of physical systems, both local and global, ranging from the Sun to the Earth. Four global MHD (GMHD) codes are available on CCMC, and have been widely used throughout the community. They are the Lyon-Fedder-Mobarry (LFM) (J. G. Lyon, Fedder, and Mobarry, 2004), Grand Unified Magnetosphere Ionosphere Coupling Simulation (GUMICS-4) (P. Janhunen et al., 2012), Block Adaptive Tree Solar-wind Roe Upwind Scheme (BATS-R-US) (Powell et al., 1999) and Open General Geospace Circulation Model (OpenGCCM) (Joachim Raeder et al., 2008).

The amount of literature using these codes is extensive. They range from studying the magnetosphere as a whole (e.g. Palmroth et al. (2001), Facskó et al. (2016), and T. I. Pulkkinen, Palmroth,

and McPherron (2007)), to focussing on specific regions such as the magnetosheath and bow shock (e.g. Samsonov, Sibeck, and Imber (2007) and Oliveira and Joachim Raeder (2014)) or tail (e.g. Sergeev, Tsyganenko, and Angelopoulos (2008)). Their low computational cost can allow for long simulation durations, for example, Facskó et al. (2016) provides the longest global MHD simulation to date, simulating the magnetosphere continuously for a year.

Though most effort has been directed at reproducing the dynamics and behaviour of the Earth's magnetosphere, global MHD codes have also been used to study other planets, for example the magnetospheres of Mercury (e.g. Kabin (2000), Ip and Kopp (2002), and Jia et al. (2015)), Jupiter (e.g. Ogino, Walker, and Margaret G. Kivelson (1998), Moriguchi et al. (2008), and Chané, Saur, and Poedts (2013)) and Saturn (e.g. Fukazawa et al. (2007), Kidder, Winglee, and Harnett (2009), and Jia et al. (2012)).

1.4.2 Common Features of GMHD Codes

When designing a global magnetosphere MHD code, there are many aspects of the magnetosphere system to consider. Each of the aforementioned codes differ in their choice of simulation grid, types of solvers and numerical methods used, and the formulation of the MHD equations.

In the development of those codes, a number of computational difficulties were found which caused non-standard formulations of the MHD equations. A difficulty in simulating the magnetosphere is the high magnetic field strength and low plasma β region in the inner magnetosphere. In this region, the field strength is dominated by the planet's intrinsic magnetic field, which follows an $1/r^3$ relationship.

The large gradient in the field causes truncation errors which creates unphysical current. This effect is reduced by using high order spatial gradients (J. G. Lyon, Fedder, and Mobarry, 2004) and an adapted or adaptive grid to provide high resolution in regions of large gradients. Alternatively, the MHD equations can be recast to split the magnetic field into two components, an external magnetic field $\vec{B}_1$ and a dipole field $\vec{B}_0$ (Tanaka, 1994), where,

$$\begin{aligned}\vec{B} &= \vec{B}_0 + \vec{B}_1, \\ \nabla \times \vec{B}_0 &= 0\end{aligned}$$

The curl-free nature of the dipole field means the current can be solely determined from $\vec{B}_1$.

The large magnetic field strength poses a particular problem in a code conservative in total energy, ϵ_T . As the large magnetic energy contribution must be subtracted from the total energy in order to obtain the pressure, as given in equation (1.45), the pressure is not guaranteed to be positive. Use of low order solvers and the need to interpolate within the grid (as needed in a staggered grid) can exacerbate the problem.

$$P = (\gamma - 1) \left(\epsilon_T - \frac{\rho v^2}{2} - \frac{B^2}{2\mu_0} \right) \quad (1.45)$$

Solutions to this include solving the energy equation in a non-conservative form, either using the internal energy (ϵ_i) or the plasma energy ($\epsilon_P = \epsilon_i + \rho v^2/2$) (J. G. Lyon, Fedder, and Mobarry, 2004), overcoming the need to subtract the large magnetic energy. In codes which employ a split magnetic field, a reduced total energy only containing the external magnetic field component in the total energy can be used, $\epsilon'_T = \epsilon_i + \rho v^2/2 + B_1^2/2\mu_0$.

In regions of high magnetic field strength and low mass density, the Alfvén speed becomes very large. As all GMHD codes use explicit solvers, their timestep is limited by the Courant-Friedrichs-Lewy (CFL) condition: the timestep cannot be larger than the transit time of the fastest wave across a cell. Hence high Alfvén speeds cause the timestep to become restrictively small. For this reason, codes typically set their inner boundary at $\sim 3R_p$. Another technique is by implementing the Boris limiter (Boris, 1970), which sets an artificial upper limit on the Alfvén speed by adding extra effective inertia to the plasma. This effectively damps fast Alfvén waves, ensuring stability at time-steps larger than would otherwise be possible.

A significant issue in any MHD code is the divergence free constraint: $\nabla \cdot \vec{B} = 0$. Solving the MHD equations, even in a conservative form, does not guarantee $\nabla \cdot \vec{B}$ remains zero, which can lead to errors in the overall solution (J.U Brackbill and Barnes, 1980; Toth, 2000). There are a number of methods to further overcome this, ranging from ‘keeping $\nabla \cdot \vec{B}$ small’ (Powell et al., 1999), to periodically ‘cleaning’ $\nabla \cdot \vec{B}$ from the solution (Jeremiah Brackbill, 1985), to staggering the grid variables, placing electromagnetic quantities on the grid faces or edges and the fluid quantities on the cell centres (Yee, 1966).

Boundary Conditions

The boundary conditions are an important facet in numerical simulations. In GMHD simulations of the magnetosphere, the outer boundaries are usually in the solar wind. These simulations

do not solve the MHD equations completely down to the atmosphere, stopping at an inner boundary approximately $3R_E$ from the Earth.

The outer boundaries consist of Von Neumann, or free-flow, boundary conditions for all quantities: the gradient of the quantity over the boundary is zero. This allows plasma to flow freely out of the box. The exception is the Sunward boundary, where the solar wind parameters are applied. By setting the values of ρ , $\vec{v}$, T , and $\vec{B}$ on the Sunward boundary, it allows new solar wind parameters to be advected into the simulation, with the exception of the B_x component. Since the solar wind velocity is predominantly in the $\vec{x}$ direction, B_x does not advect with the rest of the fluid. To overcome this, GMHD codes typically initialise the simulation with a fixed, background B_x .

The type of solar wind driving in GMHD simulations falls approximately into two categories. Ideal events usually consist of applying idealised solar wind parameters and studying how the magnetosphere responds. For example, Laitinen et al. (2007) used GUMICS-4 to investigate the position of reconnection on the dayside magnetopause under different idealised solar wind conditions. In that work, they either kept the plasma parameters constant and rotated the IMF, or vice versa. This kind of analysis is useful to identify specific causes of magnetospheric phenomena, as different causes (e.g. a particular IMF direction) and their effects (e.g. position of reconnection) can be more easily identified.

In reality, the solar wind is much more variable than such ideal events. This variability is ultimately what drives variability in the magnetosphere. Global processes such as the Dungey cycle rely on changes inflicted at the magnetopause to propagate through the magnetosphere, which has associated timescales. If these timescales are of the same order or longer than the timescales associated with the variability in the solar wind, this variability is likely to have an impact on the magnetosphere as a whole. Hence most codes can take observed solar wind parameters from an upstream monitor, such as ACE, Wind or DSCOVR, located at the 1st Lagrange Point (L1), approximately $200 R_E$ upstream (E. Stone et al., 1998).

The Sunward boundary sits typically $30 R_E$ Sunward of the Earth (Joachim Raeder, 2003). This is far enough upstream to ensure that the bow shock does not reach the Sunward boundary under typical conditions. However, it is much closer to the Earth than the position of the upstream monitors. The solar wind needs to be propagated from the upstream observer to the Sunward edge of the simulation domain in order to feed the solar wind into the simulation. There are a number of solar wind propagation techniques (e.g. Weimer et al. (2003)), though

most codes either uses a simple ‘ballistic’ propagation (a translation with the solar wind speed), or assume a phase front whose x component is a linear combination of the y and z magnetic field components (J. G. Lyon, Fedder, and Mobarry, 2004).

The inner boundary is usually defined by a dipole magnetic field and interfaced with an ionosphere code, where the field-aligned currents and flows are mapped from the inner boundary to the boundaries of the ionosphere.

From these flows and currents, an electric potential is solved over a separate ionospheric grid, and then mapped back out as boundary conditions for the field and fluid velocity (as $\vec{E} \times \vec{B}$ drift) (Joachim Raeder et al., 2008; P. Janhunen et al., 2012; Ridley, Gombosi, and Dezeew, 2004).

Most MHD codes (P. Janhunen et al., 2012; J. G. Lyon, Fedder, and Mobarry, 2004; Ridley, Gombosi, and Dezeew, 2004) treat the ionosphere as an electrical circuit, using a two dimensional thin shell approximation,

$$\nabla \cdot (\boldsymbol{\Sigma} \cdot \nabla \phi) = -j_{\parallel} \sin \delta. \quad (1.46)$$

This requires solving equation (1.46) for the electrostatic potential ϕ over the surface of a sphere, where $\boldsymbol{\Sigma}$ is the conductance tensor, $j_{\parallel}$ is the field-aligned current density, and $\cos \delta = \hat{B} \cdot \hat{r}$.

The conductance tensor ($\boldsymbol{\Sigma}$) is composed of the height integrated Hall and Pedersen conductances (Σ_H and Σ_P) for the ionosphere. These are usually calculated using empirical relations (e.g. the Knight relation (Knight, 1973), which estimates the conductance due to electron precipitation) or from an ionosphere-thermosphere model, which calculates the conductances from self-consistently from electron-neutral collisions (Joachim Raeder, 2003).

1.4.3 Global MHD Codes Verification

Having developed these codes for the magnetosphere purpose, they have been heavily used to simulate the magnetosphere under a wide range of solar wind conditions.

Given the fact that MHD is not a complete description of the physics present in the magnetosphere, it is very important to verify the results of the simulation against experimental data wherever possible.

One method is to compare the simulation with spacecraft observations with ‘virtual spacecraft’. By applying a solar wind profile as observed by an upstream monitor and sampling the simulation

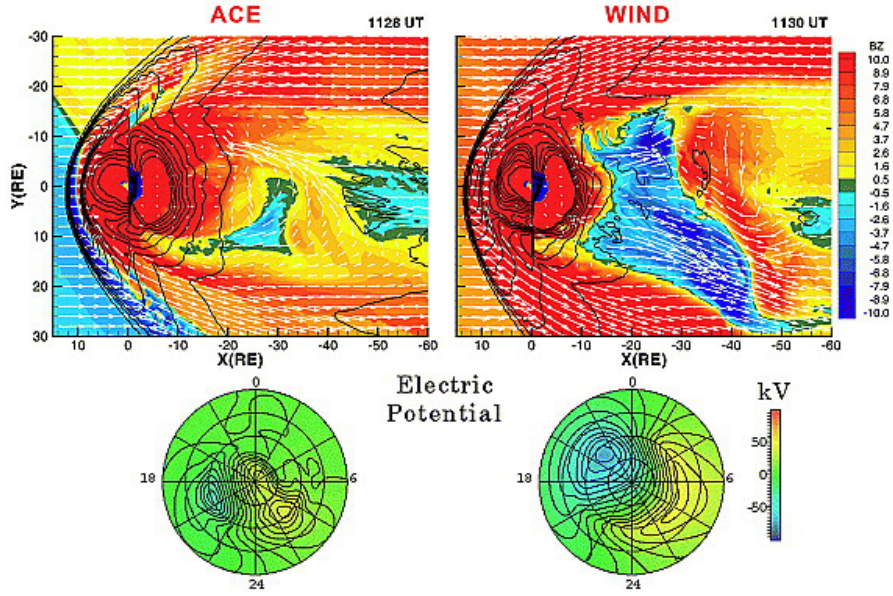


Figure 1.10: An OpenGGCM simulation of a magnetic cloud event on 24th May 2000, showing a cut in the ecliptic plane, with B_z shown as the filled colour contours, flow velocity as white arrows, and thermal pressure as the solid black contour. The ionospheric potential is shown on the bottom row. On the left, OpenGGCM is driven by ACE data; on the right, driven by WIND data. Both columns show the simulated magnetosphere at the same time, but have vastly different B_z , not only in the solar wind, but also in the tail. Figure from (Ashour-Abdalla et al., 2008).

at the same location and time of a real spacecraft, the observed time series can be compared to the simulated one (e.g. Lopez et al. (2007), Facsó et al. (2016), and Honkonen et al. (2013)). However this must be done with care. One of the reasons for this is the inaccuracy of the solar wind input into the GMHD simulation. Figure 1.10 shows two vastly different results from OpenGGCM simulations driven by solar wind observed by different spacecraft at the same time (Ashour-Abdalla et al., 2008).

Considering each GMHD code solves equivalent MHD equations, but with different approaches, Honkonen et al. (2013) tests all of the 4 mentioned codes against an event in February 2004 where the IMF B_z varies between -8 nT to 8 nT. They compare the simulations to observations by using virtual spacecraft, sampling the observed parameters in the simulation at the location of Cluster 1 and Geotail spacecraft. They also compare the ionosphere to SuperDARN data, looking at the Cross Polar Cap Potential. In general, these codes perform best on the dayside, but this decreases with distance downtail. Furthermore, there are qualitative differences in

large-scale plasmoid formation between the four codes (figure 1.11).

The good agreement on the dayside is generally because MHD simulations do well in areas which are dependent on conserved quantities such as mass, momentum and energy (Janhunen and Palmroth, 2001). However, they do not perform well in areas of parallel vs perpendicular physics, or when multi temperature physics is needed (Janhunen and Palmroth, 2001).

The other method is to compare features and parameters in the GMHD model with empirical models (García and Hughes, 2007; Gordeev et al., 2013; Gordeev et al., 2015). Gordeev et al. (2015) is one of the more comprehensive studies, comparing all four previously mentioned GMHD codes against 6 key magnetospheric parameters: the sub-solar and high latitude magnetopause position, the lobe magnetic field strength, the plasma sheet pressure, and ionospheric cross polar cap potential and total field-aligned current. They find that GMHD models agree with empirical models on the whole, although each different model has different strengths and weaknesses. They perform very well in some aspects such as the magnetopause position, but not performing as well in the magnetotail.

1.4.4 Code Coupling and Extended MHD

Given the simplifications made in MHD, there have been a number of attempts to include physics into GMHD simulations not present in MHD. This usually consists of coupling to local models, whose boundaries and internal dynamics are driven by the GMHD simulation to some extent.

The most common case of code coupling is the ionosphere, which acts as the inner boundary to the GMHD simulation, and is driven by the field-aligned currents generated in the magnetosphere. Additionally, there have been more recent attempts at including extended MHD, such as the Hall effect, or embedding local kinetic simulations.

An increasing amount of effort has been aimed at improving GMHD simulations in the inner magnetosphere. The inner magnetosphere is host to a significantly non-Maxwellian population of particles trapped on closed field lines, creating the ring current. Particle drifts in this region, such as curvature and gradient drift, are much faster than the $\vec{E} \times \vec{B}$ drift provided in ideal MHD, and hence dominate the dynamics in this region. Furthermore, those drifts are proportional to the particle energy, and hence are not captured by single fluid MHD (De Zeeuw, 2004). The

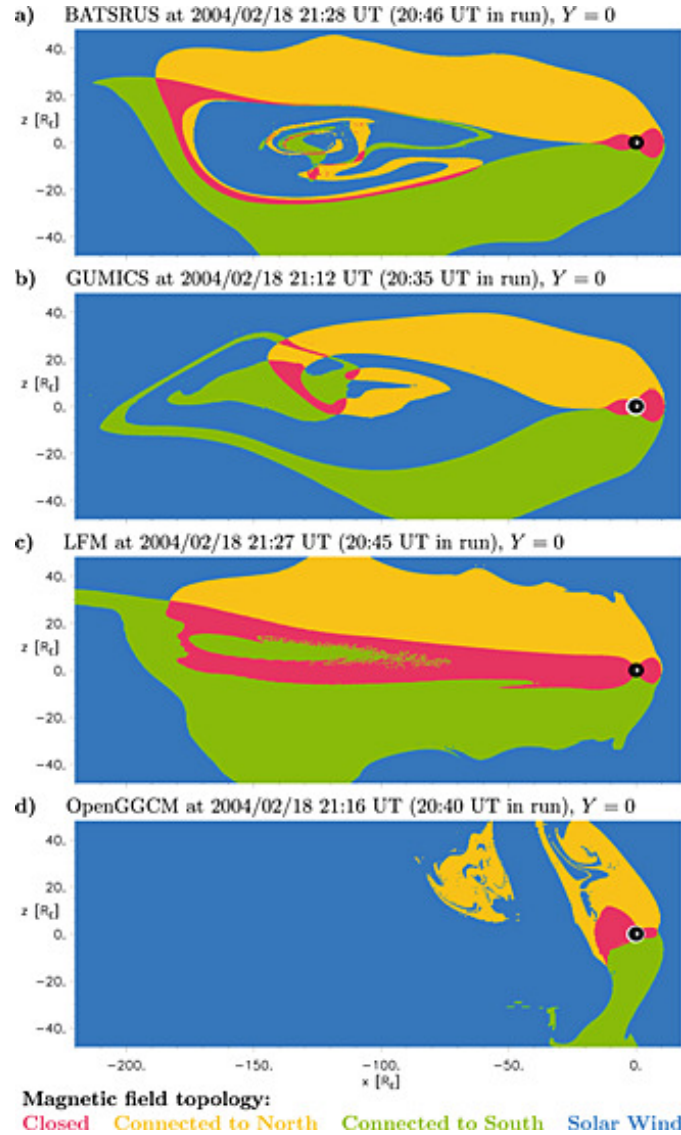


Figure 1.11: Differences in magnetic field topology between four MHD codes under the same solar wind input, with $B_z < 0$ at this time. The main differences are seen in the tail, with BATS-R-US (a), GUMICS (b) and LFM (c) showing closed (red), north- (yellow) and south-pole (green) connected field lines extending far down into the tail. BATS-R-US and GUMICS are arguably the most similar in this case, with twisted open, partially open and solar wind field lines regions intertwined. OpenGGCM (d) shows the majority of the field in the tail has already reconnected. Figure from Honkonen et al. (2013)

ring current is understood to be important in producing some of the current systems in the magnetosphere, which in turn have impacts on the global magnetosphere system.

To capture some of this missing physics, BATS-R-US can be coupled to the Rice Convection Model (RCM) (De Zeeuw, 2004), which includes the effect of the ring current. This is fed back into the GMHD simulation by adding the extra thermal pressure to the MHD solution.

Further development of BATS-R-US has included adding a Hall term (Daldorff et al., 2014) and anisotropic ion pressure (Meng et al., 2012). In more recent work, BATS-R-US has been coupled with an Embedded Particle-In-Cell code (EPIC). Rather than simulating the whole magnetospheric plasma using PIC, only a small region of the magnetosphere is targeted, allowing a full kinetic treatment in regions such as the dayside magnetopause or the magnetotail.

1.5 Outstanding Questions

The interaction between the solar wind and the magnetosphere is not fully understood, largely due to highly non-linear processes spanning over multiple spatial and temporal scales. The outer boundaries of the magnetosphere are the interface of the solar wind: understanding how the solar wind interacts with them is crucial in understanding how they inflict changes on the magnetosphere. Variability, both in the solar wind and the planetary field, can cause dynamics in the outer boundaries, determining the orientation of surfaces, and the position of reconnection. These are important as they determine a number of both large- and small-scale processes, which affect the magnetosphere globally.

Understanding this cause and effect between the solar wind and magnetospheric processes with spacecraft observations is a difficult task, as they only provide local measurements and are relatively sparsely located within the magnetosphere. Furthermore, they may have potential biases, with spacecraft positions potentially favouring one region over others, or only observing transient features.

Therefore, the aim of this thesis is to apply the Gorgon MHD code to planetary magnetospheres in order to better understand how the solar wind, and the rotation of the dipole affect the motion of the outer boundaries and generation of flux ropes.

2 The Gorgon Magnetohydrodynamic Code

Gorgon is a three-dimensional Eulerian resistive MHD code developed by the Plasma Physics Group at Imperial College London. It was originally designed to simulate laboratory plasmas: dense, collisional plasmas found in Z-pinches (Chittenden et al., 2004; Jennings, 2006; Jennings et al., 2010), laser plasma interactions (R. A. Smith et al., 2007) or magnetic tower jets produced in astrophysical laboratories (Ciardi et al., 2007).

This chapter is aimed at explaining the computational aspects and numerical methods of the code, as well as the magnetospheric implementation and comparison to other models. Given the original Gorgon code, my role was to adapt the existing MHD and hydrodynamic solvers to simulate planetary magnetospheres. Hence, in order to separate my development contributions with the already established code, they are listed here for clarity. They are as follows:

- Implemented standard numerical tests and evaluated Gorgon's capabilities
- Implemented the dipole magnetic field source with dipole precession
- Determined the inner boundary and its effect on the simulated magnetosphere
- Implemented a variable solar wind input, both idealised and from an upstream spacecraft (e.g. ACE, DSCOVR)
- Performed benchmarking of the code in a magnetospheric context

Gorgon solves the resistive MHD equations (equations (2.1) to (2.6)) on a three-dimensional uniform Cartesian grid using finite volume methods. Due to its roots in laboratory plasmas, Gorgon has features not usually found in GMHD magnetosphere codes: for example, a vacuum

wave equation solver and extra heating terms in the energy equation. The equations are as follows,

$$\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \vec{v}) = 0 \quad (2.1)$$

$$\frac{\partial \vec{p}}{\partial t} + (\vec{v} \cdot \nabla) \vec{p} = -\nabla (P_p + P_e) + \vec{J} \times \vec{B} \quad (2.2)$$

$$\frac{\partial \epsilon_p}{\partial t} + \nabla \cdot (\epsilon_p \vec{v}) = -P_p \nabla \cdot \vec{v} - \Delta_{pe} \quad (2.3)$$

$$\frac{\partial \epsilon_e}{\partial t} + \nabla \cdot (\epsilon_e \vec{v}) = -P_e \nabla \cdot \vec{v} + \eta |\vec{J}|^2 - \Lambda + \Delta_{pe} \quad (2.4)$$

$$\frac{\partial^2 \vec{A}}{\partial t^2} = -c^2 \nabla \times \nabla \times \vec{A} + c^2 \mu_0 \vec{J} \quad (2.5)$$

$$\text{where } \eta \vec{J} = -\frac{\partial \vec{A}}{\partial t} + \vec{v} \times \vec{B} \quad (2.6)$$

Equation (2.1) describes the mass continuity equation, with mass density ρ , momentum $\vec{p}$ and velocity $\vec{v}$. Equation (2.2) describes the momentum conservation equation, containing forces due to the electron and ion pressures $P_{p,e}$, and the $\vec{J} \times \vec{B}$ force, with current density $\vec{J}$ and magnetic field strength $\vec{B}$. Equations (2.3) and (2.4) describe the proton and electron internal energy equations, with the proton and electron internal energies, ϵ_p and ϵ_e . Equation (2.5) and equation (2.6) describe the magnetic induction equation and a simple Ohm's law, where $\vec{A}$ is the vector potential, and c is the speed of light and μ_0 is magnetic permeability of free space.

Though these equations are theoretically the equations of resistive MHD, they differ from the MHD equations given previously, equations (2.1) to (2.6), in a number of ways. The first is the use of conservative variables: the mass density ρ , momentum $\vec{p}$ and internal energy ϵ are all conserved quantities under certain conditions. This allows the use of solvers which are designed to conserve these variables. Secondly, the energy equations (2.4) to (2.3) describe the general evolution of the internal energy, rather than assuming an adiabatic energy equation as in equation (1.18). Finally, the magnetic induction equation is cast in vector potential form, and includes the full Ampère-Maxwell law which includes the displacement current.

The core code is built around a hydrodynamic solver, with the ability to switch the magnetic solver on or off dependent on the situation. Hence, the code is split into two parts: a hydrodynamic solver, where the fluid properties of the plasma are updated by pressure and magnetic forces; and a magnetic solver, where the magnetic field is advected and diffused through the plasma, or propagated electromagnetically through a vacuum.

2.1 Hydrodynamics Solver

The hydrodynamic part is contained in equations (2.1) to (2.3), which describe the time-evolution of the fluid momentum and energy. These are solved in a Lagrangian-remap scheme, consisting of a Lagrangian step where the plasma is accelerated and heated, followed by a remap where the plasma is advected throughout the Eulerian grid.

The Lagrangian step solves equations (2.2) to (2.3) for each cell on the Eulerian grid. This involves calculating the forces and heating terms in each cell, and updating the momentum, proton and electron energies, using equations (2.7) to (2.8), which are equivalent to equations (2.2) to (2.3) without the advection terms.

$$\frac{d}{dt} (\rho \vec{v}) = -\nabla (P_p + P_e) + \vec{J} \times \vec{B} \quad (2.7)$$

$$\frac{d\epsilon_p}{dt} = -P_p \nabla \cdot \vec{v} - \Delta_{pe} \quad (2.8)$$

$$\frac{d\epsilon_e}{dt} = -P_e \nabla \cdot \vec{v} + \eta |\vec{J}|^2 - \Lambda + \Delta_{pe} \quad (2.9)$$

The force equation equation (2.7) calculates the pressure force and $\vec{J} \times \vec{B}$ force, and applies the Boris correction to overcome issues of high Alfvén speeds. Devised by Boris (1970), the Alfvén correction damps high velocity Alfvén waves by adding extra artificial inertia to the plasma. This manifests itself in the code by multiplying the force terms on the right hand side of equation (2.7) by a factor of $(1 + (v_A/c_{\max})^2)^{-1}$, where v_A is the Alfvén speed and $c_{\max}$ is the maximum Alfvén speed, typically set to 2×10^6 m/s.

The energy equation in Gorgon is split into two parts, the proton equation (2.8) and electron equation (2.9) energy equations. The splitting was originally to account for extra heating terms that are present in collisional plasmas, such as Ohmic heating ($\eta |\vec{J}|^2$), optically thin radiation losses Λ , and proton-electron energy exchange (Δ_{pe}). These features are still available in the code, but are either disabled, or negligibly small. Gorgon has the capability to include complex equations of state. However, for the magnetosphere application, a fully ionized ideal hydrogen plasma is used, where the energy density, $\epsilon_{p,e}$, is related to the pressure by the ideal gas law, $P_{p,e} = (\gamma - 1)\epsilon_{p,e}$, where $\gamma = 5/3$.

In order to capture shocks, Gorgon uses a Von Neumann artificial viscosity (Benson, 1992), of the form in equation (2.10), where c_s is the local sound speed, Δv is change in velocity over the

cell interface, and α is a dimensionless parameter controlling the magnitude. This has the effect of decelerating the plasma through shocks, and providing the irreversible heating. Numerically, it also reduces unphysical overshoots and oscillations present downstream of a shock region.

$$P_{\text{visc}} = -\alpha \rho c_s \min(0, \Delta v) \quad (2.10)$$

Such artificial viscosities are relatively common in MHD codes and hydro-codes (e.g. Arber et al. (2001) and Robinson et al. (2008))

Once the Lagrangian calculation is performed, the quantities are then advected through the Eulerian grid. This takes the form of a remap operation, where the expanded fluid parcel calculated in the Lagrange step is spread, or remapped, back onto the original grid points. This is done in a flux conservative manner as described by Van Leer (1977), using a 3rd order advection routine.

The high order advection routine also allows the shock to be captured over a few grid cells. However, during the momentum remap, it is only conservative in flux, not kinetic energy. This creates a situation where the kinetic energy is not strictly conserved through the remap step. This is most prominent in the vicinity of shocks, where the advection routine switches to low order to avoid spurious oscillations. A fix for this is known as the DeBar correction (DeBar, 1974; Niederhaus et al., 2009), which attempts to recoup the change in energy caused by the remap. It calculates the kinetic energy before the remap step, which is then advected with the rest of the fluid quantities, and hence is conserved. This is compared to the kinetic energy from the remapped momentum and mass density: the difference between the two is added to the ion energy. This ensures better energy conservation at shocks.

2.2 Magnetic Field Solver

The magnetic field solver is split into two steps, advection and resistive diffusion/vacuum solver. Gorgon uses the vector potential, $\vec{A}$, to calculate the magnetic field, where $\vec{B} = \nabla \times \vec{A}$. Hence, the vector potential representation of the magnetic induction equation is used, as shown in equation (2.5).

The vector potential, along with a staggered grid, is used to ensure the magnetic divergence condition, $\nabla \cdot \vec{B} = 0$, is satisfied throughout the simulation. In a staggered grid, first formulated

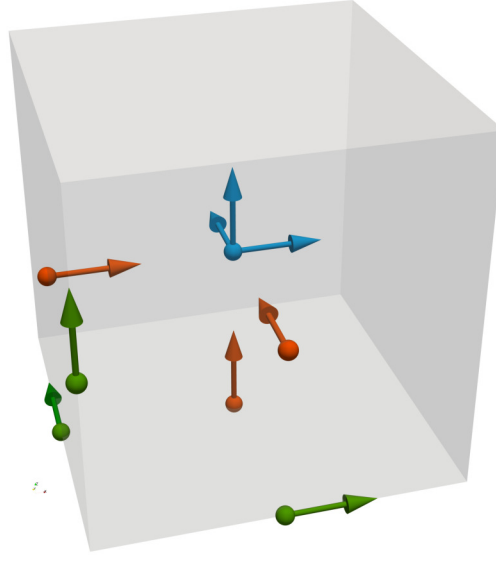


Figure 2.1: The staggered grid used in Gorgon. The cell is shown as the white box. Each sphere denotes the position of the cell (blue), face (red) and edge (green) centred variable. Cell centred variables are co-located in the centre of the cell. Face and edge centred variables are not co-located: each $\hat{x}$, $\hat{y}$ and $\hat{z}$ vector are placed on different faces/edges. For example, $\vec{B}$ is a face centred variable; the B_x component is located on the $y - z$ face of the cell.

by Yee (1966), all electromagnetic variables, $\vec{A}$, $\vec{B}$, $\vec{J}$, and $\vec{E}$, are not located in the same place as the fluid quantities, ρ , $\vec{p}$, ϵ_e , ϵ_p . This is illustrated in figure 2.1, which shows the different positions of cell, face and edge centred variables. The cost of a staggered grid is that averaging needs to occur in order to calculate variables which combine both magnetic and fluid quantities. For example, the $\vec{J} \times \vec{B}$ force on the fluid applies on the cell centre, which requires averaging the face centred $\vec{B}$ and the edge centred $\vec{J}$ vectors. This alone does not preserve $\nabla \cdot \vec{B} = 0$ entirely: formulations of the MHD equations using $\vec{B}$ are not necessarily conservative in $\nabla \cdot \vec{B}$ and, if care is not taken, magnetic divergence can build up throughout the simulation, leading to unphysical results (J.U Brackbill and Barnes, 1980). Using the vector potential reduces this error, as the vector potential allows the divergence condition to be satisfied automatically to machine accuracy: $\nabla \cdot \vec{B} = \nabla \cdot (\nabla \times \vec{A}) = 0$. As mentioned in section 1.4, some codes utilise a divergence cleaning method, which periodically remove divergence from the simulation, or solvers specially designed to be conservative in $\nabla \cdot \vec{B}$; using the vector potential and staggered grid removes the need for such algorithms.

2.2.1 Magnetic Advection

The advection equation, derived directly from the ideal MHD induction equation (2.11) in vector potential formulation is shown in equation (2.12). This needs to be conservative in magnetic flux, and thus uses the same van Leer advection routine as in the remap step of fluid equations:

$$\frac{\partial \vec{B}}{\partial t} = \nabla \times \vec{v} \times \vec{B} \quad (2.11)$$

$$\Rightarrow \frac{\partial \vec{A}}{\partial t} = \vec{v} \times \vec{B} \quad (2.12)$$

2.2.2 Resistive Diffusion

The second part solves the Ampère-Maxwell law, including Ohm's law, shown in equation (2.5). This allows the magnetic field to diffuse, and also propagate through vacuum regions. It is solved in the form of a set of coupled partial differential equations between $\vec{A}$ and $\vec{E}$:

$$\frac{\partial \vec{E}}{\partial t} = c^2 \nabla \times \vec{B} - \frac{\mu_0 c^2}{\eta} \vec{E} \quad (2.13)$$

$$\frac{\partial \vec{A}}{\partial t} = -\vec{E} \quad (2.14)$$

where equation (2.13) is Ampère-Maxwell's law, and equation (2.14) is the definition of the vector potential from the electric field $\vec{E}$ in a quasi-neutral plasma ($\nabla \phi = 0$, where ϕ is the scalar electric potential).

By assuming that $\nabla \times \vec{B}$ stays constant within the integration timestep Δt , equation (2.13) is analytically integrated to obtain the iterative equation (equation (2.15)) used to solve the magnetic field diffusion (Jennings, 2006).

$$\vec{E}(\Delta t) = \vec{E}_0 e^{-\frac{\mu_0 c^2}{\eta} \Delta t} + \frac{\eta}{\mu_0} \nabla \times \vec{B} \left(e^{-\frac{\mu_0 c^2}{\eta} \Delta t} - 1 \right) \quad (2.15)$$

Using equation (2.15), discretised formulations of the resistive diffusion equations are obtained, which are used to update the magnetic and electric field:

$$\vec{A}_{n+1} = \vec{A}_n - \Delta t \left(\vec{E}_n e^{-\frac{\mu_0 c^2}{\eta} \frac{\Delta t}{2}} + \frac{\eta}{\mu_0} \nabla \times \vec{B}_n \left(e^{-\frac{\mu_0 c^2}{\eta} \frac{\Delta t}{2}} - 1 \right) \right) \quad (2.16)$$

$$\vec{E}_{n+1} = \vec{E}_n e^{-\frac{\mu_0 c^2}{\eta} \Delta t} + \frac{\eta}{\mu_0} \nabla \times \vec{B}_{n+1} \left(e^{-\frac{\mu_0 c^2}{\eta} \Delta t} - 1 \right) \quad (2.17)$$

Table 2.1: Limits of the resistive diffusion solver

		$\vec{A}_{n+1}$	$\vec{E}_{n+1}$
Vacuum Limit	$\eta \rightarrow \infty$	$\vec{A}_n + \Delta t \left(\vec{E}_n - \Delta t c^2 \nabla \times \vec{B}_n \right)$	$\vec{E}_n - \Delta t c^2 \nabla \times \vec{B}_{n+1}$
Plasma Limit	η	$\vec{A}_n + \Delta t \frac{\eta}{\mu_0} \nabla \times \vec{B}_n$	$\frac{\eta}{\mu_0} \nabla \times \vec{B}_{n+1}$

The current is deduced from Ohm's Law, $\vec{J} = \vec{E}/\eta$.

The limits of the exponential functions in equations (2.16) to (2.17) provide a gradual transition between resistive diffusion and electromagnetic wave propagation (Jennings, 2006). The equations then reduce to the relations shown in table 2.1.

Electromagnetic waves propagate at the speed of light, c , which limits timestep. This timestep is typically much smaller than the fluid timestep. Hence the resistive diffusion solver is sub-cycled over the fluid timestep. However, internal tests have shown that the speed of light can be reduced within the simulation: typically, a speed of 4×10^6 m/s (Ciardi et al., 2007).

Gorgon defines a vacuum density, which forces the vacuum solution in the magnetic field solver and puts a lower limit on the range of plasma density in the simulation. Cells with mass densities below this threshold are said to be a 'vacuum region'. The mass density and momentum are all set to zero, and the resistivity is set such that the wave solution applies to this region.

2.3 Numerical Tests

In this section, the Gorgon code is applied to some common test cases, testing its capability to handle waves and shocks. Simulations of linear waves, shock tubes, and the Orszag-Tang MHD vortex are performed, and tested against the appropriate metric (e.g. an analytical solution). These tests have been designed to see how well the code copes in areas such as waves and shocks, and have been used extensively in other codes to determine their performance (e.g. Powell et al. (1999), Londrillo and Del Zanna (2000), Arber et al. (2001), and J. M. Stone et al. (2008)). For all MHD test problems, the resistive diffusion solver has been switched off, defaulting to ideal MHD. This is for two reasons: all these test problems are ideal MHD problems, and the resistive diffusion requires sub-cycling which resolves the speed of light, creating restrictively small timesteps due to the small simulation domain (of order ~ 1 m).

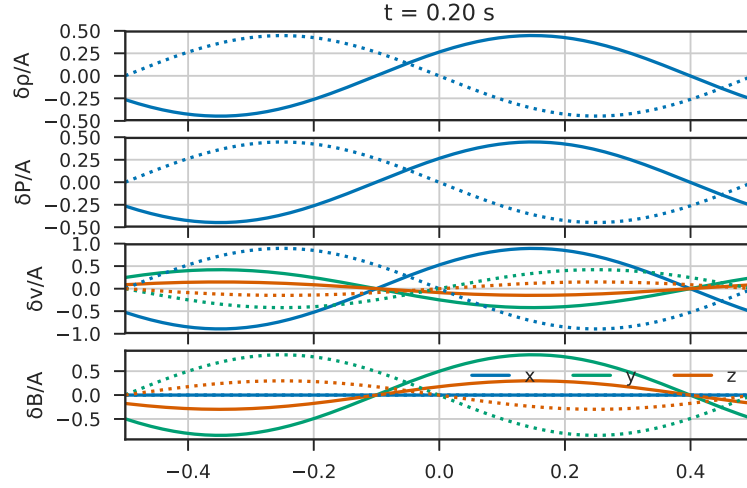


Figure 2.2: A fast magnetosonic wave ($c_f = 2$) wave test, propagating to the right. From top to bottom, it shows the deviation in mass density, pressure, velocity and magnetic field, normalised over the wave amplitude, A , with the initial conditions as dotted lines.

2.3.1 Linear Waves

In the linear wave test, small amplitude waves are initialised in the simulation domain and allowed to propagate according to the MHD equations being solved, as illustrated in figure 2.2. Small amplitude waves have well defined analytical solutions: their initial condition translated through the box at the wave speed. The propagation of linear waves tests one of the code's most basic aspects: how it transports mass, energy and magnetic flux through the system. Furthermore, it can highlight errors in the code by comparing symmetries: e.g. right-going waves should behave similarly to left-going waves. The results for the one-dimensional case are shown here, but it can be extended to three dimensions.

The three different types of MHD waves are tested (Alfvén, fast and slow magnetosonic), as well as the contact discontinuity, in a one-dimensional simulation. Each case starts in a background medium at rest, except the contact discontinuity which tests how well the code advects fluid. The general method is as follows, and has been used in multiple papers (e.g. Gardiner and J. M. Stone (2008) and J. M. Stone et al. (2008)). A wave of wavelength $\lambda = 1$ m is set up in a box

of length 1 m with the following initial conditions:

$$\begin{aligned}\rho &= \rho_0 + \alpha R_\rho \sin 2\pi x, \\ \vec{p} &= \vec{p}_0 + \alpha \vec{R}_p \sin 2\pi x, \\ \varepsilon &= \varepsilon_0 + \alpha R_\varepsilon \sin 2\pi x, \\ \vec{B} &= \vec{B}_0 + \sqrt{\mu_0} \alpha \vec{R}_B \sin 2\pi x,\end{aligned}$$

where $\rho_0 = 1 \text{ kg/m}^3$, $\vec{p}_0 = \vec{0}$ ¹, $P = \gamma^{-1} \text{ Pa}$ and $\vec{B}_0 = \sqrt{\mu_0}(\hat{x} + \sqrt{2}\hat{y} + 1/2\hat{z}) \text{ T}$. This gives the following wave speeds: $c_s = 1 \text{ m/s}$, $v_A = 1 \text{ m/s}$, $v_f = 2 \text{ m/s}$ and $v_s = 1/2 \text{ m/s}$. The amplitude of the wave is set to be much smaller than the background value: $\alpha = 10^{-6}$. The value of R is determined by the type of wave being simulated. The values used here can be found in Gardiner and J. M. Stone (2008). Figure 2.2 shows an example of a right travelling fast magnetosonic wave. The initial condition is shown as the dotted line, and the solution as $t = 0.2 \text{ s}$ is shown as the solid line. As expected, the solution is a simple translation of the initial condition to the right.

The simulation is run for one wave transit time, $t = \lambda/c_{\text{wave}}$ (c_{wave} is the relevant wave speed), where the solution should be identical to the initial condition. This is then compared to the initial condition by finding the L1 norm for every conserved variable, v , as given in equation (2.18). A total error is obtained by adding each variable L1 norm in quadrature, as given in equation (2.19). Since the magnetic field in SI units is much smaller in magnitude than the other variables, it is scaled by a factor of $\sqrt{\mu_0}$. This is done for a number of different grid sizes to see how the code performs under high and low resolutions.

$$\sigma_v = \sum_i |v_i^{\lambda/c} - v_i^0| \quad (2.18)$$

$$L_1 = \sqrt{\sum_v \sigma_v^2} \quad (2.19)$$

Figure 2.3, shows the total error of the linear wave test for each wave type, at different resolutions, ranging from $N = 12$ to 960 cells. The Alfvén wave and contact discontinuity modes perform the best, showing a power law relationship with the number of cells, as seen in other simulations (J. M. Stone et al., 2008). The two magnetosonic waves do not perform as well, levelling out in error at high resolutions. Whilst the error does level off, it is still ~ 2 orders of magnitude smaller than the wave amplitude. This aspect is not seen in other simulations (e.g. J. M. Stone

¹In the contact discontinuity, the background momentum is non-zero, $\vec{p}_0 = \hat{x}$

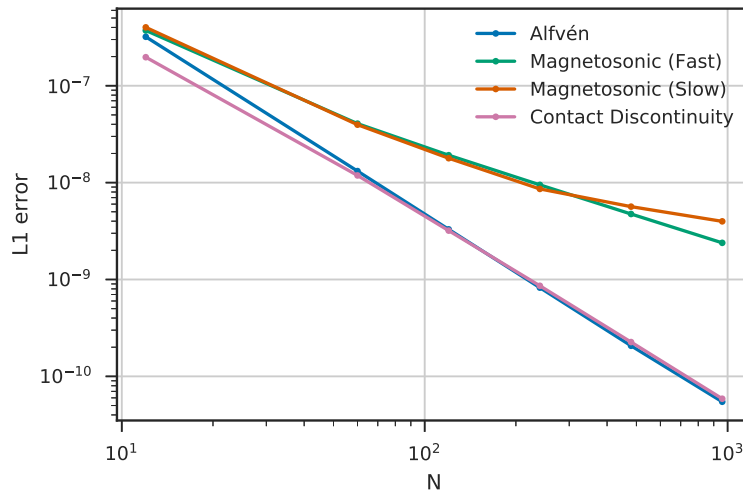


Figure 2.3: The L1 error as a function of number of cells in the simulation for each 1D small amplitude wave: the Alfvén, slow and fast magnetosonic waves, and the contact discontinuity. All L1 errors decrease with number of cells, with the Alfvén and contact discontinuity performing best. The magnetosonic waves appear to tail off as higher grid resolutions.

et al. (2008)), which also show power law relations. This is likely due to the artificial viscosity used in Gorgon (equation (2.10)), which acts to heat and decelerate the plasma in shocks. Whilst its effect is small in a smooth initial condition such as a low amplitude wave, it still provides a small amount of force and heating which distorts the waveform.

2.3.2 Shock Tubes

Another key aspect of such codes is their ability to capture shocks: a code which can capture waves correctly does not necessarily handle shocks well due to their highly non-linear nature. A number of shock tube set-ups have been devised to test the code's capability to handle shocks, which usually consists of two fluid regions of different parameters (ρ , P , etc.), separated by a discontinuity.

One such shock tube, Sod's shock tube (Sod, 1985), is a hydrodynamic test case. The initial conditions to the left and right of the discontinuity are given in table 2.2. In the literature, including Sod's original research (Sod, 1985), this initial condition is usually run with an ideal

Table 2.2: The initial conditions for the left and right regions in Sod’s and Brio-Wu’s shock tubes. Sod’s shock tube is unmagnetised $\vec{B} = 0$. Brio-Wu’s shock tube has the same plasma parameters as Sod’s, but with an embedded magnetic field.

	ρ (kg/m ³)	P (Pa)	$\vec{v}$ (m/s)	$\vec{B}$ (T) ⁽¹⁾
Left	1	1	0	$\sqrt{\mu_0} (0.75\hat{x} + \hat{z})$
Right	0.1	0.125	0	$\sqrt{\mu_0} (0.75\hat{x} - \hat{z})$

gas constant of $\gamma = 2$: this is done here too, in order to provide general comparison with the literature.

The system starts at rest, with a large pressure gradient across the discontinuity, which rapidly accelerates the gas forming a shock. The solution decomposes into three regions, consisting of a fast moving shock and a contact discontinuity moving in the $+x$ direction, and a rarefaction wave propagating the $-x$ direction. This is shown in figure 2.4, with the analytical solution show as the black dots. The blue and green curve shows two Gorgon simulations at different resolutions, 120 grid cells and 1200 grid cells respectively.

Overall, the simulation agrees well with the analytical solution, capturing the magnitude of each of the regions and the position of the shock. This suggests the code satisfies the shock jump conditions. It also captures the rarefaction wave well, whose profile is non-linear. The code deviates from the analytical solution in the vicinity of the shocks and discontinuities. This is expected as the grid spacing is finite and the shock position does not necessarily lie exactly on a grid cell. The numerical solution smears the shock over approximately 4 grid cells, as is best seen by the low resolution (blue) curve in figure 2.4. This is a function of the numerical scheme used, and, in the case of Gorgon, the size of the artificial viscosity. The viscosity only controls the width of the shock: the position is still determined by the upstream and downstream parameters.

The Brio-Wu shock tube extends Sod’s shock tube to include a magnetic field (Brio and C. Wu, 1988). Starting with the same initial conditions as Sod’s, the Brio-Wu shock tube adds a constant component normal to the discontinuity (B_x in this set-up), and a tangential component of opposite direction either side of the discontinuity. The full initial conditions are shown here in table 2.2: these have been chosen for direct comparison with Londrillo and Del Zanna (2000) ($\gamma = 5/3$).

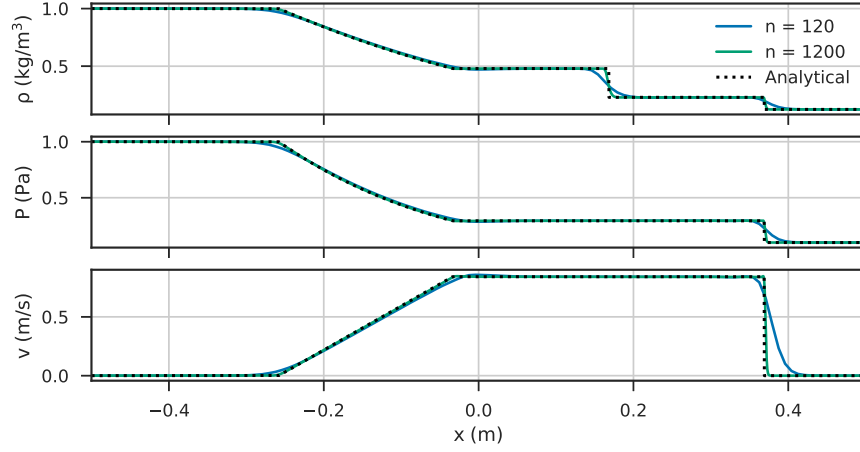


Figure 2.4: The 1D Sod's shock tube problem, comparing the analytical solution (black dots), to two Gorgon simulations at two different resolutions: blue with 120 grid cells, green 1200 grid cells. The numerical solution agrees with the analytical, capturing the pre and post shock plasma parameters, including the rarefaction wave. The numerical solution deviates from the analytical in the shock vicinity, where the shock is smeared to multiple grid cells. Additionally, there are slight over and undershoots at the edges of the rarefaction wave. Both are more pronounced the lower the resolution.

The presence of a magnetic field increases the number of types of waves present: Alfvén, slow and fast magnetosonic waves. This causes the evolving solution to be much more complex than in the hydrodynamic case. The numerical solution as simulated by Gorgon is shown in figure 2.5, at a resolution of 800 cells.

When compared to the solution presented in Londrillo and Del Zanna (2000), Gorgon agrees well, with the value of each plateau between the shocks and rarefactions agreeing. Importantly, the magnetic field component normal to the shock, (B_x), stays constant, implying $\nabla \cdot \vec{B}$ correctly handled. This shows that Gorgon can handle MHD shocks as well. However, the Gorgon solution exhibits a larger overshoot than seen in Londrillo and Del Zanna (2000), especially on the slow shock at $x \sim 0.6$ in figure 2.5, and shows more structure on the compound shock. This can be mitigated by using a larger artificial viscosity.

Overall, the Gorgon code has shown the capability to handle shocks. With the use of the artificial viscosity and Van Leer advection, the shock width is kept to within approximately four grid cells. There are codes that can capture shocks with smaller widths, but typically use more complicated solvers, like an approximate Riemann solver (e.g Powell et al. (1999) and J. M. Stone

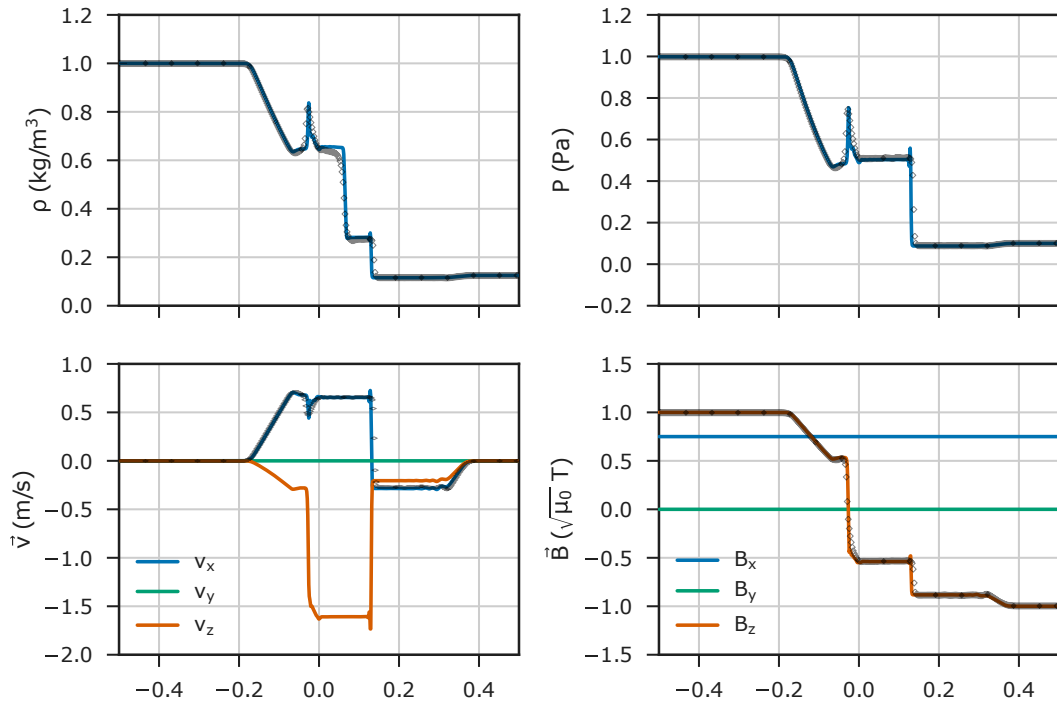


Figure 2.5: The 1D Brio-Wu shock tube problem, the magnetised extension of the Sod shock tube, with 800 grid points. By adding a magnetic field, the solution becomes much more complex with additional regions. The solution consists of 4 regions: a fast rarefaction wave (between $-0.18 \lesssim x \lesssim -0.08$), a slow compound wave (at $x \sim -0.01$), a contact discontinuity (at $x \sim 0.05$), a slow shock (at $x \sim 0.15$), and another fast rarefaction wave (at $x \sim 0.35$). The solution is compared to the Londrillo and Del Zanna (2000) MHD code (with 400 grid points), shown in the black diamonds.

et al. (2008)) or a scheme with higher order spatial gradients (J. G. Lyon, Fedder, and Mobarry, 2004), rather than the Lagrange-Remap formulation in Gorgon.

2.3.3 Orszag-Tang Vortex Test

The Orszag-Tang Vortex Test is a commonly used test problem for MHD codes as it gives a good sense of how codes handle waves, shocks and $\nabla \cdot \vec{B} = 0$. Originally devised by Orszag and Tang (1979) to study turbulence in MHD systems, and has been used as a test problem for a number of MHD codes (e.g. J. M. Stone et al. (2008), Arber et al. (2001), and Londrillo and Del Zanna (2000)).

The simulation starts with a uniform mass density and thermal pressure, such that the sound speed, $c_s = 1$. Embedded into the plasma is a rotational velocity field, with smaller rotational magnetic fields within. This is given by

$$\begin{aligned}\vec{v} &= -\sin(2\pi y) \hat{x} + \sin(2\pi x) \hat{y}, \\ \vec{B} &= -B_0 \sin(2\pi y) \hat{x} + B_0 \sin(4\pi x) \hat{y},\end{aligned}$$

where $B_0 = \sqrt{\mu_0/4\pi}$ T. This gives an Alfvén speed of $v_A = 1/\gamma = 3/5$.

The results of a 2 s simulation, at a resolution of 720×720 cells, is shown in figure 2.6. The initial condition shown in the top row of figure 2.6 shows a uniform mass density, and the rotational velocity and magnetic fields. As the simulation progresses, the solution becomes more complex, mainly due to the non-linear wave interactions and the formation of shocks.

Unfortunately, the Orszag-Tang Vortex does not have an analytical solution. Instead, a common approach is to compare it to other codes (Orszag and Tang, 1979; Londrillo and Del Zanna, 2000; Arber et al., 2001; J. M. Stone et al., 2008). Figure 2.7 shows such a comparison with the simulation in Londrillo and Del Zanna (2000). Whilst the overall structure of the thermal pressure agrees, there are some notable differences. These differences are unlikely to be significant. They are likely due to the different solvers and methods used, and is typical for such cross-simulation comparison.

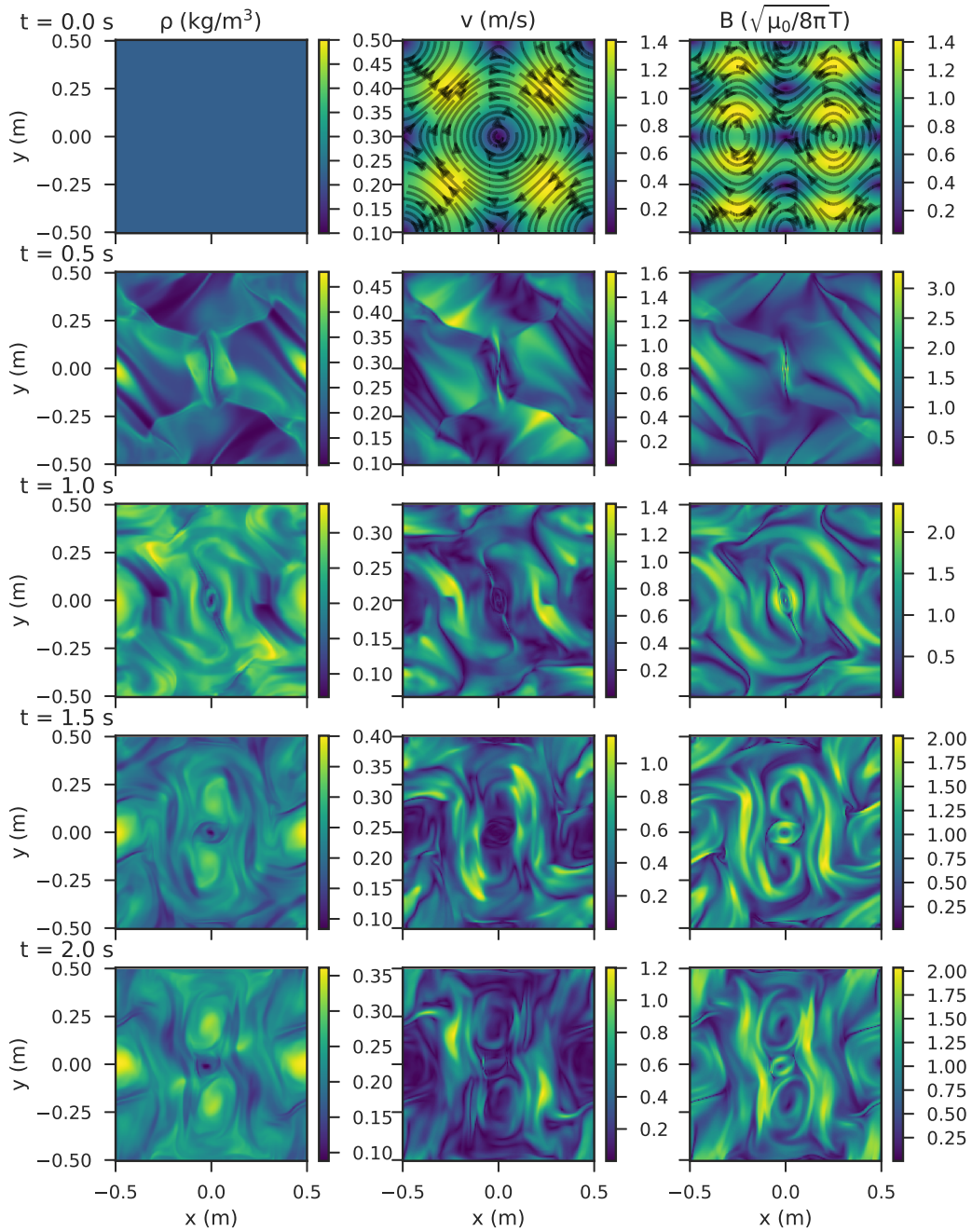


Figure 2.6: Two-dimensional numerical simulation of the Orszag-Tang Vortex. The three columns show the mass density, velocity and magnetic field, coloured by magnitude. Each row shows the solution at different times. The initial condition is shown in the top row, with velocity and magnetic streamlines shown in black. The initial condition consists of a velocity vortex, embedded within smaller magnetic vortices. As the simulation progresses, the solution becomes complex, causing strong non-linear wave interactions and shocks.

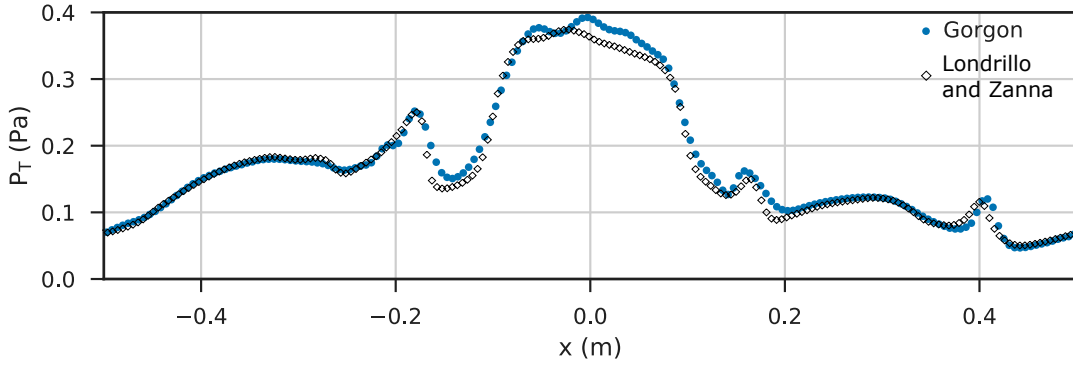


Figure 2.7: Comparison to Londrillo and Del Zanna (2000)’s MHD simulation of the Orszag-Tang Vortex, showing the thermal pressure along the $y = 0.4277$ line at $t = 0.5s$. The two codes show good agreement with the large-scale features, both showing peaks at $x \sim [-0.2, 0, 0.2, 0.4]$. The magnitudes are slightly different, which is likely a consequence of the different solvers and limiters used.

2.4 Comparison to Other GMHD Codes

Whilst there are a number of GMHD simulations, the focus will be on comparing Gorgon to the four ‘original’ codes the Lyon-Fedder-Mobarry (LFM) (J. G. Lyon, Fedder, and Mobarry, 2004), Grand Unified Magnetosphere Ionosphere Coupling Simulation (GUMICS-4) (P. Janhunen et al., 2012), Block Adaptive Tree Solar-wind Roe Upwind Scheme (BATS-R-US) (Powell et al., 1999) and Open General Geospace Circulation Model (OpenGCCM) (Joachim Raeder et al., 2008). These have been widely used, and have a proven track record. Such MHD often have a common heritage: for example, the teams developing GUMICS-4 and BATS-R-US shared information initially and hence have some similarities (e.g. they both use a version of Roe’s approximate Riemann solver).

All five codes solve the MHD equations, with Gorgon, BATS-R-US and OpenGCCM solving the resistive equations, and LFM and GUMICS-4 the ideal equations ($\eta = 0$). However, each of them differ in their implementation of these equations, with different solvers, grid and boundary conditions. These have been summarised in table 2.3.

Table 2.3: A comparison of the four main GMHD codes and Gorgon, comparing differences in grid, equations and solvers. Cartesian (Cart.); Staggered (Stag.); Constrained Transport (CT); Vector Potential (Vector Pot.).

Code	Grid	Energy Eq.	$\nabla \cdot \vec{B}$ Control	Split Field
BATS-R-US	Block-Adapted Cart.	Total	8-Wave Riemann Solver	Yes
GUMICS-4	Cell Adaptive Cart.	Total	Projection Cleaning	Yes
LFM	Deformed Spher.	Plasma	Stagg. Grid, CT	No
OpenGGCM	Stretched Cart.	Plasma	Stagg. Grid, CT	No
Gorgon	Uniform Cart.	Internal	Stagg. Grid, Vector Pot.	No

2.4.1 Choice of Solvers

All four of the CCMC codes use different finite volume, ‘shock-capturing’ solvers. They involve calculating the fluxes at each cell interface, which are used to update the solution in a conservative manner. Similar to Gorgon, these fluxes are calculated with a high order scheme, switching to a low order scheme in the vicinity of shocks. This section will note some major points regarding each of these solvers; the exact details is out of the scope of this thesis.

Due to their close communication in early development, BATS-R-US and GUMICS-4 use a variation of Roe’s approximate Riemann solver (Roe, 1981) to solve the MHD equations. However, they have since diverged, having developed different variations of the solver. Riemann solvers were designed to solve the Riemann problem: two regions of fluid separated by a discontinuity. The shock tubes in section 2.3 are an example of a Riemann problem. Hence, these are particularly good at resolving shocks, as they incorporate the relevant waves into the calculation of the interface fluxes. However, they are harder to implement and adapt to different physics.

GUMICS-4 still uses Roe’s solver, but it switches to the diffusive Harten-Lax-van Leer (HLL) solver (Harten, Lax, and Leer, 1983) when non-physical states occur. Though its 7-wave Riemann solver preserves $\nabla \cdot \vec{B}$ to machine accuracy, $\nabla \cdot \vec{B}$ tends to accumulate over time, which is then removed using the projection method (J.U Brackbill and Barnes, 1980). This ‘cleaning’ of divergence is not performed every timestep, but does account for $\sim 10\%$ of total runtime (P. Janhunen et al., 2012).

BATS-R-US however, has developed a slightly different version of the MHD equations which

preserve $\nabla \cdot \vec{B}$ numerically (Powell et al., 1999). Its 8-wave Riemann solver incorporates $\nabla \cdot \vec{B}$ source terms, which are not necessarily zero numerically. The solver includes an eighth wave mode which corresponds to the $\nabla \cdot \vec{B}$ wave, the idea being that any magnetic divergence is advected away and eventually out of the simulation domain.

LFM and OpenGGCM use solvers which introduce minimal $\nabla \cdot \vec{B}$ into the simulation, and is kept to round-off error using a staggered grid (J. G. Lyon, Fedder, and Mobarry, 2004; Joachim Raeder et al., 2008). Gorgon employs the same technique, using a staggered grid and the vector potential.

2.4.2 Energy Equation

These codes also differ in their choice of energy equation. BATS-R-US and GUMICS-4 are fully conservative. Both solve total energy equations, with internal, kinetic and magnetic energy contributing: $\varepsilon_T = B^2/2\mu_0 + \rho v^2/2 + P/(\gamma - 1)$. Non-physical states are a possibility in regions where the thermal energy is much smaller than the magnetic and kinetic energies, since the subtraction of the large total and magnetic energies can create negative internal energies and hence pressures. This is particularly present in the inner magnetosphere, where the plasma beta is very high. GUMICS-4 mitigates this by switching solvers when these non-physical states occur.

In codes with a staggered grid, like LFM, OpenGGCM and Gorgon, this formulation is more problematic as the numerical error associated with interpolation can create a large mismatch in magnetic and total energy. Hence LFM and OpenGGCM use the plasma energy equation, $\varepsilon_P = \rho v^2/2 + P/(\gamma - 1)$, which only includes the internal and kinetic energies and are not fully conservative in energy.

Gorgon uses an internal energy equation, which does not include the kinetic nor magnetic energy. This means total energy is not fully conserved, though it is usually conserved within 10% (Ciardi et al., 2007). However, energy conservation is greatly improved using the deBar correction, which enforces total energy conservation in the hydrodynamic case.

2.4.3 Time-stepping and Grid

All of the codes use explicit solvers, meaning the timestep is limited by the Courant-Friedrichs-Lewy (CFL) condition,

$$\Delta t < \min (\Delta x_i / v_i), \quad (2.20)$$

where Δt is the timestep, Δx_i is the cell size, v_i is the fluid speed plus the fastest wave speed and i denotes each grid cell. In Gorgon, this criteria is satisfied, but since its grid is uniform Cartesian, the cell size Δx is constant throughout the simulation domain, and the timestep is only dependent on the fastest wave speed.

The BATS-R-US uses a block adapted Cartesian grid (Powell et al., 1999). This total grid consists of a set of smaller connected Cartesian grids of different resolution which communicate with each other to create the whole simulation grid. The LFM grid, shown in figure 2.8a, is a deformed spherical grid (J. G. Lyon, Fedder, and Mobarry, 2004). Near the planet, the grid has a high resolution, almost spherical grid. Moving outwards, the grid deforms to capture the asymmetry of the magnetosphere along the x_{GSE} axis, keeping a high resolution on the dayside, but elongating in the x_{GSE} direction in the tail. On this grid, the LFM code typically uses 8th order differencing schemes in order to calculate gradients, which is significantly higher than in other codes. OpenGGCM's grid is most similar to Gorgon's, employing a stretched Cartesian grid with high resolution near the planet and low resolution in the solar wind and far down-tail. The high resolution of these grids near the planet can exacerbate the low timestep due to the high Alfvén speed in the inner magnetosphere. Therefore, these codes use the Alfvén correction (otherwise known as the Boris limiter) (Boris, 1970), to cap the Alfvén speed.

Finally, GUMICS-4 uses Adaptive Mesh Refinement (AMR) on a Cartesian grid and temporal sub-cycling, which allows GUMICS to be able to run efficiently on a local PC. However, unlike the other global magnetosphere codes, it is not parallelised. In AMR, the simulation grid adapts to the local parameter regime. The advantage of this is that it can target certain regions of the magnetosphere that need a high spatial resolution, splitting the grid cells locally into smaller cells. This allows the grid to be large in regions such as the solar wind, and small over the bow shock, inner magnetosphere and in the tail, potentially reducing the number of calculations needed. This can be seen in the two-dimensional cut of the GUMICS-4 grid in figure 2.8b, where the grid has refined in order to capture tail reconnection, and the high field magnetic gradients in the inner magnetosphere. The choice on whether to refine, or re-coarsen, a grid cell

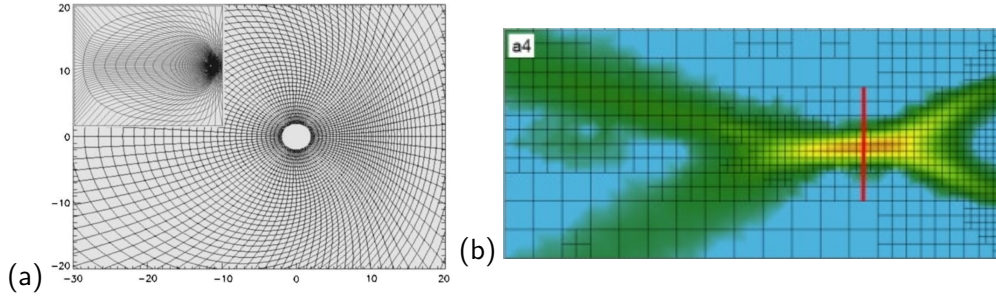


Figure 2.8: (a) The LFM adapted (deformed) grid (J. G. Lyon, Fedder, and Mobarry, 2004). (b) The GUMICS-4 adaptive grid (P. Janhunen et al., 2012). The colour shows the current density magnitude in the magnetotail sliced along the $x_{\text{GSE}} - z_{\text{GSE}}$ plane, with the Earth to the right of the figure. The grid becomes very fine in regions of high current density and closer to the Earth, allowing the tail current sheet to be more accurately resolved.

is decided by a number of dimensionless MHD parameters which capture the relevant gradients. A consequence of this is that only a 1st order advection scheme is used, which causes fast waves moving across the grid to diffuse rapidly (P. Janhunen et al., 2012). Temporal sub-cycling also reduces the simulation time. GUMICS-4 uses explicit solvers: its timestep is bounded by the CFL condition. By specifying a global timestep, typically $\Delta T = 1\text{s}$, the update for each cell is sub-cycled using a smaller timestep which satisfies the CFL condition locally. This method is unstable however if the CFL condition changes rapidly over a timestep (P. Janhunen et al., 2012).

2.5 Simulating the Magnetosphere

2.5.1 Boundary Conditions

The outer boundaries in Gorgon are free-flow (von Neumann) boundary conditions on all boundaries except the Sunward boundary. The simulation typically spans from $-30 R_E$ Sunward (x_{GSE}) and $70 R_E$ to $100 R_E$ anti-Sunward, and $30 R_E$ to $40 R_E$ in the dawn and dusk (y_{GSE}) and z_{GSE} direction. The grid size is typically $0.2 R_E$ to $0.5 R_E$. The solar wind parameters are applied to the Sunward boundary ($-x_{\text{GSE}}$), setting ρ , $\vec{v}$, T_i , T_e , B_y and B_z on the edge of the Sunward x boundary. Since the velocity of the solar wind is predominantly in the $\hat{x}$ direction, B_x is not advected into the simulation domain. To apply a B_x component, it needs to be set when

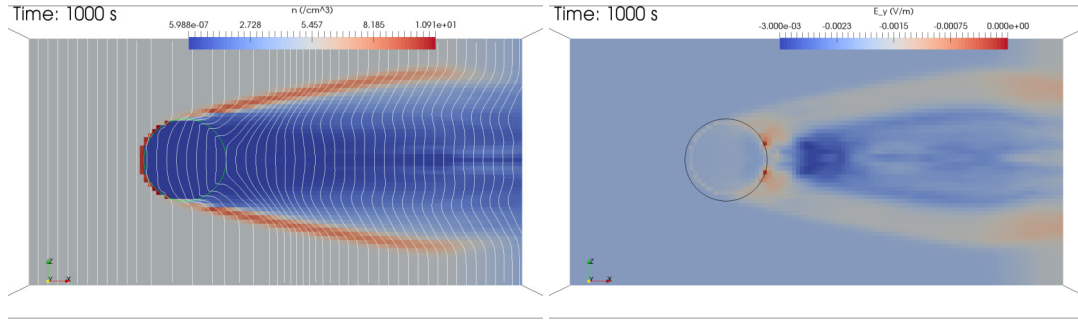


Figure 2.9: Simulation of the inner boundary without a magnetosphere, sliced through the centre of the planet along magnetic field lines. The left shows the number density and magnetic field streamlines (in white). The right shows the $\hat{y}$ component of the electric field. The circle denotes the size of the boundary. Solar wind is input from the left of the box. The vacuum region creates Alfvén wings on the upper and lower halves of the figure, shown as a decrease in the electric field.

initialising the simulation, as shown in figure 2.10, and is constant (in the solar wind) throughout the simulation, as is done in other GMHD codes (J. Raeder et al., 2001).

As explained in the introduction, in reality the magnetosphere transitions seamlessly into the ionosphere. This is not possible in GMHD codes, largely due to the dipolar magnetic field. The finite uniform grid size will not be able to resolve the magnetic field and currents, leading to increasing numerical errors. Furthermore, the high Alfvén speed due to the increasing magnetic field strength reduces the timestep to a prohibitively small amount. Due to these numerical constraints, the MHD equations are solved to an inner boundary, typically set at a radius $3 R_p$ to $4 R_p$ depending on the grid resolution.

Gorgon’s inner boundary consists of a spherical vacuum region containing a dipole magnetic field source. Any plasma which enters this region, is numerically removed. This allows the region to stay as a vacuum, meaning the magnetic field propagates using the electromagnetic wave solver. Within this region, a dipolar magnetic field is set within a smaller radius, typically $1 R_p$ to $2 R_p$. The magnetic field can propagate out from the planet toward the plasma region, where it switches back to normal MHD propagation. This allows the magnetic field to change on the inner boundary, creating a resistive sphere.

In order to verify this boundary’s resistivity, the simulation is run without a dipole in place, just a spherical vacuum region. In figure 2.9, it is seen that the inner boundary forms Alfvén wings, which are formed by a supersonic plasma flowing over a resistive body. As the solar

wind interacts with the vacuum region, the numerical resistivity breaks the frozen-in condition, allowing the plasma to diffuse through the magnetic field, or vice versa. This creates a kink in the magnetic field line, which propagates outwards as an Alfvén wave. Since the rest of the field is still frozen-in, the Alfvén wave is carried downstream by the solar wind, creating Alfvén wings. The process of Alfvén wing formation is described in Neubauer (1980), which relates the angle of the wing to the solar wind flow and the Alfvén Mach number, and the electric field in the Alfvén wing to the solar wind electric field and the object’s conductance. Hence, the conductance of Gorgon’s inner boundary can be estimated by sampling the electric field in the wing. It is found to be roughly $\sim 1 \Omega^{-1}$, depending on the grid size. This conductance is comparable to the ionospheric conductances usually used in GMHD simulation.

Whilst this approach does create a working magnetosphere, it is missing vital aspects of the ionosphere that cause a host of magnetospheric phenomena. For example, the Hall effect prevalent in the ionosphere is neglected. Finally, it does not aid space weather prediction of the ionosphere, as one cannot obtain a field-aligned current distribution over the ionosphere.

Initialisation

The standard procedure used throughout this thesis is to begin with an empty simulation box: all plasma is set to the numerical ‘vacuum density’. This allows the planetary dipole to expand into the simulation domain. Simultaneously, the solar wind is allowed to enter from the Sunward ($+x_{\text{GSE}}$ boundary). Initialisation typically takes two hours of simulation time. The simulation coordinate system is consistent with Geocentric Solar Ecliptic (GSE), where the planetary dipole is located at the origin.

The initialisation process is shown in figure 2.10. Here, a typical slow solar wind of $n_{\text{sw}} = 5/\text{cm}^3$, $\vec{v}_{\text{sw}} = -400\hat{x}$ km/s, $T_i = T_e = 5$ eV, and IMF of $\vec{B} = -5\hat{z}$ nT. The planetary dipole is aligned to the z axis. The initial condition is shown in the first panel, at $t = 0$ min. In this specific case, a background B_x is specified at the start of the simulation, as denoted by the horizontal black lines.

As time increases, the dipole expands out of the centre, shown by the dipolar field lines (black) near the red and white circles. The solar wind interacts with the planetary field, creating a cavity in the mass density, as is shown by the circular dark blue region impinging on the green solar wind at $t = 4$ min. This is the beginning of the magnetopause. At later times, a dark yellow

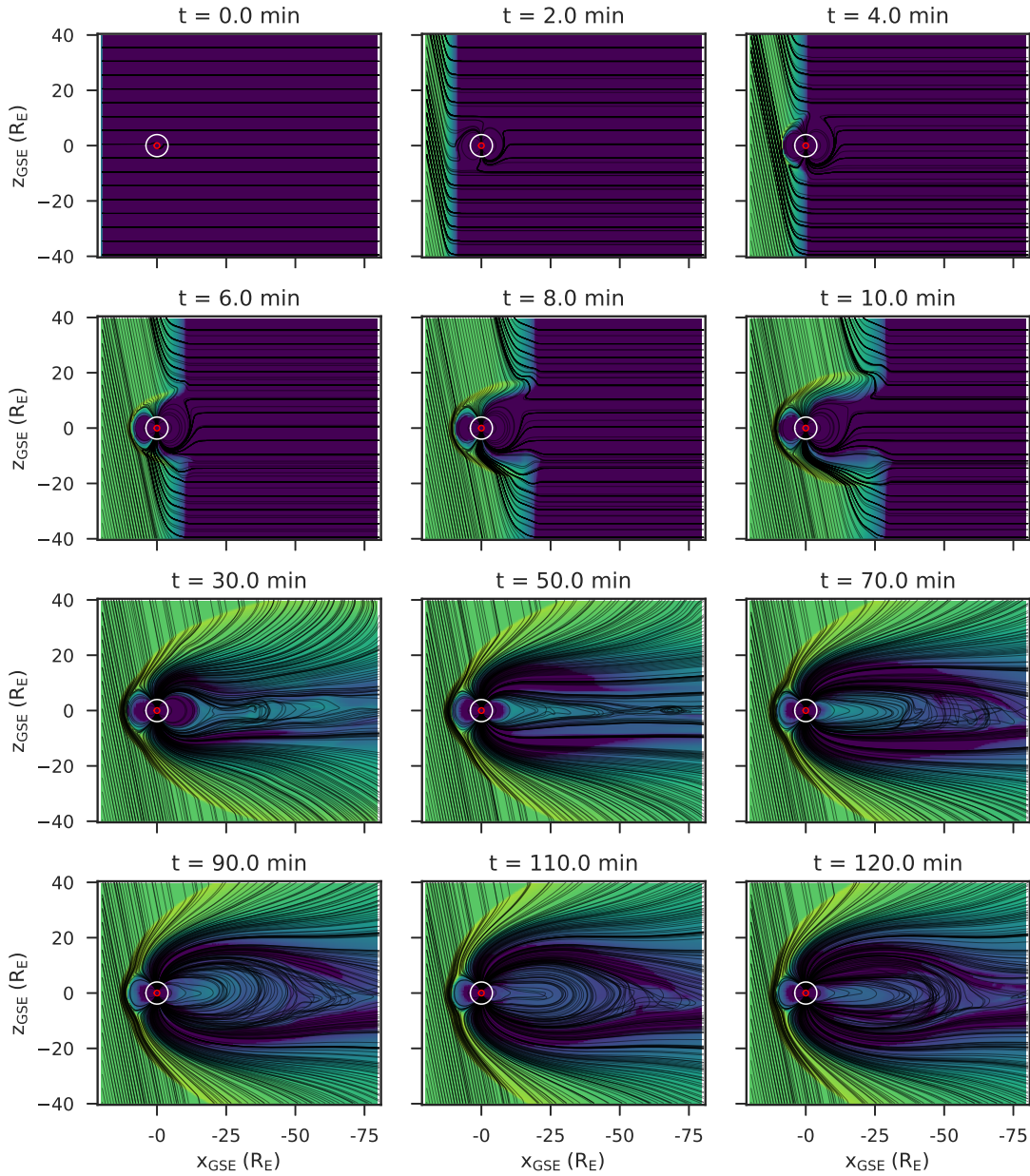


Figure 2.10: Initialisation of the magnetosphere. Each panel shows a snapshot in time of the simulation, sliced through the $x - z$ (noon-midnight) plane. The colour denotes log of mass density: blue is low, and green is high. The black lines are magnetic field lines. The two inner circles show the inner boundary (white) at $4 R_E$, and the Earth (red).

region of high mass density forms Sunward of the magnetopause. This is the magnetosheath, which is bounded by the bow shock further upstream.

At $t = 30$ min, the simulation already begins to resemble the full magnetosphere, with the magnetotail forming, containing the plasma sheet with elongated dipole field lines, and the low density magnetotail lobes. Once it reaches the two hour mark, the simulation has settled to a quasi-steady state.

2.5.2 Benchmarking Gorgon

In Gordeev et al. (2015), six key parameters are extracted from each GMHD code: the sub solar magnetopause position, high-latitude magnetopause position, plasma sheet pressure, lobe magnetic field strength; and two ionospheric parameters, total field-aligned current and the cross polar cap potential. These key parameters have been chosen as they can be seen as indicators of magnetospheric dynamics: for example, the lobe magnetic field strength is an indicator of the amount of magnetic flux in the tail (McPherron, 1991).

The values of these key parameters are obtained from a set of 19 six hour simulation runs with various different solar wind and IMF parameters. These parameters have been chosen to be statistically representative of the real solar wind, along with a few extreme events. In each run, a constant solar wind (e.g. constant ρ_{sw} , T_{sw}) is applied, but the IMF rotates and changes in magnitude every two hours. Each of these key parameters are then compared to empirical models for those specific solar wind and IMF parameters (e.g. the sub-solar magnetopause position is compared to the Shue et al. (1998) model).

This allows the determination of three skill scores for each GMHD simulation: the correlation coefficient, the slope of a linear fit and the prediction efficiency. The prediction efficiency is a measure of both the variation in the two datasets and the difference in absolute value. It is given by,

$$PE = 1 - \frac{\langle (X_G - X_M)^2 \rangle}{\sigma_G^2} \quad (2.21)$$

where the angle brackets $\langle \dots \rangle$ denote the arithmetic mean, X_G is the GMHD value, X_M the empirical model value, and σ_G^2 the variance in the GMHD values. A “perfect” score would be a correlation coefficient, slope and prediction efficiency all of 1.

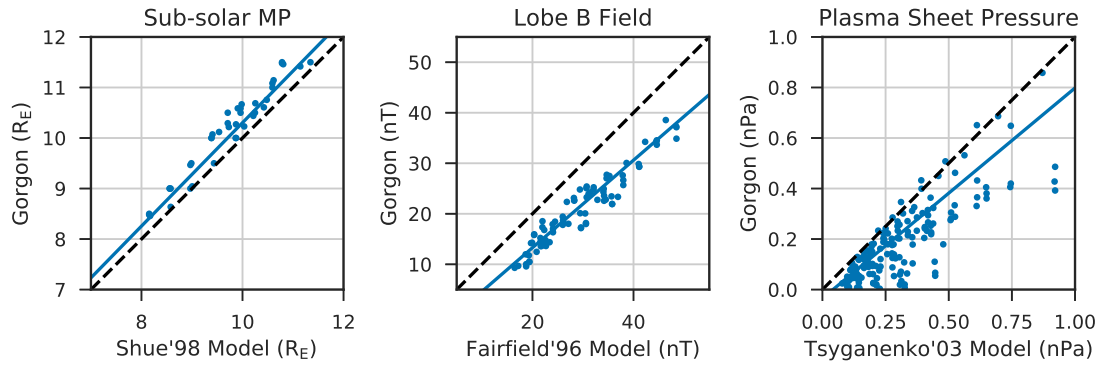


Figure 2.11: Benchmarking the Gorgon MHD code against empirical models, comparing the sub-solar magnetopause distance, lobe magnetic field strength, and the plasma sheet pressure. Each panel shows the empirical model plotted against the Gorgon value, with the blue dots showing the 2 hour average of the 5 min cadence dataset. The linear fit, shown in blue, is performed on the 5 min dataset. The black dashed line denotes the one-to-one, “perfect” relationship between Gorgon and empirical model.

This methodology is applied to the Gorgon code, where the sub-solar magnetopause position, lobe magnetic field strength and plasma sheet pressure is analysed. To obtain the sub-solar magnetopause distance, the magnetopause current sheet is identified by a peak in the current density. The lobe magnetic field strength are sampled at two locations in the northern lobe: $(-15, 0, 8) R_E$ and $(-25, 0, 8) R_E$. The plasma sheet pressure is sampled along four points, positioned at $x = (-12, -15, -20 \text{ and } -25) R_E$ and $y = 0$. The plasma sheet pressure is then chosen to be the largest thermal pressure in the magnetotail, along the z axis. Each parameter is sampled at a 5 min cadence, except when the IMF turns southward, where the cadence increases to 1 min to capture potential substorm effects (Gordeev et al., 2015).

Figure 2.11 shows the relationship between the Gorgon value and the empirical model: Shue et al. (1998) for magnetopause distance, D. H. Fairfield and Jones (1996) for the lobe magnetic field strength, and Tsyganenko (2003) for the plasma sheet pressure. The x axis shows the empirical model with the y axis the Gorgon value. The one-to-one relationship between model and Gorgon is shown as the black dashed line. Ideally, each of the two hour averages (blue dots) and linear fit of the five minute dataset (blue line) should lie exactly on the black dashed line. The sub-solar magnetopause distance performs particularly well. The linear fit scales well, with a fit slope of ~ 1 , being approximately parallel to the one-to-one relationship. The two hour averages show little spread from the fit. In general, Gorgon tends to overestimate the magnetopause position, with the fit line and all data points lying above the one-to-one

relationship. This could be due to the $0.5 R_E$ grid size: the discretisation of the magnetic field creates non-zero current, even in a current free dipole field. This adds to the pressure balance in the magnetosphere, pushing the magnetopause further out.

The lobe magnetic field strength is considerably smaller than the one-to-one relationship, with the fit and points lower than the black dashed line. It scales relatively well however, with a fit slope of approximately 0.8. The discrepancy here is likely due to a lack of ionosphere in the simulation, which affects the amount of open flux in the magnetotail (Kamide et al., 1998), and hence the lobe magnetic field strength.

The plasma sheet pressure shows the most spread in two hour data points, though Gorgon always underestimates the plasma sheet pressure. Despite this, the slope of the fit scales relatively well, with a slope approximately 0.8.

The skill scores for Gorgon are shown in figure 2.12 (blue), compared to those of the four GMHD codes tested in Gordeev et al. (2015): BATS-R-US (green), GUMICS-4 (orange), OpenGGCM (pink), and LFM (yellow). Gorgon performs well in the slope and correlation coefficient, scoring above 0.8 in all three key parameters. The prediction efficiency has more mixed results. Out of the three parameters, Gorgon performs best in prediction the sub-solar magnetopause position, with slope, correlation coefficient and prediction efficiency near 1. Whilst Gorgon performs well in the slope and correlation coefficient in the plasma sheet pressure, its prediction efficiency is considerably lower. It performs the worst in the lobe magnetic field strength, with a negative prediction efficiency. The low values of the prediction efficiency are due to that score taking into account difference the absolute value of the parameter: Gorgon systematically has lower plasma sheet pressures and lobe magnetic field strengths than predicted by the other models.

When comparing Gorgon between the other codes, it scores the same, if not higher in all three skill scores. The exception to this is the lobe field strength prediction efficiency, which is considerably lower than the other codes. As previously mentioned, this could be due to a lack of ionosphere.

2.6 Summary

The main aims of this chapter is twofold: to explain the numerical methods Gorgon utilises to solve the MHD equations and verify its performance when subjected to simplified test problems;

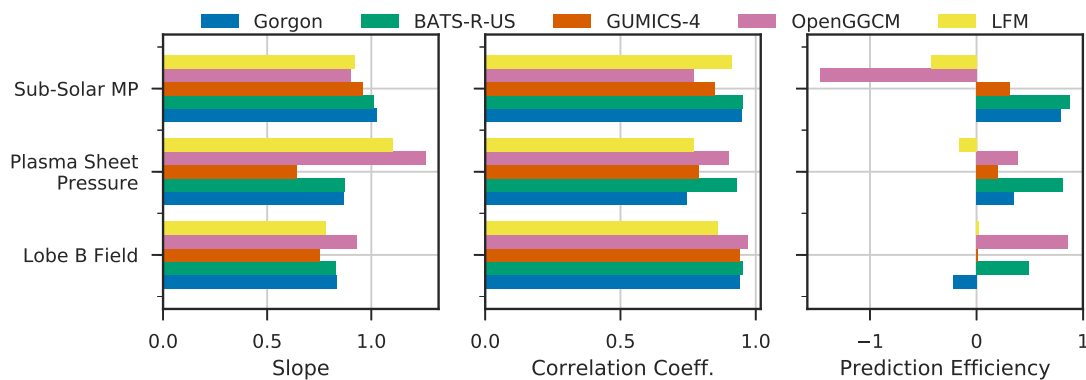


Figure 2.12: Skill scores, fit slope, correlation coefficient and prediction efficiency for the three parameters tested, sub-solar magnetopause position, lobe magnetic field strength and plasma sheet pressure. The blue is Gorgon, from the data shown in figure 2.11. The rest of the data are from the other four GMHD codes, from Gordeev et al. (2015). A perfect score for all three scores is 1.

to compare to existing MHD codes to Gorgon; and to explain its magnetospheric implementation and how well it performs in a magnetospheric context.

By subjecting Gorgon to standard MHD test problems, it has shown that it does solve initial condition problems correctly, and comparably well to other MHD codes. Gorgon was subjected to three widely used standard MHD test problems: small amplitude (linear) waves, shock tubes, and the Orszag-Tang vortex. The linear waves test shows that Gorgon can simulate small amplitude waves, which include the three MHD wave modes (magnetosonic slow and fast, and the Alfvén wave). In order to test the non-linear regime, Gorgon simulated Sod’s (hydrodynamic) and Brio-Wu’s (magnetohydrodynamic) shock tubes, showing that Gorgon handles shocks well. Finally, the Orszag-Tang vortex tests the code’s ability to handle both shocks and waves simultaneously using a complex two-dimensional initial condition. The solution is compared to a different MHD code (Londrillo and Del Zanna, 2000), which agrees with Gorgon. These tests gives confidence in the fact that Gorgon can correctly solve the MHD equations, and be used in more complex problems.

The Gorgon code has also been subjected to benchmarking in the magnetospheric context. This was done using the method described in Gordeev et al. (2015), which compares key magnetospheric parameters to empirical models. Here, it performed just as well as other global magnetosphere codes, scoring especially high in the skill scores of the sub-solar magnetopause (figure 2.12, top bar chart). Compared to the three parameters tested, Gorgon performs best

when predicting the sub-solar magnetopause position, and reasonably well in the plasma sheet pressure. This suggests that Gorgon can correctly simulate the outer boundaries of the magnetosphere, such as the bow shock, magnetosheath and magnetopause. However, it did not as perform well in the lobe magnetic field strength, implying some discrepancy on the inner boundary. This is expected, as MHD is not expected to perform particularly well in the inner magnetosphere (De Zeeuw, 2004). Furthermore, Gorgon is not coupled to any ionosphere or ring current models, which are expected to play a role in improving this metric (Gordeev et al., 2015).

3 Highly Inclined Rotating Dipoles

The work in this chapter has been published in *Journal of Geophysical Research: Space Physics* in Mejnertsen et al. (2016).

3.1 Introduction

Study of magnetospheres is not only constrained to Earth. Five out of the eight planets in the solar system have appreciable internal magnetic fields which extend into interplanetary space, creating a magnetosphere (Bagenal, 1992). Each planet has a different combination of solar wind parameters, dipole and rotation axis tilts, internal plasma sources (such as moons and the ionosphere), which makes modelling each planet interesting and challenging.

The outer planet magnetospheres, Jupiter, Saturn, Uranus and Neptune, tend to have much larger magnetospheres relative to their planet size, compared to the inner planet magnetospheres, Earth and Mercury. This is shown in figure 3.1, where Mercury's magnetosphere can easily fit into Earth's magnetosphere, Earth's inside Saturn's, and Saturn's inside Jupiter's magnetosphere. This is prominently due to two factors: the solar wind dynamic pressure decreases with distance from the Sun (Slavin and Holzer, 1981), and their strong internal magnetic field strength, although the presence of internal drivers can add to overall size of the magnetosphere as is the case with Jupiter (Bagenal, 1992).

The drivers of dynamics and variability in these magnetospheres can be roughly divided into four categories: solar wind variability, internal plasma sources, the rotation and tilt of the planetary magnetic field, and the ionosphere.

In the case of Earth and Mercury, the dynamics are mostly driven by the solar wind due to their proximity to the Sun (e.g. Brain et al. (2010) for Mercury). They lack internal drivers, such as

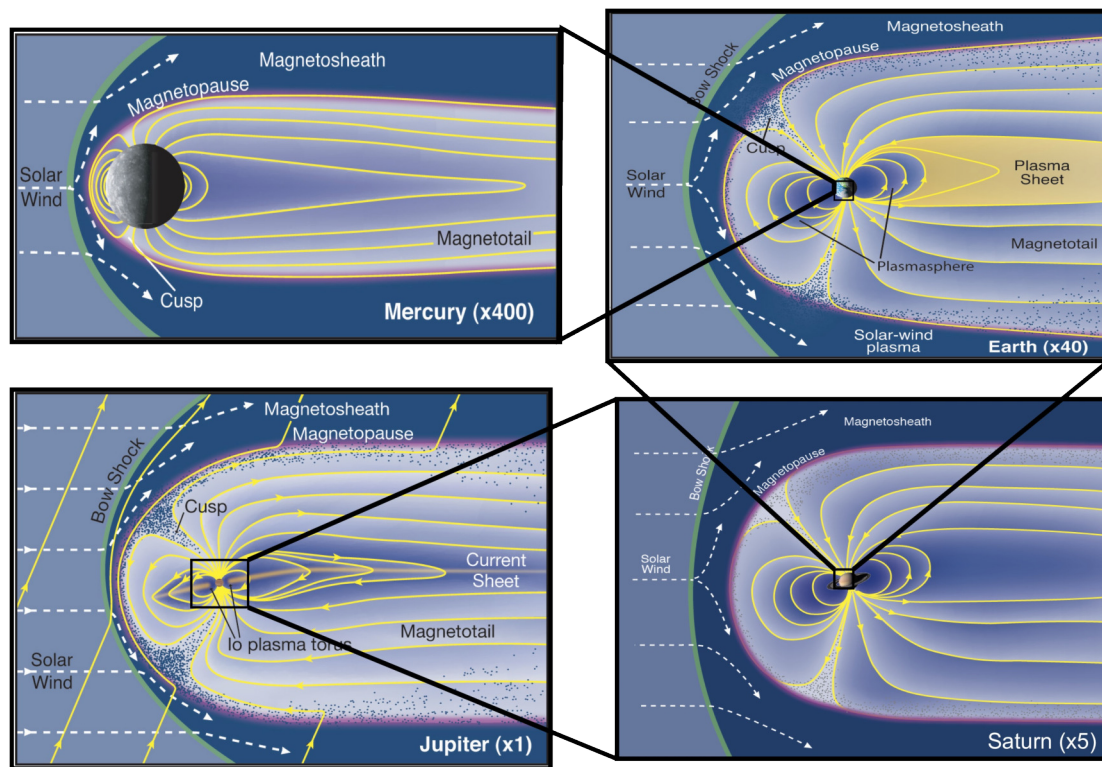


Figure 3.1: Comparison of Mercury's, Earth's, Saturn's and Jupiter's magnetospheres in increasing order of size. Credit: Fran Bagenal and Steve Bartlett¹.

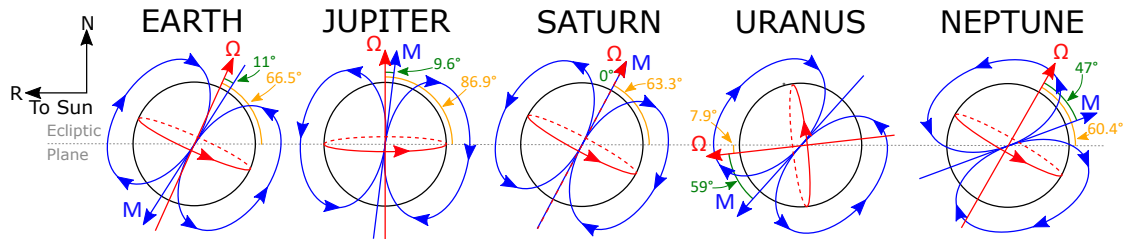


Figure 3.2: Rotation and dipole axes of the magnetised planets. Earth, Jupiter and Saturn all have approximately parallel dipole and rotation axes, which are almost perpendicular to the ecliptic plane (direction into the page). Uranus and Neptune have large angles between the rotation and dipole axes, as well as large angles between the rotation axis and the ecliptic plane. Adapted from Bagenal (1992).

active moons, and their rotation axis is approximately perpendicular to the solar wind flow, and approximately parallel to the magnetic dipole axis.

Jupiter and Saturn on the other hand are much further away from the Sun, and as a consequence the solar wind driving is much weaker. They do however have a number of active moons which drive the magnetosphere internally (e.g. M. G. Kivelson et al. (1997), Krupp et al. (2004), and Mitchell et al. (2009)). Due to this, the effect of solar wind driving on these magnetospheres much less understood (e.g. Cray et al. (2005) and Clarke et al. (2009)).

Finally, the Ice Giants, Uranus and Neptune, are the furthest magnetised planets from the Sun. Whilst they do have moons, their magnetospheric impact is relatively unknown (C. S. Arridge et al., 2014; Masters et al., 2014). However, their dipole and rotations are significantly tilted with respect to each other, and the solar wind flow (Connerney, 1993). This is illustrated in figure 3.2. This creates a magnetosphere strongly driven by this daily variation, which changes with season (M. Schulz et al., 1995; Masters, 2014; Masters, 2015; Desch et al., 1991).

To better understand the general physics of magnetospheres formed by highly-inclined, rapidly-precessing dipoles, Gorgon has been adapted to simulate Neptune’s magnetosphere, using a precessing dipole, as observed during the Voyager 2 flyby.

3.1.1 The Ice Giant Magnetospheres

Within the solar system, the planets Neptune and Uranus are examples of ‘ice giants’, where an important component of their internal composition are volatile ices such as water, ammonia and methane (e.g. Guillot (2005), Christopher S. Arridge et al. (2012), and Masters et al. (2014)). Both

planets are very important for a whole host of solar system, planetary and heliophysics science goals, and test our theoretical understanding of planetary formation (e.g. Tsiganis et al. (2005)), dynamo physics (e.g. Soderlund et al. (2013)), and magnetospheric dynamics (Bagenal (1992)). Due to their similar internal structure, both Uranus and Neptune's magnetospheres are very similar (C. S. Arridge et al., 2014; Masters et al., 2014). The rotation axes of both planets are tilted relative to the normal of their orbital plane ($\sim 97.5^\circ$ at Uranus and $\sim 23^\circ$ at Neptune) (Bagenal, 1992; Holme and Bloxham, 1996). Both planets' magnetic fields also exhibit considerable higher order moments compared to the other giant planets (C. T. Russell and Dougherty, 2010). Such large magnetic moment tilts are thought to be generated by a complex dynamo within each planet which is not fully understood. It is thought both Uranus and Neptune are composed of a mixture of ices and rocks (Connerney, 1993; Soderlund et al., 2013), which causes the formation of complex, high order, non-axisymmetric magnetic fields (Ness et al., 1989; Connerney, Acuña, and Ness, 1991; Holme and Bloxham, 1996).

The primary source of experimental observations concerning the ice giants comes from Voyager 2, which is the only spacecraft thus far to visit Uranus and Neptune in 1986 and 1989 respectively (E. C. Stone and Miner, 1986; E. C. Stone and Miner, 1989). These flybys provided the bulk of the evidence that the ice giants are fundamentally different from Jupiter and Saturn (Kurth and Gurnett, 1991; Connerney, 1993; B. H. Mauk and Fox, 2010). In addition to their composition, both exhibit magnetic moments that are significantly tilted when compared to the planetary rotation axis ($\sim 60^\circ$ at Uranus and $\sim 47^\circ$ at Neptune).

Whilst they are thought to have similar origins and internal structure, there are some notable differences. The specific set of rotation axis and magnetic dipole axis tilts create magnetospheres which vary differently to one another on a daily and seasonal basis, potentially creating differences in how effective solar wind coupling is and the movement of plasma with the magnetosphere.

On its trajectory out of the solar system, Voyager 2 encountered Uranus first (E. C. Stone and Miner, 1986), orbiting at a radius of 19.2 AU. At the time of the encounter, Uranus's rotation axis was approximately in the Sunward direction (E. C. Stone and Miner, 1986). With the magnetic dipole axis tilted by 60° degrees with respect to the rotation axis, creating a daily precession of the magnetosphere about the Sun-Uranus line. This creates a magnetosphere which can be approximated to a steady state within the reference frame rotating with the magnetic dipole (Tóth et al., 2004), and strongly modulates reconnection (Masters, 2014).

On approach to Neptune, Voyager 2's trajectory took it first through the bow shock, then exiting the magnetosheath into the magnetosphere through the dayside cusp region near the sub-solar point (Ness et al., 1989; Szabo et al., 1991). At this time the magnetic moment was found to be approximately parallel to the solar wind flow. Voyager 2 remained inside the magnetosphere for more than 38 hours and thus more than 2 Neptunian days. Compared to the other outer planets, Neptune's magnetosphere was found to be relatively empty (Belcher et al., 1989), although Triton, with its dense atmosphere (Broadfoot et al., 1989), may act as a source (John D. Richardson et al., 1991; Zhang, John D. Richardson, and Sittler, 1991). However, the role of Triton as a plasma source in Neptune's magnetosphere is ultimately not yet known (Masters et al., 2014).

On the basis of the measurements made by Voyager 2, it was concluded that the magnetosphere goes back and forth from an Earth-like to a pole-on configuration, as shown in figure 3.3. With the magnetic axis almost perpendicular to the solar wind flow, the Earth-like configuration is thought to resemble the Earth's magnetosphere, with the standard dayside plasma region, cusps, and tail containing a plasma sheet. In the pole-on configuration, the magnetic axis is almost parallel to the solar wind direction, causing a rarefied cusp region on the dayside, and northern and southern current sheets separated by a rarefied cusp region in the tail (Ness et al., 1989; Belcher et al., 1989; R S Selesnick, 1990; R. Lepping, 1994; M. Schulz et al., 1995). Although they appear to be two separate current sheets in the 2D slice shown in figure 3.3, it is thought they may be connected, forming a cylindrical structure (Voigt and Ness, 1990; Michael Schulz and McNab, 1996).

Due to its distance ~ 30 AU from the Sun, on approach to Neptune Voyager 2 found both the solar wind density and magnetic field strength to be low, with average values of $6.78 \times 10^{-3} \text{ cm}^{-3}$ and 0.18 nT respectively (C. S. Arridge, 2015). The solar wind speed was largely similar to values observed in the inner heliosphere, $v_{\text{avg}} \sim 461 \text{ km s}^{-1}$ (C. S. Arridge, 2015). This, in combination with Neptune's relatively large magnetic dipole moment ($2.2 \times 10^{17} \text{ Tm}^{-3}$, 25 times greater than that of Earth), means that the magnetosphere is relatively large and that the rotation period cannot be ignored in the context of the flyby. Its rotation leads to two distinct configurations of the magnetosphere, which alternate every half rotation (~ 8 hours) (Belcher et al., 1989; R S Selesnick, 1990; Bagenal, 1992).

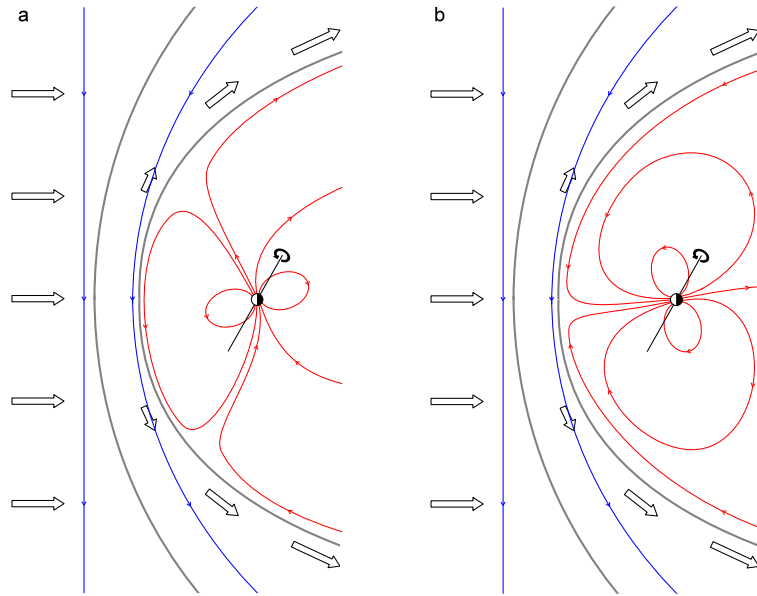


Figure 3.3: The two different configurations of Neptune's magnetosphere at solstice: Earth-like (left) and pole-on (right). The plane of this image contains Sun-Neptune line, and the normal to the solar ecliptic plane. The arrows denote the solar wind flow, with the blue lines showing the IMF. The red lines show the magnetic field of Neptune. The two black lines are, from left to right, the bow shock and magnetopause. Figure from Masters (2015).

3.1.2 Previous Modelling

Following the Voyager 2 observations, several attempts have since been made to model Neptune's magnetosphere analytically and to predict aspects of its dynamics.

For example, by using a simple model for the convection electric field (R. S. Selesnick and J. D. Richardson, 1986) and calculating the convection velocity inside the corotating frame, R S Selesnick (1990) found that the rotation could cause a particle accelerator effect, which may transport particles either toward or away from the planet.

Regarding the energy input into Neptune's magnetosphere, magnetopause reconnection is thought to be significant (Desch et al., 1991), but highly dependent on season, time of Neptune day, and interplanetary magnetic field (IMF) orientation (Masters, 2015), and is expected to be 'bursty' and infrequent (Huddleston et al., 1997). Subsequent processes within the magnetosphere, such as tail reconnection, may suffer complications due to winding of magnetic field lines in a similar mechanism as proposed by Cowley (2013) for Uranus, who argued a seasonal dependence on the ability to close open flux in the Uranian system. This is shown in figure 3.4.

Previous attempts at modelling Neptune's magnetosphere have been performed by Voigt and Ness (1990) and M. Schulz et al. (1995). Voigt and Ness (1990) used a two-dimensional magneto-hydrodynamic (MHD) equilibrium magnetosphere to explore the energy balance within the magnetosphere. They used an empirical model of the magnetosphere (Voigt, 1981) to verify the accuracy of their MHD model. By finding linear solutions to the Grad-Shavranov equation, they suggest the free energy within Neptune's magnetosphere remains constant throughout Neptune's daily rotation. M. Schulz et al. (1995) modelled Neptune's magnetic topography using the source surface model (Michael Schulz and McNab, 1996) for different angles of attack (Ψ), i.e. the largest angle between the solar wind velocity and dipole moment. Both models (Voigt, 1981; M. Schulz et al., 1995) predict for large Ψ , the magnetotail undergoes a distinct change in topology: the tail current sheet no longer connects to the magnetopause current sheet, but closes in on itself, forming a cylinder. However, neither model is intrinsically time dependent. This is shown more specifically in M. Schulz et al. (1995), who note that it is not fully self consistent as their model does not achieve thermodynamic equilibrium.

More recently, global three-dimensional MHD simulations have been performed on Uranus' magnetosphere. Tóth et al. (2004) used an adapted version of BATS-R-US to simulate Uranus during solstice, making use of Uranus' rotation axis pointing Sunward to cast the MHD equations in a

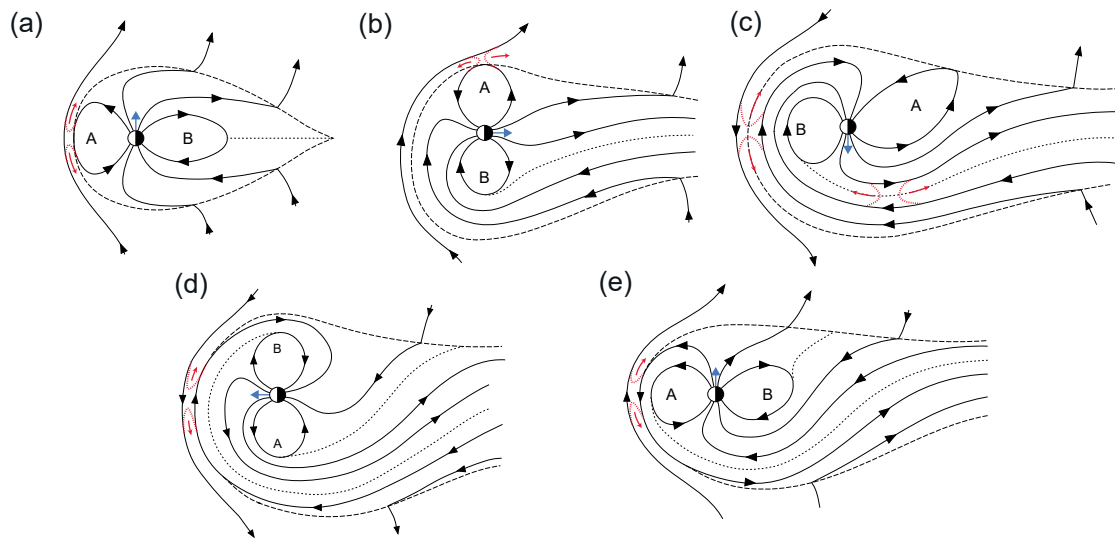


Figure 3.4: The effect of a precessing dipole inhibiting a ‘Dungey Cycle’-like process at Uranus during equinox. Panels (a) through (e) are sequential. In panels (a) through (b), the IMF reconnects with region A, opening flux on the dayside. As Uranus rotates, region B moves to the dayside, and A into the tail (panel c). The flux opened in panel (a) now drapes from the southern lobe, over the closed field line region B, and into the north pole, inhibiting the opening of the flux in region B. Figure from Cowley (2013).

rotating coordinate system. In this rotating frame, Uranus' magnetosphere is well approximated by a steady state solution. They focus on the twisting of the magnetotail due to the precession of the internal magnetic field, calculating a twisting rate of $\sim 0.4^\circ/R_U$ ($R_U = 25\,559$ km, Uranian planetary radius). Cao and Paty (2017) studies the coupling between the solar wind and the Uranian magnetosphere at different times of year. They find a 'switch-like magnetosphere' occurs both at equinox and solstice, where the rotation of the planet moves the magnetosphere from open to closed, showing a strong dependence on reconnection.

3.1.3 Open Questions and Motivation

Whilst the Voyager 2 mission answered many questions about the Ice Giant magnetospheres, it also raised many more. There are open questions as to how the magnetosphere responds to the changing solar wind and planetary magnetic field, and whether processes such as reconnection can create global circulation of flux within a precessing magnetosphere.

The study of both Neptune and Uranus may provide insights for other planets in the solar system, where the dipole tilt is not as extreme. Furthermore, a significant number of exoplanets are thought to be Neptune-like (Batalha et al., 2013), potentially providing insights there too.

Although there have been science cases put forward for new missions to visit both Uranus (C. S. Arridge et al., 2014) and Neptune (Agnor et al., 2009; Hansen, Hammel, and Spilker, 2009; Christophe et al., 2012; Masters, 2014), in-situ observational data for the foreseeable future remains limited to the Voyager 2 data set. In order to further the understanding of Neptune's magnetosphere, modelling and simulations must be developed, which would provide predictions for future missions.

The initial aim of this chapter is to understand the structure of the magnetosphere of a highly-inclined precessing dipole, specifically: the change in shape and location of the outer magnetospheric boundaries on the dayside; the possible location and variation of magnetic reconnection on the magnetopause; and the changes in structure of the magnetotail as a function of planetary rotation.

Table 3.1: Solar wind conditions at Neptune, comparing the values observed by Voyager (used in simulation), and values predicted by scaling laws. ¹From Slavin and Holzer (1981) and Masters (2015).

	n_P (/cm ³)	T (eV)	v_{sw} (km s ⁻¹)	B (nT)
Voyager 2	4.6×10^{-3}	0.5	$-400\hat{x}$	$0.07\hat{y} + 0.11\hat{z}$
Predicted from Scaling ¹	2×10^{-2}	2.5	450	0.11

3.2 Adapting the Gorgon code for Neptune

A region of $-110 < x_{NSO} < 40$, $-60 < y_{NSO} < 60$, $-60 < z_{NSO} < 60$ (distances in Neptune radii, $R_N = 24.6 \times 10^3$ km) is simulated. The simulation axes are equivalent to the Neptune Solar Orbital (NSO) coordinate system, with the x direction Sunward, y opposite to the direction of Neptune's orbital motion, and z completes the right-handed set. This is the Neptune equivalent of the Geocentric Solar Ecliptic (GSE) coordinate system commonly used at Earth. To clarify the terminology used in this chapter, north is defined to be in the $+z$ direction (i.e. northern regions are where $z > 0$), and south is in the $-z$ direction. Dusk is defined to be in the $+y$ direction, and dawn in the $-y$ direction. Simulations were performed with a resolution of $0.3 R_N$.

The solar wind conditions are kept steady throughout the simulation. They are based on the observations made by Voyager 2 just prior to crossing the bow shock, as shown in the 'Voyager 2' row in table 3.1 (Szabo and Ronald P. Lepping, 1995).

It is worth noting that these solar wind conditions are approximately 75% lower (with the exception of the speed) than the average expected value (Slavin and Holzer, 1981) (as shown in the bottom row of table 3.1) due to the variable nature of the solar wind. Hence the size and shape of the magnetosphere is expected to differ from typical Neptunian magnetosphere (Masters, 2015). Furthermore, running the simulation with a steady solar wind does not capture such variability. In order to compare with Voyager 2 observations, the lower than average solar wind parameters are used.

The simulation is initialized with the same method as in section 2.5.1. The vacuum density is set to $\rho_{vac} = 10^{-26}$ kg/m³, to account for the lower densities present in the Neptunian system. The precessing dipole source begins at northern midnight: angled at co-latitude (from $+z$ axis) of 18.7° and azimuth (from $+x$ axis) of 16.1° . It rotates about an axis angled at a co-latitude of 28.3° and azimuth of 16.1° , which corresponds to the orientation of Neptune's rotation axis at

the season Voyager 2 arrived. The solar wind is input from the start, filling the empty box and interacting with the precessing dipole to first form the magnetosheath and magnetopause. It runs for a total of 307,500 s (5.3 Neptune days).

Since the properties of Neptune's ionosphere and the impact of Triton on the Neptune's magnetosphere are very poorly understood and highly uncertain, and since the focus of our investigation is not on the inner magnetosphere, a Triton-like plasma source has not been included into the simulation, and the ionosphere is treated as a resistive shell.

3.2.1 Modelling Neptune's Magnetic Field

The significant non-alignment of Neptune's and Uranus' dipole and rotational axes presents a different challenge to modelling their magnetospheres. In this regard, Uranus is more simple because the rotational axis is approximately contained within the planet's orbital plane.

At the time of the Voyager 2 flyby of Uranus, the rotational axis was approximately anti-parallel to the solar wind flow (E. C. Stone and Miner, 1986). The consequence of this is that if one changes into a frame rotating with the dipole, the solar wind flow is unaffected, but the IMF rotates around the x axis. This symmetry was exploited in the global simulations of Uranus' magnetosphere performed by Tóth et al. (2004). In this work, they accounted for the rotation of the magnetic axis by utilizing the rotational symmetry of Uranus' magnetosphere at the time of Voyager 2's flyby, recasting the MHD equations to include rotational effects, effectively solving the MHD equations in a corotating frame. In the case of Neptune, this approach cannot be used because as discussed in the introduction, the rotational axis does not lie in the Neptune orbital plane, being offset by $\sim 23^\circ$. The simplest approach to modelling Neptune's interaction with the solar wind requires a precessing dipole (i.e. whose axial orientation changes with time) to be included specifically in the code.

This is implemented by rigidly rotating the magnetic dipole that is set within the vacuum region inside the inner plasma boundary. The rotation is given by,

$$\frac{d\vec{M}}{dt} = \vec{\Omega} \times \vec{M}, \quad (3.1)$$

where $\vec{M}$ is the dipole moment and $\vec{\Omega}$ is the rotation axis, where $|\Omega| = 2\pi/T$ with rotation period, T . Throughout the simulation, $\vec{\Omega}$ is kept constant.

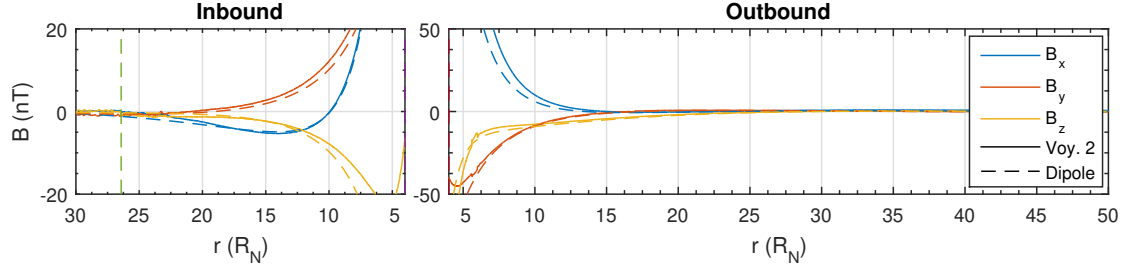


Figure 3.5: Comparing the Voyager 2 magnetic field data with the precessing dipole approximation, used at the inner boundary. Both on the inbound and outbound trajectories, the dipole approximation agrees well with the Voyager 2 data at distances $10 R_N$ to $20 R_N$. It begins to diverge near the planet, where higher order terms dominate, and near the dayside magnetopause, where the magnetic field is not expected to be dipolar.

Though Neptune’s magnetic field is distinctly non-dipolar near Neptune, the field becomes roughly dipolar at distances greater than $4 R_N$ (Ness et al., 1989; Connerney, Acuña, and Ness, 1991). This work is only concerned with the global structure of the magnetosphere and the behaviour of the outer magnetospheric boundaries: hence, the planetary field is approximated as a planet-centred dipole, with strength $2.2 \times 10^{17} \text{ T m}^3$ ($0.14 \text{ G } R_N^3$) (Connerney, Acuña, and Ness, 1991). This follows the approach of previous theoretical studies (Michael Schulz and McNab, 1996; Voigt and Ness, 1990). Although the Voyager 2 observations indicate that this dipole is offset by $0.5 R_N$ from the centre of the planet, centring the dipole simplifies the model but does not impact on the basic features of the simulated magnetosphere. This is shown in figure 3.5, comparing the dipole approximation used as the inner boundary of the simulation to the magnetic field observed by Voyager 2 (Connerney, Acuña, and Ness, 1991).

3.3 Overview of Neptune’s Magnetosphere

In this section, the basic features and structure of the magnetosphere are discussed, focussing on the overall configuration, the outer-magnetospheric boundaries, magnetopause reconnection, and tail dynamics. Figure 3.6 shows mass density slices in the $y = 0$ plane and $z = 0$ plane every half rotation from 2.5 days into the simulation. By day 4 to day 5, the simulation has reached a quasi-periodic state: the output generated from this interval is analysed to avoid any phenomena associated with the initialization of the simulation.

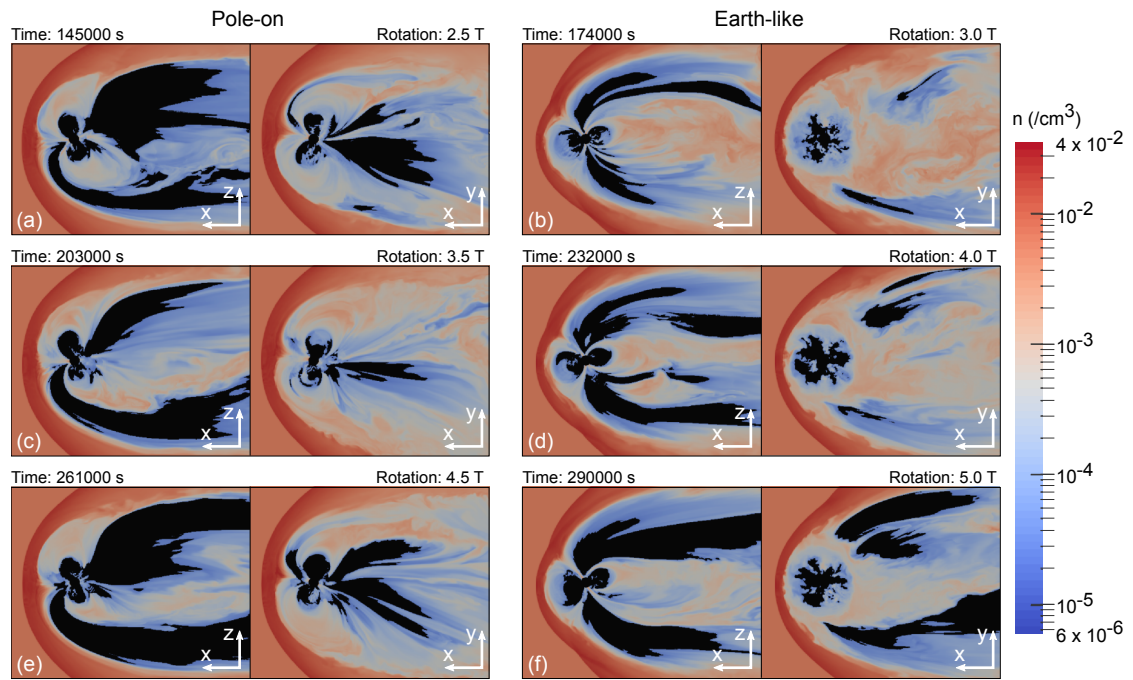


Figure 3.6: Slices of mass density at $y = 0$ and $z = 0$ (left and right of each subplot respectively) of every half rotation from 2.5 to 5.0 in subplots (a) to (f). The black regions denote where the density is at the floor density. This shows that the simulation has reached a quasi steady state.

Figure 3.6 already shows that Neptune’s magnetosphere undergoes major changes on a daily basis, from the pole-on configuration on the left, to the Earth-like on the right.

3.4 The Outer Boundaries

3.4.1 Identifying the Outer Boundaries

The location of the bow shock in the simulation can be obtained by locating the point where the plasma is compressed, indicating an increase in mass density. By scanning for this increase along the x direction for each value of y or z , 2D slices of the bow shock position are obtained, as shown in figure 3.9.

There are a number of methods to locate the magnetopause (Gordeev et al., 2015). Common ones include identifying the magnetopause current layer — seeking the peak in current density — or identifying the significant drop in mass density over the magnetopause surface. These methods are commonly used on the dayside, where the compressed nature of the magnetosphere creates high peaks in current density and sharp cut-offs in the mass density. In the tail, these methods become more problematic. The magnetopause current peak tends to be smeared over multiple grid cells, becoming smaller in magnitude, and the drop in mass density is not as pronounced.

Another method involves finding the fluopause: the surface where there is minimal solar wind entry into the magnetosphere. Palmroth et al. (2003) uses momentum streamlines to find the fluopause, calculating streamlines and searching for ‘voids’ in the position of these streamlines. These ‘voids’ are indicative of two different regions of plasma: magnetospheric plasma and magnetosheath plasma. By locating the ‘voids’ and binning them in a cylindrical grid, it gives a full three-dimensional shape of the magnetopause.

The method used here is similar to Palmroth et al. (2003). Momentum streamlines are generated upstream of the bow shock, and are calculated until they either leave the box or the inner boundary. This is done in the standard technique, by solving the streamline equation numerically using the momentum vector field,

$$\frac{d\vec{x}_s}{ds} = \frac{\vec{p}}{|\vec{p}|} \quad (3.2)$$

where s the path of the streamline, whose position is given by $\vec{x}_s$. The full path is found by integrating over ds from a given starting position (seed) using the fourth order Runge-Kutta

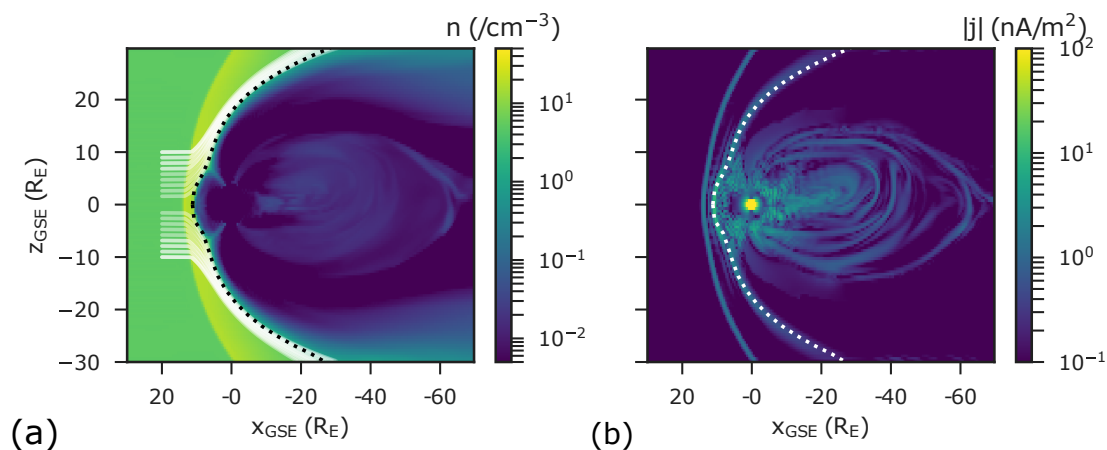


Figure 3.7: Three-dimensional magnetopause finding algorithm. Both figures show 2D cuts in the $x - z$ plane. (a) shows some of the streamlines (white lines) created to find the magnetopause position (black dashed line), plotted over a slice of number density. Momentum streamlines are calculated from positions upstream of the shock, and follow the momentum vector field until they reach the simulation boundaries. (b) shows the magnetopause position (white dashed lines) and the current density magnitude below it.

(RK4) integration scheme. At each position along the path, the value of the vector field is interpolated using tri-linear interpolation, the three-dimensional analogue to one dimensional linear interpolation. Typically, a step size is 1/10 of the grid cell size is used.

As the streamlines meet the magnetopause, they are deflected around the magnetopause surface. Figure 3.7 shows the streamlines generated in the solar wind (shown as the white lines in figure 3.7a) which deflect around the magnetopause surface. This defines the magnetopause surface.

In order to generate a surface from this, the streamline points are binned into a parabolic grid. Illustrated in figure 3.8, the parabolic grid has advantages over a cylindrical grid, as it is able to capture the radial nature of the magnetopause on the dayside, and the cylindrical nature of the magnetotail.

Each streamline is binned into a predefined grid of values of u and θ . The third coordinate, v , is determined by taking the minimum value in v for each of the streamlines in that particular bin. This creates a three-dimensional magnetopause surface. A slice in the $x - z$ plane is shown in figure 3.7b, plotted over a pseudo-colour plot of current density. The magnetopause position, indicated by the white dashed line, fits well with the magnetopause current sheet position, and

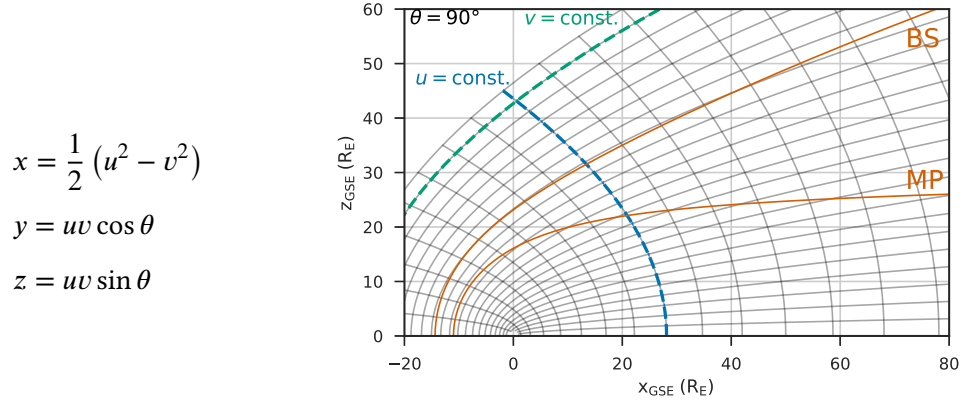


Figure 3.8: Illustration of the parabolic grid used in the magnetopause finding algorithm. It can capture the radial nature of the magnetopause (denoted by the orange line named ‘MP’), on the dayside, whilst becoming more cylindrical in the tail.

captures the indentations in the cusps.

3.4.2 Variability in Outer Boundaries

Szabo and Ronald P. Lepping (1995) used the Voyager 2 inbound crossing to categorise the shock as “a low β , high Mach number, and strong quasi-perpendicular shock”, which was moving away from the planet at the time of the crossing. Using remote observation of the shock, Iver H. Cairns et al. (1991) deduced the position and flaring of the shock is controlled by the rotation of Neptune’s magnetic dipole. The crossing of the bow shock by Voyager 2 was measured both by its plasma experiment (Belcher et al., 1989) and its magnetic field experiment (Ness et al., 1989). Two cylindrically symmetric parabolic models of the bow shock were obtained, which largely agree but have minor differences in stand-off distance and flaring. They are derived taking the two crossings of the bow shock as observed by Voyager 2 and hence do not account for the rotation of the dipole, nor the solar wind conditions, hence only give a rough guide for Neptune’s bow shock.

Figure 3.9 shows the cuts of the shock on the dayside in the $y = 0$ and $z = 0$ planes for the Earth-like and pole-on configurations, as well as the two empirical models by Belcher et al. (1989) and Ness et al. (1989). A good agreement is shown in figure 3.9 between the empirical models and the pole-on configuration with regards to the location and shape, although the simulation predicts a blunter bow shock than the models. According to John D. Richardson

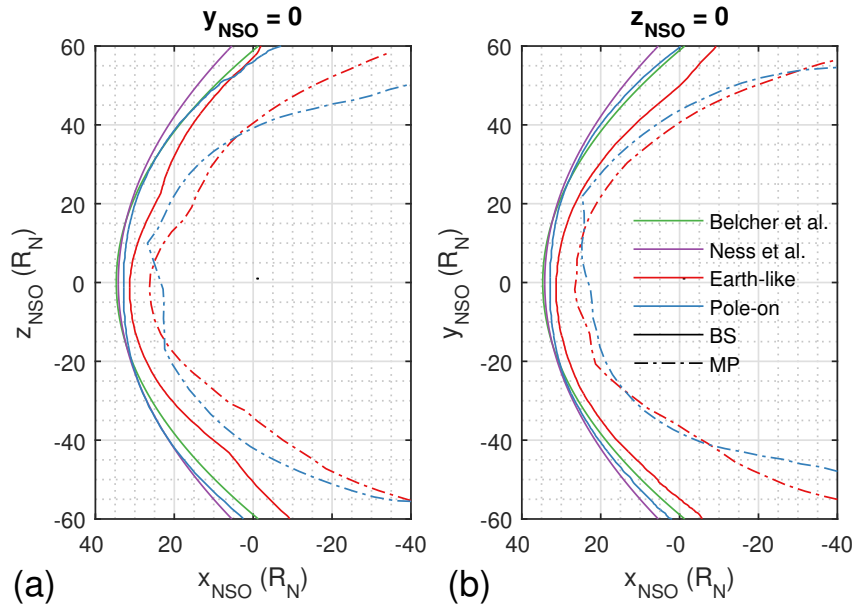


Figure 3.9: Bow shock and magnetopause profiles in the $y = 0$ (a) and $z = 0$ (b) planes. Also shown are the empirical Belcher et al. (1989) and Ness et al. (1989) bow shock models. The pole-on bow shock agrees well with the empirical models. It is much more blunt than the Earth-like bow shock, which follows the shape of the magnetopause.

et al. (1994), Voyager 2 was in the magnetosheath for approximately 3.5 hours; the observed bow shock location is therefore closer to the stand off distance of the pole-on bow shock.

Furthermore, the simulation reveals the bow shock is not cylindrically symmetric in either the pole-on or Earth-like configuration, being deflected in both the $y = 0$ and $z = 0$ planes. This asymmetric flaring could be due to the flaring of the magnetopause due to the IMF orientation, which is known to alter the shape in Earth's magnetopause (Lin et al., 2010). The Earth-like configuration in figure 3.9a has a cusp-like dip at approximately $z = 20 R_N$, which reflects the more complex magnetopause surface shape.

The Earth-like stand-off distance is roughly $4 R_N$ closer to Neptune than in the pole-on case, and exhibits a similar amount of flaring. This result confirms that the shock is controlled by Neptune's precessing dipole, as explored by Iver H. Cairns et al. (1991), but the simulation finds opposite behaviour: the stand-off distance is larger for the pole-on configuration because the magnetopause is more blunt.

3.4.3 Dayside Reconnection

Reconnection at Neptune has been studied before by Desch et al. (1991), R S Selesnick (1990), Huddleston et al. (1997), and Masters (2015), using either models or remote observations. Early studies focus on purely anti-parallel reconnection (e.g. R S Selesnick (1990)) as the only form of reconnection: now it is well understood that reconnection can still occur for smaller magnetic shear angles (e.g. Götz Paschmann, Marit Øieroset, and T. Phan (2013) and J P Eastwood et al. (2013)). The simulation shows evidence of reconnection occurring on the magnetopause, and that it is modulated by the daily rotation (see below). Though reconnection is fundamentally a kinetic process, whose entire dynamics cannot be captured by MHD, it is generally thought that global MHD simulations with numerical resistivity accurately predict the location of reconnection sites (Komar, Fermo, and Cassak, 2015).

To illustrate the reconnection process, figures 3.10 to 3.11 show snapshots in time of the magnetopause. 3D magnetic field lines and a slice of the velocity magnitude of the plasma are shown. The slice contains the IMF orientation (which is constant), and the solar wind flow direction, and centred through Neptune's origin ($x = y = z = 0$). It also shows the dipole moment direction. Reconnection is most likely to occur for anti-parallel magnetic field configurations, which geometrically will be located in this plane if they arise. In fact, reconnection is found to be confined to the northern hemisphere as is shown in figures 3.10 to 3.11.

The 8 panels show the progression from just before reconnection occurs in figure 3.10a, to when it stops in figure 3.11h, with 5,000 s ($= 0.086$ Neptune days) between each figure. Reconnection 'turns on' as the dipole moment rotates from the Earth-like to the pole-on configuration, approximately 0.2 Neptune days (τ_N) before pole-on, shown in figure 3.10c. There are high velocity jets shown in light-red visible at the northern magnetopause surface, with kinked magnetic field lines. After that, in figure 3.10d to figure 3.11f, reconnection continues to occur with the same high velocity jets and highly kinked magnetic field lines forming and moving downstream. These kinked field lines loop around themselves, indicating a flux rope type structure.

In figures 3.11g and later, approximately $0.2 \tau_N$ after pole-on, the field lines are smooth and no jets are visible, suggesting no reconnection is occurring in this plane. At this time, the local shear between the magnetospheric field and the IMF is 120° on the northern dayside. Hence, the

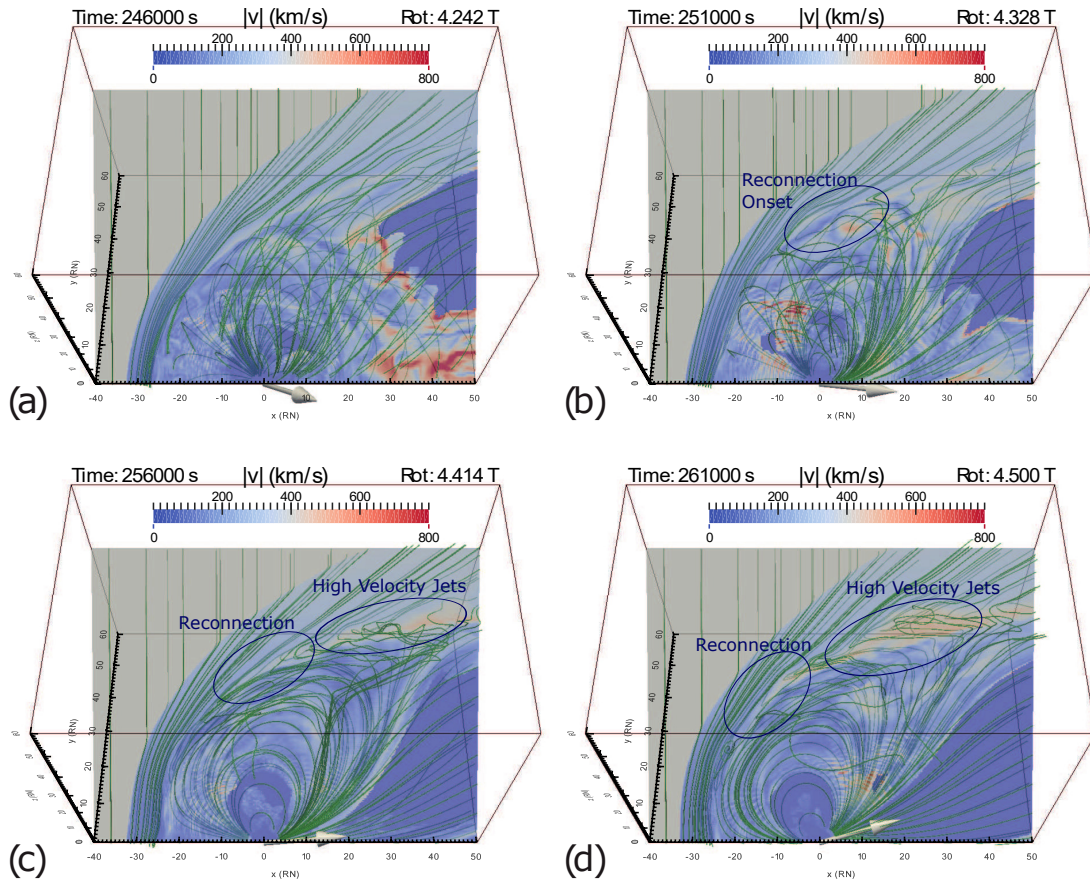


Figure 3.10: Onset of reconnection in the northern region of the dayside magnetosphere. Reconnection ‘turns on’ in panel (b) as the dipole rotates towards an Earth-like configuration, as shown by the kinked field lines and high velocity plasma.

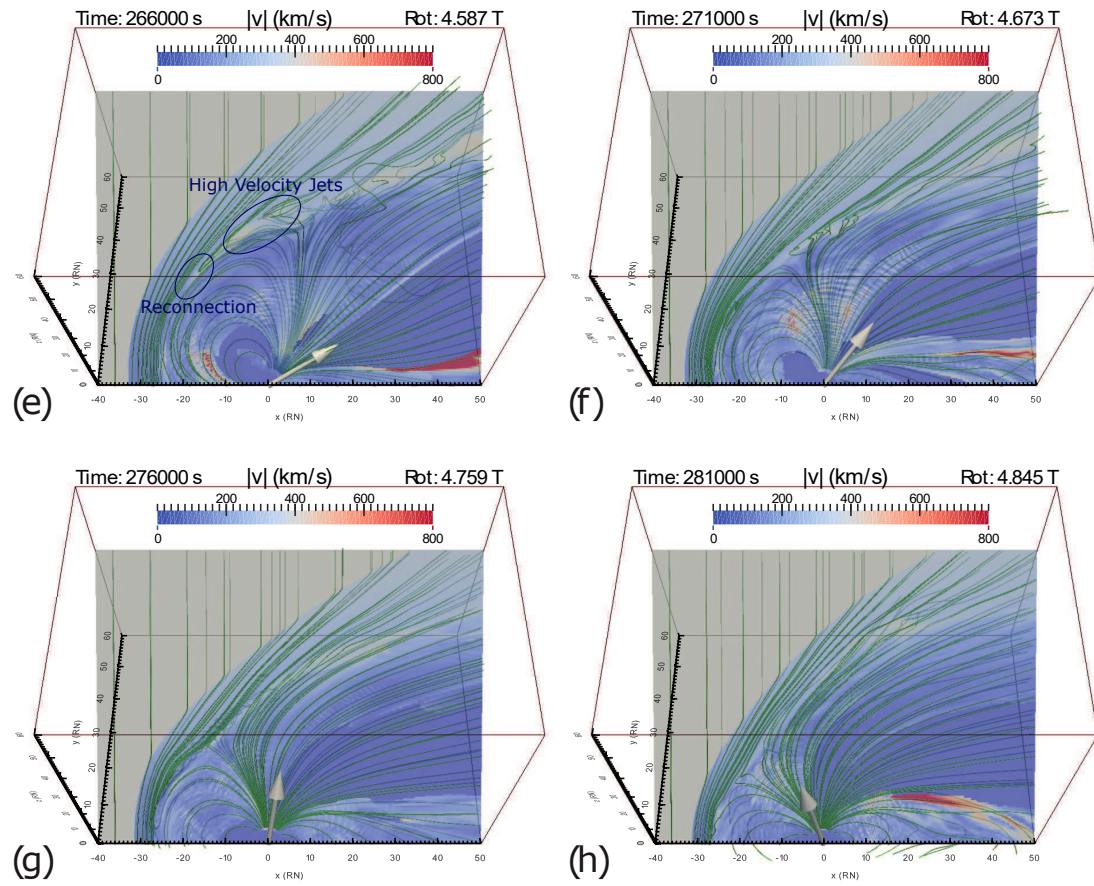


Figure 3.11: Reconnection persists until the northern lobe field reaches the magnetopause, stopping reconnection. It is seen until panel (e), where the field lines smoothen and no high velocity jets are seen.

reconnection appears to be controlled by the shear angle which, for the IMF condition observed, is largest for the configuration midway between the pole-on and Earth-like configurations. This gives a window of approximately $0.5 \tau_N$ (8 hours) during Neptune's rotation where reconnection occurs. This reconnection site is in some ways similar to the case at Earth with lobe reconnection under northward IMF. However, the effect of daily rotation heavily regulating the reconnection may suggest that there is little or no global circulation.

3.5 Magnetospheric Reconfiguration

Figure 3.12a shows a cut of the plasma density in the $y = 0$ plane. Overlaid are magnetic field lines, whose colour indicates the location in the y direction (blue $y < 0$, red $y > 0$). This corresponds to time-step $t = 290,000$ s, which is at the start of the fifth day. At this time, the angle between the magnetic dipole and the solar wind flow is the nearest to being perpendicular. The solar wind flows from left to right, and the solar wind magnetic field points in the $+y, +z$ direction. The bow shock forms ahead of the magnetosphere, and the magnetosheath corresponds to the dark red region. Closed field lines on the dayside present an obstacle to the solar wind and two well-developed cusp regions form. These cusps are tilted asymmetrically, due to the tilt of the dipole. The plasma density in the closed dayside field region is lower than in the magnetosheath. A boundary layer of plasma persists anti-Sunward of the cusps. Inside the magnetosphere, relatively empty plasma lobes are connected to the polar regions and map into the magnetotail as expected. The magnetotail lobes sandwich the plasma sheet, which largely corresponds to closed magnetic field regions on the nightside. Figure 3.12b shows the equivalent cut in the $z = 0$ plane at the same time. Again, the overlaid lines are magnetic field lines, but their colour corresponds to their position in the z direction. The overall configuration of the magnetosphere therefore resembles that of the Earth.

Figure 3.12c-d show the configuration of the magnetosphere half a planetary rotation earlier (time-step 261,000 s) in the $y = 0$ and $z = 0$ planes respectively. In this case, the dipole is almost parallel to the solar wind velocity. The bow shock and magnetosheath still form, though due to the dipole tilt, there is a cusp near the sub-solar magnetopause, corresponding to Neptune's southern magnetic pole. North of this region, within the magnetosphere, there is a low density plasma region (compared to the sheath), which exists on closed magnetic field lines. South of this region, the magnetosphere is filled with magnetic field lines emanating from the southern magnetic pole which extend into the tail, acting as a barrier between the magnetosheath and the tail plasma sheet (the light-red plasma region on closed field lines which extends into the tail). Finally, in the northern tail region, there is also a rarefied lobe plasma region corresponding to the magnetic field emanating from Neptune's northern magnetic pole, which also acts to contain the northern closed field plasma region. The dayside closed field region is limited to the vicinity of the planet by the magnetic field emanating from the northern magnetic pole and does not extend deep into the magnetotail. Figure 3.12c and 3.12d therefore show that the structure of the magnetotail for the pole-on configuration is somewhat different from that

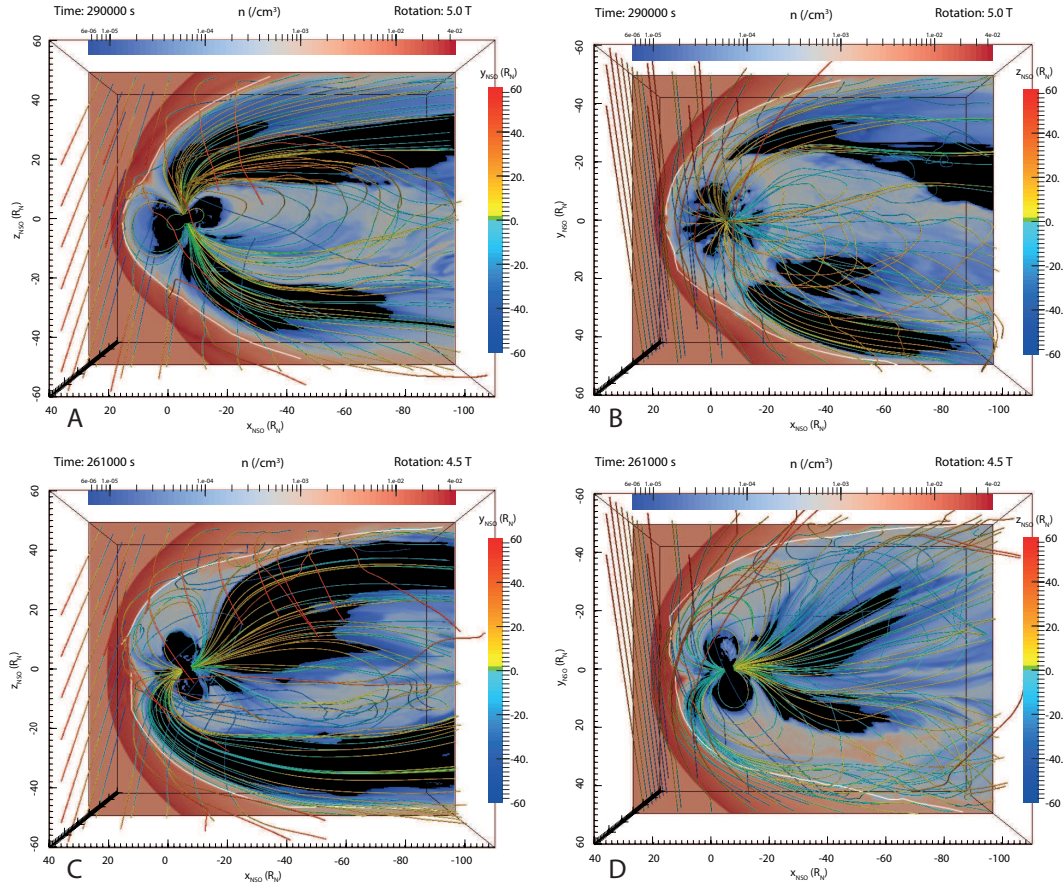


Figure 3.12: Slices of mass density in the $y = 0$ (left column) and $z = 0$ (right column) planes showing the reconfiguration of the two cases, Earth-like (top row) and pole-on (bottom row). The black regions denote where the density is at the floor density. The colour of the magnetic field lines denotes their position in the y_{NSO} - z_{NSO} plane. These show the plasma and magnetic field distribution changing dramatically between the pole-on and Earth-like configurations, moving the tail, dayside plasma regions and the positions of the lobes.

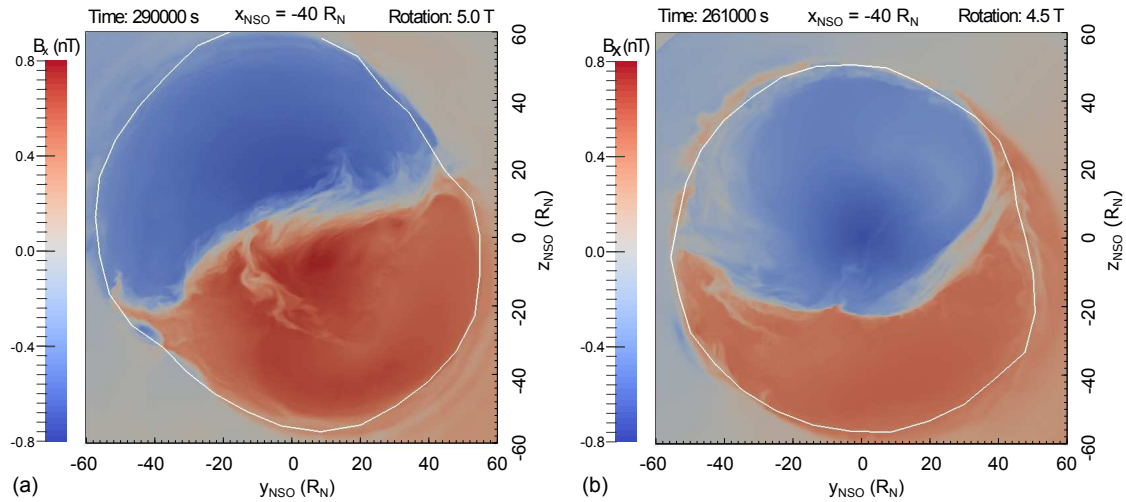


Figure 3.13: B_x component in the magnetotail, $40 R_N$ down-tail for Earth-like (a) and pole-on (b), looking down the tail. The magnetopause is given as the white line. The two configurations both show a twisted tail, with the Earth-like (a) elongating the magnetopause in the $-y + z$ direction. During pole-on (b), the tail does not close in on itself, as other models suggest, and is relatively symmetrical.

shown in figure 3.3.

Neptune's magnetic topology has been modelled with the source surface model (Michael Schulz and McNab, 1996) and empirical 'stretch' model (Voigt and Ness, 1990). It is predicted that near the pole-on case, for angles such that the solar wind velocity and dipole moments are parallel, the tail neutral sheet current is expected to close in on itself, rather than connect to the magnetopause current sheet.

By looking at the B_x component in the tail, aspects of the magnetotail shape and configuration can be inferred. Figure 3.13 shows this at a distance of $40 R_N$ downstream of Neptune, for the Earth-like and pole-on cases (see figure 3.12). The location of the magnetopause in this plane, calculated with the method explained above, is also shown. In both configurations, there are two hemispheres of oppositely directed B_x , separated by a current sheet, which is represented as $B_x = 0$. In the top left and bottom right corners of the figures, the draped B_x of the IMF can be seen. Overall, the tail is not cylindrically symmetric. In the Earth-like case, it is elongated in the $-y, +z$ direction, suggesting that the IMF acts to skew the tail in its magnetic field direction. The pole-on case shows a more symmetric tail magnetopause, which has been slightly shifted downward in the z direction.

In this simulation, the current system does not close in on itself, as can be seen in figure 3.13b. In both the Earth-like (figure 3.13a) and pole-on cases, the tail current sheet extends to each side of the magnetopause, connecting to the magnetopause current sheet.

Figure 3.13b gives the appearance that the blue region is enclosed by the red region (top left), possibly indicating a lobe like region which could suggest the tail current sheet is almost separated from the magnetopause current sheet. However, this is a region where reconnection occurs, containing the highly kinked field lines which connect the IMF to the planetary field at the dayside cusp. This time-step corresponds to figure 3.11d, where the kinked field lines can be observed at $x = -40 R_N$. These field lines do not connect to the dayside magnetic pole. In fact, they are connected to the nightside pole and are of the same topology as the blue region in figure 3.13b. Furthermore, this is supported by the position of the magnetopause, which finds that the red region at top left is outside the magnetopause.

3.6 The Voyager 2 Flyby

Voyager 2 arrived at Neptune from roughly the ecliptic plane, passing near the sub-solar bow shock and magnetopause. Neptune's rotation was such that Voyager 2 entered the magnetosphere through the cusp, in the near pole-on configuration. Neptune's gravity assist deflected the trajectory southward, where Voyager 2 crossed a plasma sheet, the magnetopause on the southern flank, and the bow shock again.

Figures 3.14 and 3.15 show slices of number density (n) and current density (j) in the plane which contains Voyager 2's trajectory, including annotations of the trajectory for the Earth-like and pole-on configurations. The positions of Voyager 2 in the simulation are also shown. Since the spacecraft is inside the magnetosphere for roughly three Neptune days, there are three positions of Voyager 2 corresponding to any particular dipole orientation. The order in which Voyager 2 is positioned at each green point is shown by the number next to it (e.g. Voyager 2 passes through point number 1 first, the point number 2, and so on.)

As shown in figure 3.14b by the first Voyager 2 position (point 1), the Voyager 2 spacecraft is very close to the magnetopause during the pole-on configuration, in keeping with the observation that Voyager 2 entered Neptune's magnetosphere through a cusp. As Voyager 2 arrives at its closest approach to Neptune, the dipole is now in its Earth-like configuration, as shown in

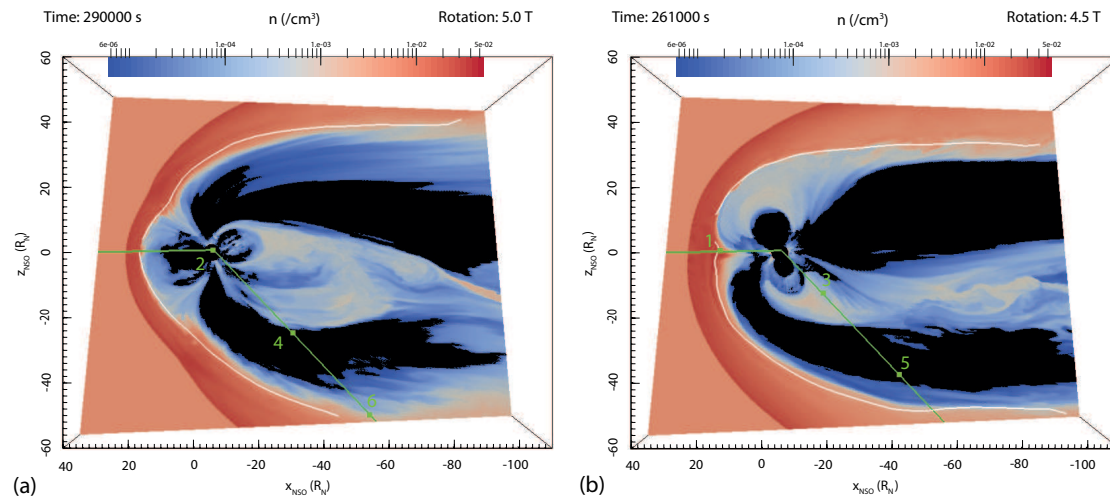


Figure 3.14: A slice of mass density containing Voyager 2's trajectory (floor density denoted in black). The numbers indicate the locations of the Voyager 2 spacecraft corresponding to the times when the Earth-like (a) and pole-on (b) configurations occurred. The magnetopause in this plane is denoted by the white line.

figure 3.14a at point 2. In the next Voyager 2 position, shown in figure 3.14b, point 3, the dipole is back to pole-on, and Voyager 2 is positioned in the magnetotail. In subsequent positions (4, 5 and 6), figure 3.14 shows that Voyager 2 passes through the blue southern lobe region, indicating a very low mass density. The outbound (or southern) magnetopause crossing is not contained within the simulation domain. This can be seen in the current density plots in figure 3.15, which shows the magnetopause as a peak in current density, indicating the magnetopause current sheet.

Voyager 2's flyby has been recreated by interpolating the magnetic field and plasma parameters along its trajectory, in accordance with its trajectory. Figures 3.16 and 3.17 show the inbound and outbound trajectories respectively. For the outbound trajectory (figure 3.17), only the magnetic field data is compared: for large parts of Voyager 2's trajectory inside the magnetosphere the ion flux was too low in order to calculate plasma parameters (John D. Richardson et al., 1991), except for in some isolated cases nearby the planet. Plasma observations are shown for the inbound trajectory, since there is complete data in the solar wind and magnetosheath, but no measurements were made after crossing the dayside magnetopause, except for an interval deep in the magnetosphere near to closest approach.

For the inbound trajectory in figure 3.16 the simulated number density (n), bulk velocity (v) and

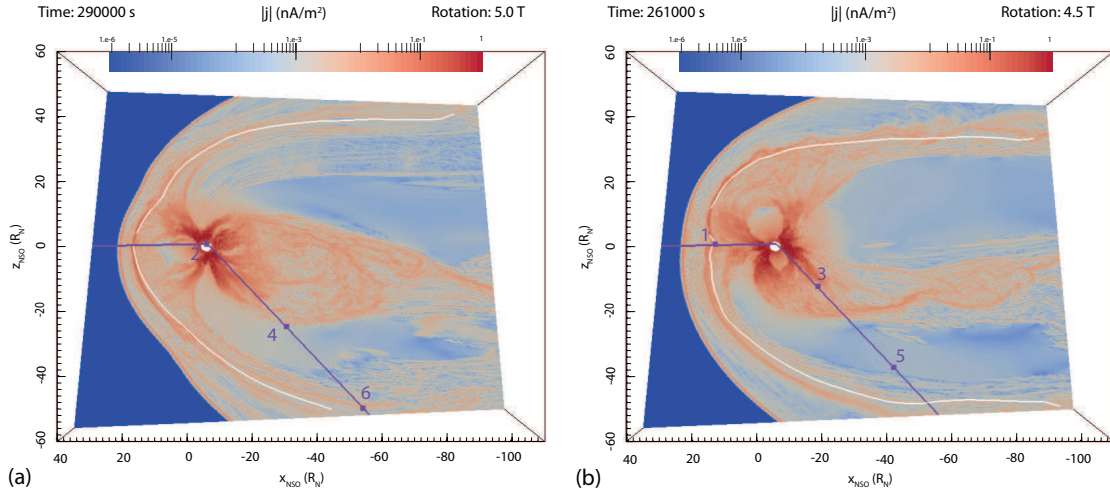


Figure 3.15: A slice of current density magnitude containing Voyager 2's trajectory. The numbers indicate the locations of the Voyager 2 spacecraft corresponding to the times when the Earth-like (a) and pole-on (b) configurations occurred. The magnetopause in this plane is denoted by the white line.

the proton and electron temperatures (T_p and T_e) are in good agreement with the Voyager 2 flyby data. No Voyager 2 ion measurements are available shortly after crossing the magnetopause, but cold dense plasma was detected closer to the planet. The simulation does not capture this, due to the fact it is initialized as an empty box, with no sources of mass inside the magnetosphere. The inbound magnetometer data agrees well with observations too, capturing the bow shock, magnetopause and the time varying magnetic field within the outer magnetosphere. Gorgon predicts a bow shock approximately $2 R_N$ closer than that measured by Voyager 2, also seen in the bow shock profiles in figure 3.9. The magnetopause is found to be in the same place, as seen in all components, but most noticeably the change in sign of the B_y component.

There is also a good agreement in all components of the magnetic field in the outer magnetosphere on the outbound flyby, as shown in figure 3.17. However, the B_x is slightly smaller in this region than in the observations. The peak in B_x and B_z at around Aug-25 17:00, is captured by the simulation and is due to the rotation of the magnetic dipole. Nearer the outbound magnetopause crossing Aug-26 02:00, there is a difference in sign in the B_y component between Gorgon and the Voyager 2 observation.

Gorgon predicts the flank magnetopause to be further out than Voyager 2, to the extent that it is not contained within the simulation domain. The change in sign in B_y at Aug-26 09:00 could have indicated a magnetopause. However, since no decrease in B_x is observed, and the magnetic

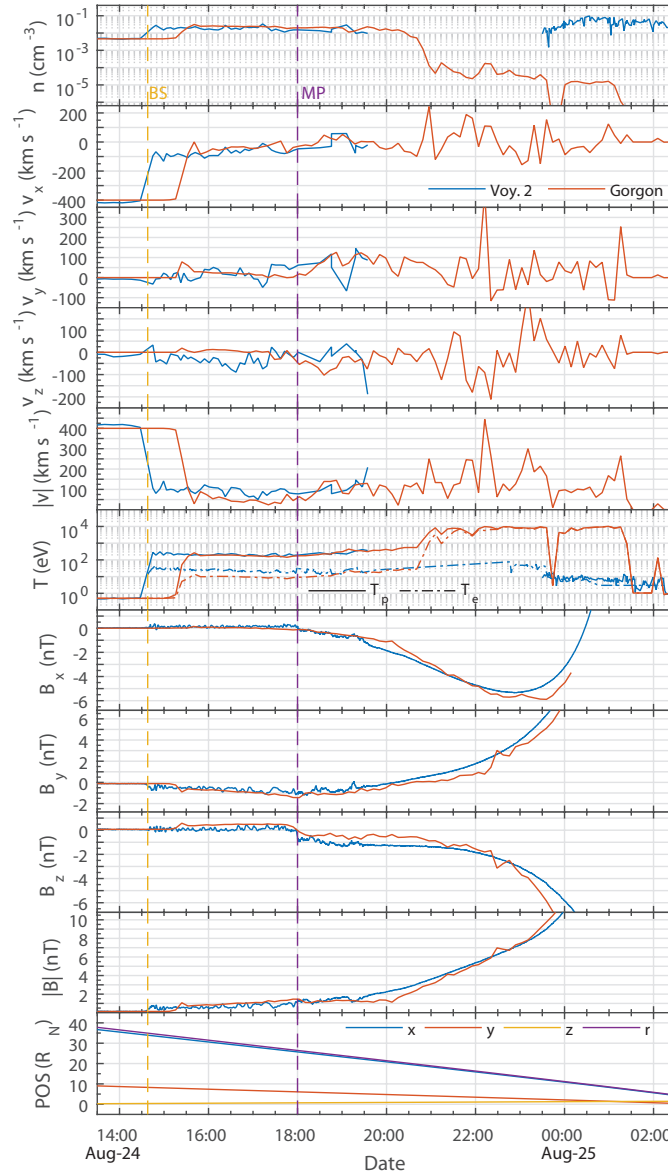


Figure 3.16: Voyager 2 observations compared to Gorgon. Starting in the solar wind, the Voyager 2 probe observes the bow shock at approximately 14:30 on Aug-24, followed by the magnetopause at 18:00. Shortly after crossing the magnetopause, the density was too low in order for Voyager 2 to measure the plasma data. Gorgon data agrees well, capturing the shock ratio well in the number density (n), velocity (v), ion temperature (T_i) and magnetic field (B), and predicting the position of the magnetopause. The location of the shock is closer to Neptune in the simulation, and Gorgon does not capture the high number density close to the planet.

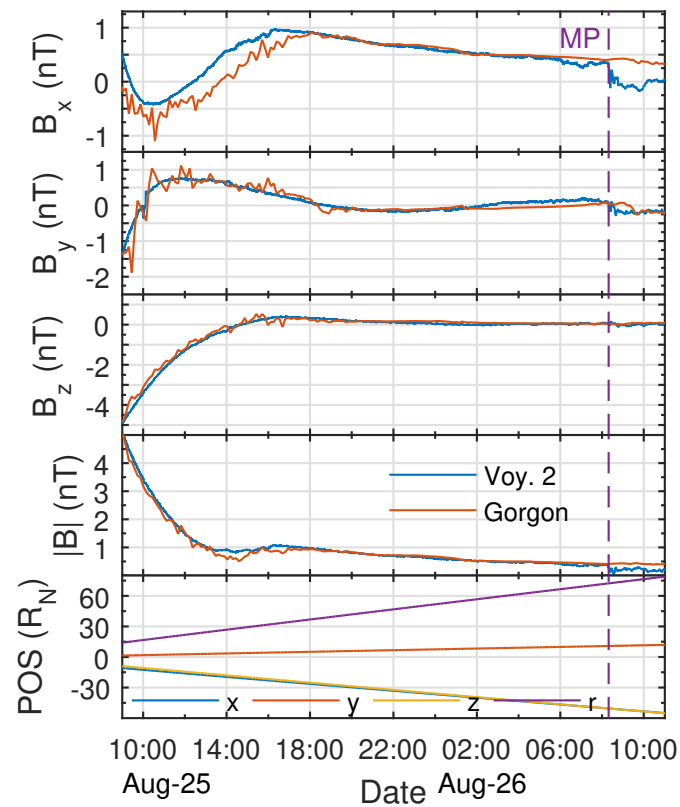


Figure 3.17: Outbound magnetometer data from Voyager 2 compared with Gorgon from 09:00 Aug-25 to when it crosses the magnetosphere on Aug-26. Gorgon agrees well with Voyager 2 observations in the tail, despite not capturing the outbound magnetopause.

field lines drawn from this region are still connected to Neptune, it is still in the magnetosphere. This is also supported by our method of finding the magnetopause using momentum streamlines. For the Voyager 2 comparison, it is worth noting that the simulation is run with a steady solar wind, and hence there is no observation on how the solar wind has varied once Voyager 2 has entered the magnetosphere, and this could account for discrepancies with the simulation and observed data.

3.7 Conclusion

The goal of this chapter is to understand the nature of the outer magnetospheric boundaries (including reconnection), the magnetotail structure and the large-scale behaviour of magnetospheres created by highly-inclined precessing dipoles. Neptune's magnetosphere was simulated and compared to the Voyager 2 flyby. Overall, the simulation performs well and shows good agreement both the shock models and the Voyager 2 flyby for the large-scale features.

The use of a global MHD code to simulate Neptune's magnetosphere has allowed a number of new results have also been obtained. It shows that Neptune undergoes global reconfiguration during its daily rotation, with effects that are intrinsically time-dependent and three-dimensional in nature. These effects would be difficult to capture in a steady state or equilibrium model, such as in previous work (Voigt, 1981; Voigt and Ness, 1990; M. Schulz et al., 1995), emphasising the need for three-dimensional global MHD simulations. Additionally, the distinct combination of the rotation axis tilt and magnetic field tilt of Neptune means its magnetosphere cannot be approximated by a steady state solution in any rotating frame, making it distinctly different to Uranus (Tóth et al., 2004).

As predicted in previous studies, the window for reconnection to occur is found to be small, occurring intermittently on the northern-dawn side region for the given IMF and season. This has a number of implications on Neptune's magnetosphere. Firstly, it suggests that a similar mechanism of inhibiting the opening and closing of magnetic flux that was proposed by Cowley (2013) for Uranus is also the case for Neptune. The small window in which dayside reconnection occurs makes a Dungey-like cycle difficult to set up for Neptune's magnetosphere. Secondly, the position of the reconnection site could suggest a reason for the lack of dynamics measured by Voyager 2 on its outbound trajectory, southward of the planet. Finally, the fact that reconnection occurs on the northern dawnside magnetopause could also be a factor in

influencing the magnetotail by preventing the dayside closed field region extending deep into the nightside magnetosphere.

Though a good agreement is found with the Voyager 2 flyby on the inbound trajectory with respect to the outer boundaries, the simulation does not capture the plasma close to the planet. The simulation was initialised empty of plasma with no mass sources other than the solar wind, hence the plasma density is expected to be low in this region. Additionally, the outbound magnetopause is not captured. However, the solar wind is kept steady at the last known parameters that were observed. The arrival of a completely new and unmeasured solar wind front may explain any discrepancy in the location of the outbound magnetopause crossing.

There are many avenues in which this work can be expanded in the future. Simulating the Voyager 2 encounter with the Neptunian system has allowed for verification of the Gorgon model. However, in reality, the solar wind is variable, which will impact the dynamics of Neptune's outer boundaries. The role of a variable solar wind on Earth's magnetosphere will be studied in chapter 4: this methodology can be applied to Neptune's magnetosphere. By fully incorporating the variability of the solar wind, one can determine whether a Dungey cycle can ever form for a significant amount of time at Neptune. Furthermore Masters (2015) showed that also the season of Neptune affects reconnection on the magnetopause: this work can be extended to investigate the full effect of the season on the three-dimensional magnetosphere. Finally, this analysis can also be applied to other highly inclined, precessing magnetospheres, using Uranus as a test case, or moving into a more general realm of varying tilts and rotations. This general case of tilts has the potential to make useful insights into the workings of exoplanets, a large proportion of which are thought to be Neptune-like (Batalha et al., 2013).

Future work to explore the inner magnetosphere may benefit from a planetary magnetic field model containing higher-order moments, which has implications closer Earth. During times when the Earth's internal dynamo is undergoing a reversal, it is thought the dipole moment of the Earth's magnetic field diminishes in strength and the higher order moments dominate (Vogt and Glassmeier, 2001; Zieger, 2004), becoming somewhat similar to Neptune's magnetic field.

Finally, there is the open question on the possible role of Triton on Neptune's magnetosphere. Whilst Triton was expected to play some role, its effect on the system could not be determined by the Voyager 2 encounter (John D. Richardson et al., 1991; Zhang, John D. Richardson, and Sittler, 1991). By implementing Triton as a potential plasma source in an MHD simulation,

useful insight can be gained by varying its effect and assessing the change in magnetospheric response as a result, allowing determination the properties of Triton.

Whilst there are no spacecraft arriving at either of the Ice Giants in the immediate future, there are science cases being made for such a mission (Christophe et al., 2012; Masters et al., 2014). In the absence of observational data, theoretical predictions and simulations are of great use in planning such missions, potentially shedding light on the dynamics of the system.

4 Motion of the Outer Boundaries under Variable Solar Wind

The work in this chapter has been published in *Journal of Geophysical Research: Space Physics* in Mejnertsen et al. (2018).

4.1 Introduction

The interaction of the solar wind, streaming out radially from the Sun, with the magnetosphere leads to the formation of a bow shock upstream of the magnetopause, with a turbulent sheath region (the magnetosheath) in between (Ness, Searce, and Seek, 1964; D. H. Fairfield, 1976; Lucek et al., 2005; Burgess and M. Scholer, 2013; Krasnoselskikh et al., 2013). The main role of the shock is to decelerate and deflect the supersonic solar wind around the magnetosphere, however, reflection of particles at the shock leads to the formation of the foreshock which leads to numerous wave-particle interactions and particle acceleration (J. P. Eastwood et al., 2005).

4.1.1 Modelling the Earth's Bow Shock Position

Early work on modelling the Earth's bow shock included quantifying the position and shape of the bow shock under different solar wind and IMF parameters, and the position and shape of the magnetopause. Following on from the Spreiter, Summers, and Alksne (1966) 'hydromagnetic' model, various empirical models were derived using a combination of spacecraft data for statistical fitting and gas dynamic theory for the basic shock shape (Filbert and Kellogg, 1979; Formisano, Domingo, and Wenzel, 1979; Farris and C T Russell, 1994; Peredo et al., 1995).

Empirical models, constructed from databases of spacecraft shock crossings, usually assume several first order approximations. Most models employ axisymmetric conic sections for the bow shock, although more complex models have been developed (Schwartz, 1998). The magnetopause is also assumed to be a similar conic section: however, Lin et al. (2010) derives a full three dimensional magnetopause shape, including cusps, from spacecraft observations. These models have shown that the bow shock shape and position is largely determined by the solar wind dynamic pressure, and to a lesser extent, the solar wind Alfvén Mach number (Peredo et al., 1995; Merka, 2005). In using empirical models, it is effectively assumed that the solar wind pressure is uniform over the bow shock surface. It is also assumed that there is an instantaneous relationship between changes in the magnetopause and bow shock location, and they do not account for the relevant signal speeds in the magnetosheath, nor the motion of the boundaries. There are a number of empirical models in the literature (e.g. Merka (2005), Iver H Cairns and J G Lyon (1995), and Jeřáb et al. (2005)), all derived using different datasets and reference surfaces. Though these models differ from one another when applied under the same solar wind conditions, they usually agree with one another to within a few R_E , as shown in figure 4.1. Comparisons of these models have been performed in numerous papers. For example, Turc et al. (2013) compared a number of models when the solar wind Alfvén Mach number was low, cross-referencing them with Cluster crossings of the bow shock. They found that, in low Alfvén Mach numbers, the Jeřáb et al. (2005) model performed best, with the Merka (2005) model performing the worst.

4.1.2 Variability in the Shock Position

The launch of Cluster (Escoubet, Fehringer, and Goldstein, 2001) led to routine calculation of shock normal speeds and orientations which can be obtained using four spacecraft timing analysis (Schwartz, 1998). Subsequent work using the Cluster four-spacecraft timing analysis has shown that model prediction of position and shape typically perform reasonably well when compared with observations (T. S. Horbury et al., 2002; Maksimovic et al., 2003; K. Meziane, Alrefay, and Hamza, 2014).

By assuming the shock is planar over the scale of the spacecraft separation (typically 600 km at the shock (Escoubet, Fehringer, and Goldstein, 2001)), one can determine the shock motion and

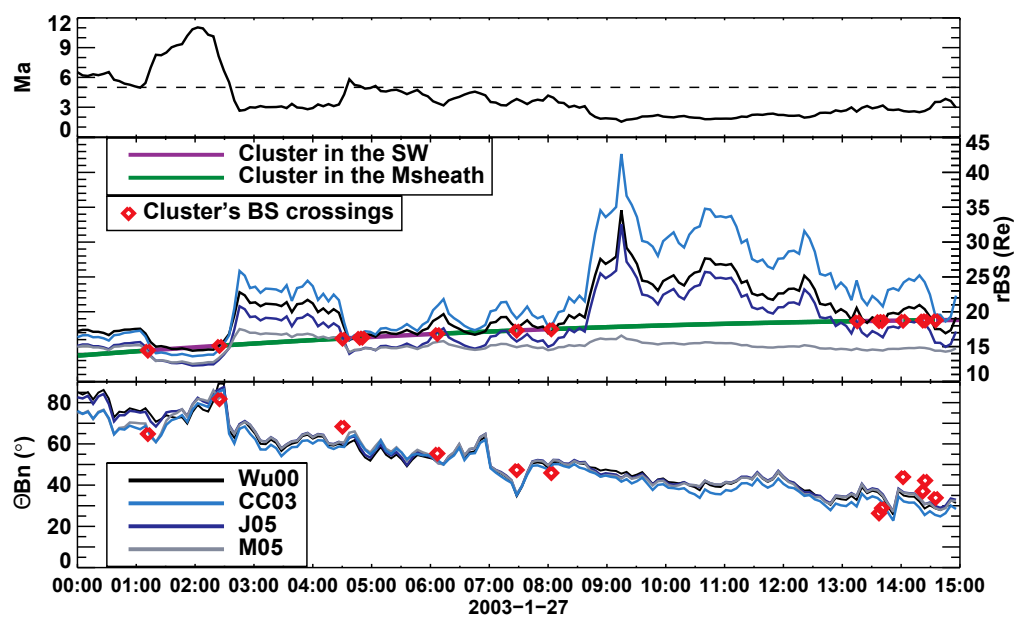


Figure 4.1: Comparison of bow shock models (D. J. Wu, Chao, and R. P. Lepping, 2000; J. F. Chapman and Iver H. Cairns, 2003; Jeřáb et al., 2005; Merka, 2005) and Cluster bow shock crossings (mid panel) under low solar wind Alfvén Mach numbers (top panel). Three of the four models agree on the general motion of the shock, but disagree in magnitude by approximately $10 R_E$ at the most. Figure from Turc et al. (2013).

normal by comparing the times each spacecraft crosses the shock, using the relation,

$$\left(\vec{V} (t_i - t_j) \right) \cdot \vec{n} = (\vec{r}_i - \vec{r}_j) \cdot \vec{n}, \quad (4.1)$$

where $\vec{V}$ is the shock velocity, $\vec{n}$ is the shock normal, and $t_{i,j}$ and $\vec{r}_{i,j}$ are the time and position of the crossing for two different spacecraft. T. S. Horbury et al. (2002) used the magnetic field data from Cluster to identify each shock, and analysed 48 crossings of the shock where the shock normal and velocity could be obtained. They found that typical shock velocities were ~ 35 km/s, and that the shock normal agrees with empirical models. Comparisons with the coplanarity method of finding shock normal, which uses single spacecraft, performed poorly. Maksimovic et al. (2003) applied a spacecraft timing analysis to a specific solar wind event that occurred in March 2001. During this event, multiple spacecraft crossings were observed. By calculating the shock velocity, the shock position could be inferred between crossings, and found that gas dynamic models agreed on shorter timescales of the order of the solar wind variability. K. Meziane, Alrefay, and Hamza (2014) extended T. S. Horbury et al. (2002)'s analysis by including more shock crossings, agreeing with T. S. Horbury et al. (2002)'s findings. It was found that slow shock speeds (< 80 km/s) follow a Gaussian distribution of width 42 km/s. As found in the derivation of the empirical models, the spacecraft observations show the bow shock velocity is largely driven by the changes in the solar wind dynamic pressure, and has little dependence on Mach number (K. Meziane, Alrefay, and Hamza, 2014). Furthermore, the speed of the bow shock, as calculated from the time derivative of empirical model shock position, agrees statistically with observation as well (Meziane et al., 2015).

Moving beyond typical solar wind variability, discrete structure (specifically, interplanetary shocks) causes rapid and significant boundary motion (Ivanov, 1964; Dryer, Merritt, and Aronson, 1967; Shen and Dryer, 1972; Grib et al., 1979). Early work included Greenstadt, Hedgecock, and C T Russell, 1972, who used the Ogo 5, Heos 1 and Explorer 33 to calculate the shock velocities. Völk and Auer (1974) applied a theoretical analysis of shock motion by considering the shock under different solar wind Alfvén Mach numbers. The observed high shock velocities under low Mach numbers originate due to an increased sensitivity to Alfvén waves, whereas high Mach numbers and associated high shock velocities are due to the variability introduced by travelling interplanetary shocks (Völk and Auer, 1974).

A number of studies have simulated the effect of an observed interplanetary shocks (IPS) on the Earth's magnetosphere, finding that the shock-bow shock interaction causes a new train of discontinuities (Šafránková et al., 2007; Přech, Zdenek Němeček, and Šafránková,

2008; Pallocchia et al., 2010; Pallocchia, 2013). As they propagate through the shock and sheath, these have been shown to create subsequent shocks, which deflect and reflect off the magnetopause. Simulations by Samsonov, Sibeck, and Imber (2007) show Earthward bow shock and magnetopause motions as the interplanetary shock impacts, but outward motions due to a reflected shock. Similar results have been observed by spacecraft (Pallocchia, 2013). Such shocks have a profound effect on the magnetosphere, creating disturbances which propagate through the magnetosphere at speeds higher than the original disturbance in the solar wind (Wang, 2005; Xiao-Cheng, You-Qiu, and Chi, 2005; You-Qiu et al., 2005; Andréová et al., 2008), and create significant responses in the inner magnetosphere (Andréová et al., 2011).

4.1.3 Open Questions and Motivations

In situ data provides a wealth of information about the local plasma properties, and the various physical processes controlling thermalisation, particle acceleration, etc. However, knowledge of the context of these in situ measurements is required to fully understand the observed behaviour. Variability in the solar wind causes dynamic movement of the outer magnetospheric boundaries as they respond to changes in pressure and interplanetary magnetic field (IMF). Consequently, empirical models of the shock shape and location, derived statistically from spacecraft observations, are crucial in this regard. They are fundamental to many spacecraft investigations of shock and foreshock physics, as they are used to place local measurements in the global context, and to correctly interpret shock geometry. However, whilst such models agree well with the average location and shape of bow shock or magnetopause, they often need to be scaled in order to agree with the specific spacecraft crossing location for any particular event (Schwartz, 1998).

Empirical models are therefore limited in the extent to which they can capture shock motion that occurs in response to rapid variation in the upstream solar wind. In particular, it is not well understood as to when these statistical models become no longer appropriate, nor is it clear as to how these models differ from reality when taking into account the dynamic transition of boundaries to new equilibrium states. Global magnetohydrodynamic simulations are a useful tool in this regard, as they can simulate the whole system self consistently.

In this chapter, Earth's magnetosphere is simulated with a variable solar wind, focussing on the motion of the bow shock. The documented case study from Cluster covering variable solar

wind from Maksimovic et al. (2003) is considered.

4.2 Simulation Set-Up and Solar Wind Input Conditions

In this work, the region of interest contains the outer magnetospheric boundaries. Hence, a simulation area spanning $x_{\text{GSE}} = -50 R_E$ to $30 R_E$, $y_{\text{GSE}} = -20 R_E$ to $20 R_E$, $z_{\text{GSE}} = -30 R_E$ to $30 R_E$ is used, with a uniform grid size of $0.2 R_E$. To avoid complications associated with dipole tilt or rotation in this analysis, it is aligned to the z_{GSE} direction. The inner boundary is set at $3 R_E$, and is centred on the origin.

One-minute resolution data observed on 31st March 2001 were used in the simulation input, corresponding to the Cluster spacecraft case-study performed by Maksimovic et al. (2003). The data were first ballistically propagated from the location of ACE to the Sunward edge of the simulation box ($x_{\text{GSE}} = 30 R_E$). Since the simulation timestep is smaller than 1 minute, the parameters were linearly interpolated between observations. The interval of interest is 17:14:00 UTC – 19:34:00 UTC. The simulation was initialized with the preceding two hours of ACE observations. Consequently, a simulation time of $t = 0$ min corresponds to 15:14:00 UTC, and the analysis interval begins at $t = 120$ min.

The simulation input data corresponding to the analysis interval is shown in figure 4.2. The proton number density (panel a of figure 4.2) ranges between $10/\text{cm}^3$ and $30/\text{cm}^3$. The average solar wind speed is approximately 550 km/s (panel b). The strength of the IMF (panel c) was larger than typical IMF strength, ranging between 20 nT to 35 nT. The IMF orientation is strongly southward, with B_z ranging between -32 nT to -20 nT. Since the solar wind velocity is predominantly in the $-x_{\text{GSE}}$ direction, B_x cannot be propagated into the box. A constant B_x is initialised throughout the simulation domain, whose magnitude corresponds to the average B_x measured by ACE. With these solar wind and IMF parameters, the solar wind Alfvén Mach number is low, averaging $M_A \sim 5$. Empirical models typically do not perform as well under low Mach numbers as there are fewer crossings in this regime: the best model for low Alfvén Mach numbers has been established to be the one proposed by Jeřáb et al. (2005) (Turc et al., 2013), and is used here to compare with the simulation output.

The state of the magnetosphere at the beginning of the analysis period ($t = 120$ min) is shown in figure 4.3 It shows the number density and the dynamic pressure in slices along the ecliptic

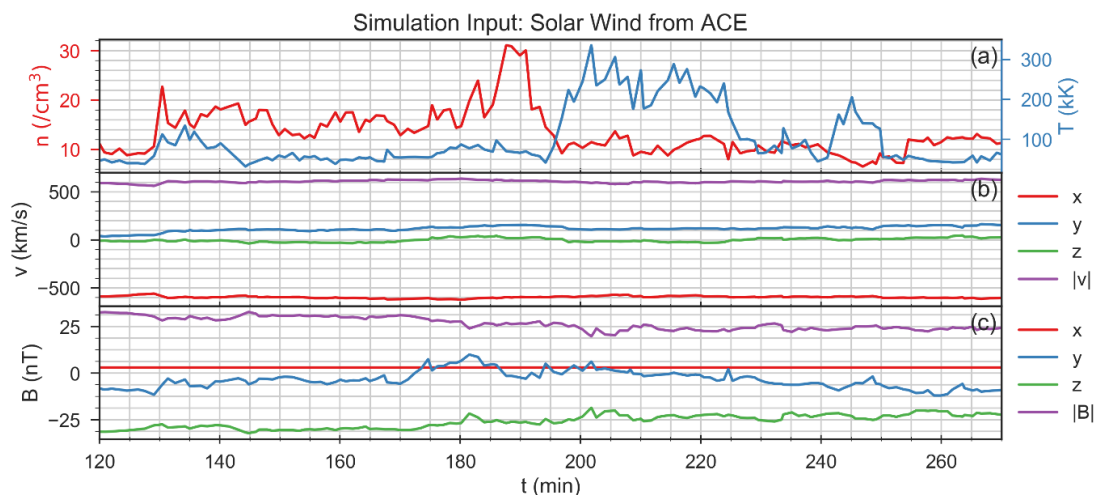


Figure 4.2: Solar wind input to the simulation, as observed by ACE (Maksimovic et al., 2003).

and noon-midnight planes. At this timestep, the simulation has formed a full magnetosphere with its characteristic features: a low density solar wind, arriving from the left, the bow shock and the turbulent magnetosheath, and a low density magnetosphere. In figure 4.3a and b there is a clear asymmetry shown in the tail position, deflecting in the positive y direction. Those panels also show structure on the magnetopause, at approximately $x = 4 R_E$, $y = -7 R_E$: the kinks produced in the magnetopause position are due to flux ropes produced on the dayside, and will be discussed in Chapter 5.

4.2.1 Locating the Bow Shock

To characterize the shape and location of the shock, these properties are examined in the $x_{GSE} - y_{GSE}$ (ecliptic) plane and the $x_{GSE} - z_{GSE}$ (noon-midnight) plane. In this sub-section, the shock finding algorithm is explained by referring to the shock in the ecliptic plane. The same method is used for the polar plane, except all instances of y_{GSE} are substituted by z_{GSE} .

First the bow shock is located along the x axis, by seeking the compression of plasma due to the shock. Technically, this is done by calculating the mass density gradient along the x axis, and identifying negative peaks, as illustrated in figure 4.4a. In the simulation, there are often many such peaks, only one of which corresponds to the bow shock. To determine which is the shock, ‘significant’ peaks are found: ones whose scaled magnitude is larger than a set value. This method is found to work well as the bow shock is a strong shock, so it is usually the largest

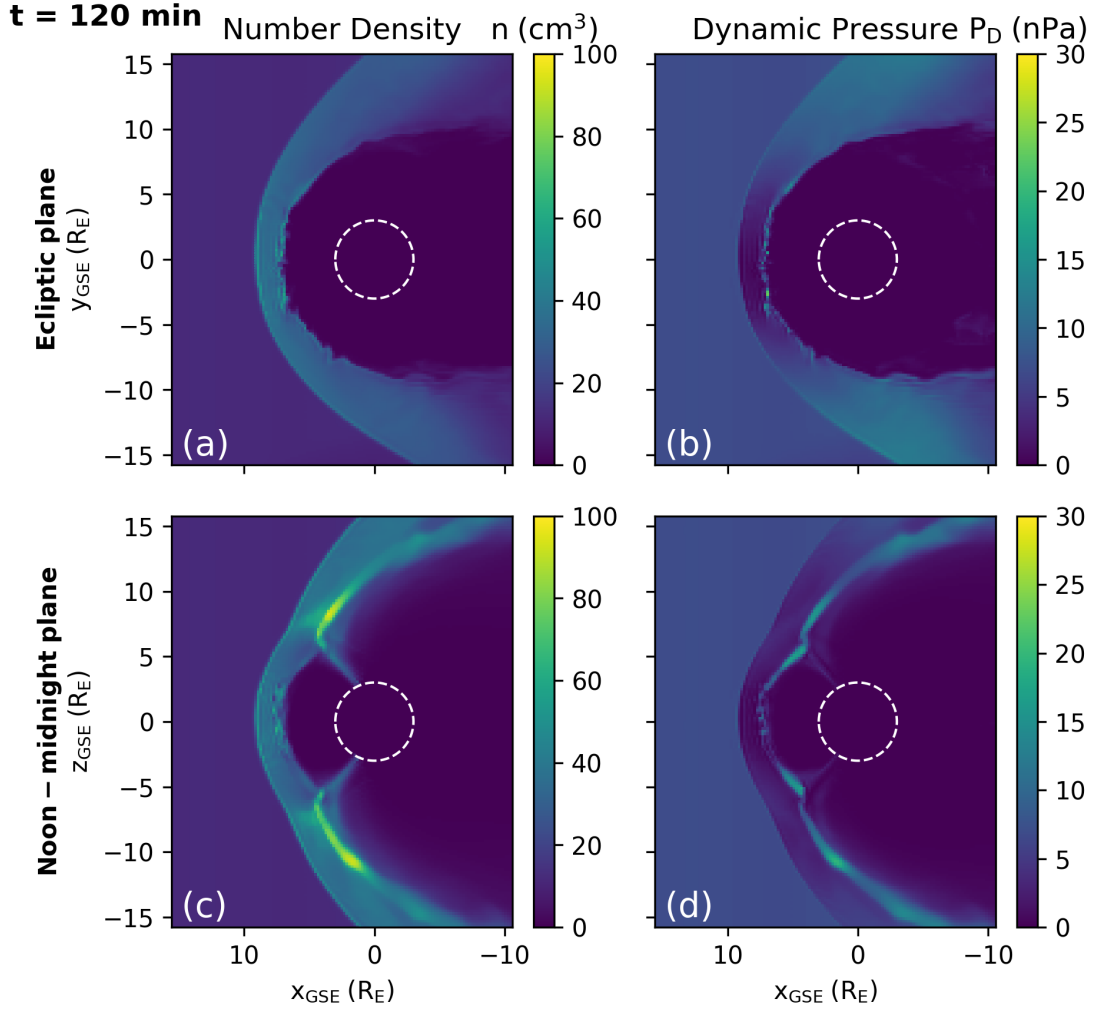


Figure 4.3: Slices of the number density (panels a and c) and dynamic pressure (panels b and d) through the Ecliptic (panels a and b) and Noon-midnight planes (panels c and d) of the simulation, two hours after initialization. The solar wind streams in from the left. The bow shock is visible by the increase in number density (panels a and c) and decrease in dynamic pressure (panels b and d). The magnetosphere is largely denoted by the regions of low number density and dynamic pressure, shown in dark blue. The white dashed line denotes the inner boundary of the simulation, at $3 R_E$. In panel d, there is an enhanced dynamic pressure along the magnetopause, which is due to reconnection jets on the dayside. In panel a and b, there is a kink in the magnetopause position at $x = 4 R_E$, $y = -7 R_E$: this is due to a flux rope propagating downstream along the magnetosphere.

such peaks. If there are multiple ‘significant’ peaks, the peak most Sunward is chosen, as there tend to be no such peaks in the solar wind.

In global MHD simulations, the position of the shock is always located on the grid cell interface, since Eulerian simulations are a set of calculations solving the MHD equations on finite sized cells. To overcome this, a sub-grid method was used, as well as smoothing of the bow shock standoff position curves. Once the location of the grid position of the bow shock has been found, the bow shock peak position and the two adjacent grid points in the x -direction are taken and a parabolic function was fitted ($d\rho/dx = ax^2 + bx + c$), passing through these three points. This is shown as the purple curve in figure 4.4a. The sub grid shock position is given as the maximum of this parabola, $x_{BS} = -a/2b$, and corresponds to the black dashed line. This process was repeated for all y_{GSE} grid values, and for each time-step. Hence the 2D shock position $x_{BS}(y_{GSE}, t)$ is obtained as a function of y_{GSE} and t . The result of this analysis is shown in figure 4.4b (red curve).

The object of this analysis is to explore the motion of the bow shock. The corresponding shock speed along the x -direction is shown in figure 4.4c (red curve). Although the large-scale motion of the shock is visible, the sub-grid parabolic fitting method does not completely eliminate the impact of the grid as the shock passes from one cell to the next. Therefore, the shock position is smoothed as a function of time using a Savitsk-Golay filter: this algorithm fits a polynomial to the data in a rolling window (blue curve in figure 4.4b). Once smoothed, a 2nd order centred difference gradient is used to obtain the shock velocity, as shown in blue in figure 4.4c. This combined filtered sub-grid method removes grid fluctuations in the shock position and velocity, leaving the underlying behaviour visible. This method is simpler and more effective than using a Rankine-Hugoniot approach, which assumes a steady-state solution that is not generally the case in this simulation, and mitigates issues which arise in how the upstream and downstream definitions are determined.

The slight overshock present in the mass density curve (red) in figure 4.4a is a numerical effect of the method used in solving the MHD equations on a grid. Many MHD codes use an artificial viscosity to prevent such overshoot, though if the viscosity is too large, the shock is smeared over more grid cells, leading to a wide shock. The parabolic fitting method and temporal smoothing used here overcomes this numerical effect and minimizes the effect of the overshoot.

Finally, the sub-solar magnetopause is located by identifying the magnetopause current sheet, and the corresponding peak in the magnitude of the current density. The sub-grid magnetopause

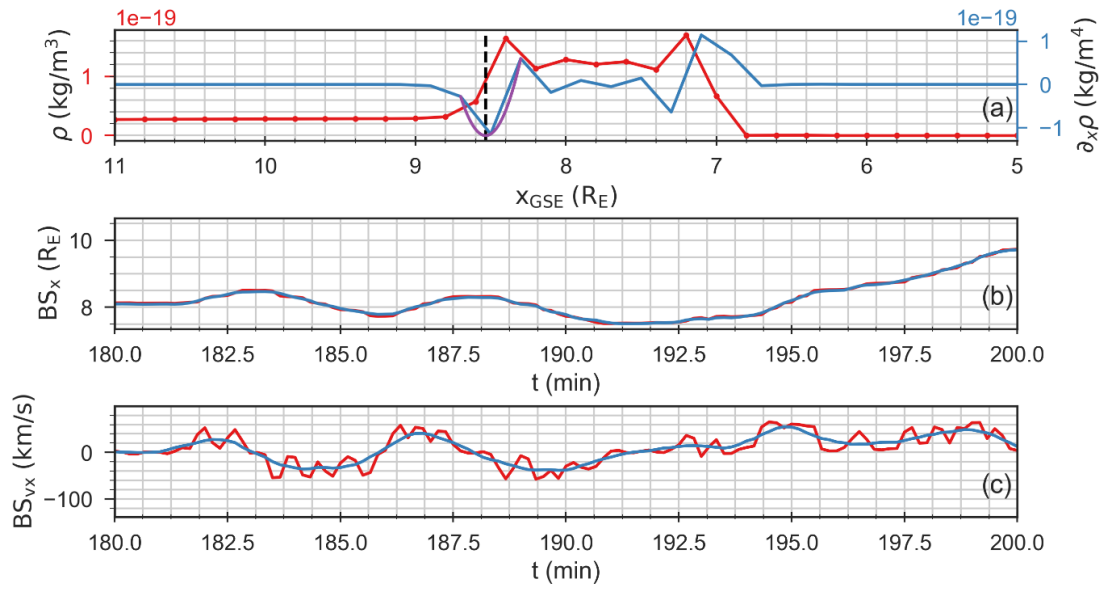


Figure 4.4: Overview of the method used to find the bow shock position and smoothing. (a) shows the mass density profile (red) and its gradient (blue) along the Sun-Earth line. The region of solar wind is to the left. The location of the shock is identified as the black dashed line. The purple curve shows the parabola through the points near the shock. Panels (b) and (c) show the effect of smoothing on the shock position and shock velocity with time.

position and speed can then be located using the same technique as for the bow shock.

The empirical models are calculated as described in Jeřáb et al. (2005) for the bow shock and Shue et al. (1998) for the magnetopause, using the solar wind conditions as input into the simulation. The velocity of the model position is calculated using the same smoothing techniques as with the simulation positions.

4.3 Results

4.3.1 Comparison along the Sun-Earth line

Figure 4.5 shows how the sub-solar bow shock and magnetopause react under variable solar wind conditions. Panel (a) shows the solar wind dynamic pressure (green) and Mach number (purple). Panel (b) shows the magnetopause and bow shock location on the Sun-Earth line, for both empirical models (blue) and Gorgon (red). The empirical models used here are the Jeřáb et al. (2005) bow shock model, and the Shue et al. (1998) magnetopause model. Figure 4.5d and figure 4.5e show the velocities of the bow shock and magnetopause in the x_{GSE} direction.

Both empirical have models been scaled so that their average location is comparable to the simulation: this has been done in order to investigate the dynamics of the outer boundaries. Empirical models almost always require scaling when used with satellite shock observations. Furthermore, the discrepancy in the size of the outer boundaries between empirical models and global MHD simulations is common, and could be attributed to a number of factors (Gordeev et al., 2015). However, it is generally thought that the relative dynamics are still applicable; more realistic inner magnetosphere models do not much improve the motion of the outer boundaries (Z. Němeček et al., 2011).

The MHD model performs well in capturing the overall motion of the bow shock and magnetopause, compared to the scaled empirical models. For example, figure 4.5a shows that large variations in the solar wind dynamic pressure occur starting at $t = 130$ min and $t = 190$ min. The sharp increase in dynamic pressure at $t = 130$ min causes the bow shock and magnetopause standoff positions to move Earthward, as expected. The pressure then reduces slightly, causing the shock and magnetopause to move Sunward. Both the empirical models and Gorgon show the boundaries responding on similar overall timescales. However, they do share common

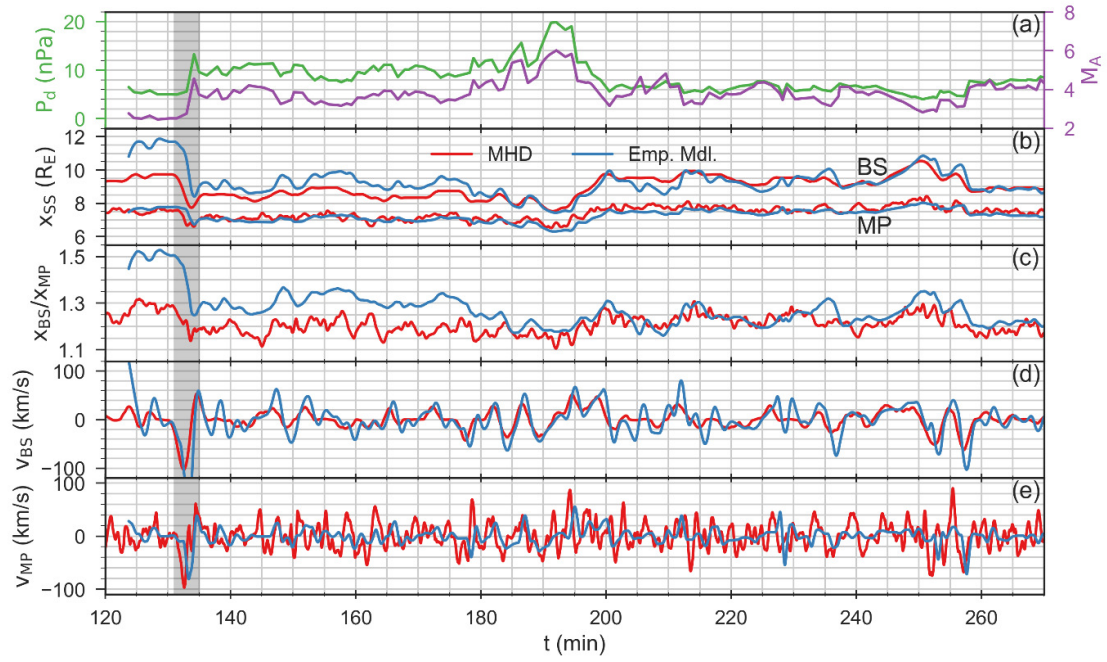


Figure 4.5: The variability of the Sun-Earth magnetopause and bow shock of the simulation (red), compared to the empirical models (Jeřáb et al. (2005) for bow shock, Shue et al. (1998) for magnetopause) (blue), showing: (a) the solar wind dynamic pressure (green) and the Alfvén Mach number (purple); (b) the magnetopause and bow shock positions; (c) the ratio between the the magnetopause and bow shock standoff distance; (d-e) the radial velocity of the bow shock and magnetopause. The grey shaded region indicates the time period analysed in figure 4.6.

features, such as the alternating magnetopause velocity peak at $t = 130$ min, corresponding to an Earthward movement, followed by a Sunward movement.

The drop in pressure at $t = 190$ min has more complex features. Leading up to the peak, from $t = 170$ min the dynamic pressure and the Mach number oscillate as they rise. This causes the bow shock and magnetopause to rapidly reposition, moving back and forth along the x_{GSE} direction. The bow shock velocities agree well, with similar peaks overlapping between Gorgon and model. During this interval, the empirical model bow shock velocities tend to be larger than the simulation velocities, which is a general trend throughout. However, the simulated magnetopause is much more variable than predicted by the Shue et al. (1998) empirical model. This is echoed in figure 4.5c which shows the ratio between the bow shock and magnetopause distances, and where there is more small-scale variability in the simulation result (red). Closer analysis of the simulated magnetopause shows that the magnetopause current sheet is often disrupted by the presence of reconnection jets and associated structure. However, this does not appear to affect the overall location of the shock, with the magnetosheath acting in some sense as a low-pass filter.

4.3.2 Overview of two-dimensional boundary dynamics

The interval $t = 131$ min to $t = 135$ min is now examined in more detail, when the bow shock moved Earthward and then Sunward over a period of 4 minutes in response to a sharp increase and small decrease in solar wind dynamic pressure (denoted by the grey region in figure 4.5). Figure 4.6(a-b) each show 8 snapshots of the dynamic pressure in the ecliptic and noon-midnight planes, with time increasing from left to right. Figure 4.6(c-d) show the positions of the shock in the ecliptic and noon-midnight planes at those times, calculated according to the procedure described in section 4.2.1. The colour of the line indicates the time. Figure 4.6e shows the corresponding position of the empirical model shock, which is cylindrically symmetric. Figure 4.5d immediately shows that the simulation bow shock shape has cusp-like features and is not a simple conic section in the noon-midnight plane, as assumed by many empirical models. Furthermore, as the bow shock moves, the position of the cusps do not stay in the same place. When comparing the shock normal between the simulation shock and the model shock, the simulation shock normal deviates on average by 8° from the model shock normal, with the majority of the deviation being within 20° : this is similar to the statistical results obtained in K. Meziane, Alrefay, and Hamza (2014).

As mentioned in the previous section, during this interval the solar wind dynamic pressure increased and then reduced. Figure 4.5(c-d) show that the sub-solar point is the first to move Earthward: this is expected as the solar wind front reaches this position on the shock first. However, the flanks of the shock during the first three snapshots (131.0 to 132.2 minutes) barely move from their original position. At intermediate times, 132.8 min to 134.0 min, the flanks are now moving anti-Sunward whereas the shock sub-solar point is approximately stationary and appears to reach equilibrium. The time from when the sub-solar point first moves and when the flank at 90° latitude moves is approximately 2 minutes, which corresponds to a speed along $-x_{\text{GSE}}$ of approximately 530 km/s, comparable to the solar wind speed. In the final two snapshots the sub-solar point moves Sunward. This is due to the decrease in solar wind dynamic pressure near the nose as can be seen in the change in colour from yellow to green in panels 5 to 8 of figure 4.6a and b. However, this change has not yet reached the flanks, which continue to move anti-Sunward. This behaviour is not seen in the model bow shock position (panel e), whose overall profile moves instantaneously and uniformly with the upstream solar wind conditions.

4.3.3 Shock Motion on the dayside flanks

To better understand the global motion of the shock, its variability and speed away from the Sun-Earth line is now examined. The radial position and velocity of the shock is focussed on as a first step. Using the method described earlier, the bow shock position is found in the Cartesian GSE coordinate system. To calculate the radial motion, the GSE position is first converted to a two-dimensional polar coordinate system centred on Earth, where r is the radial distance from the centre of the Earth, and ϕ is the azimuth from the positive x_{GSE} axis (defined such that $\tan \phi = y_{\text{GSE}}/x_{\text{GSE}}$). The location of the shock is then interpolated along lines extending radially from the origin. This gives the position of the shock as a function of ϕ . By doing this for each time-step and taking the gradient between each time-step, the radial velocity is obtained. Figure 4.7 shows the radial bow shock position along three radial lines: at $\phi = 0^\circ$ (the Sun-Earth line, shown in red), at $\phi = 40^\circ$ (blue) and at $\phi = 80^\circ$ (green).

The standoff positions (figure 4.7a) along each flank direction show similar behaviour to the Sun-Earth line. The flank shock positions have larger standoff positions, which is expected due to the overall shape of the shock (figure 4.5). Whilst the general pattern of the flank shock motion follows that of the sub-solar shock (with some delay, as noted above), it is more exaggerated,

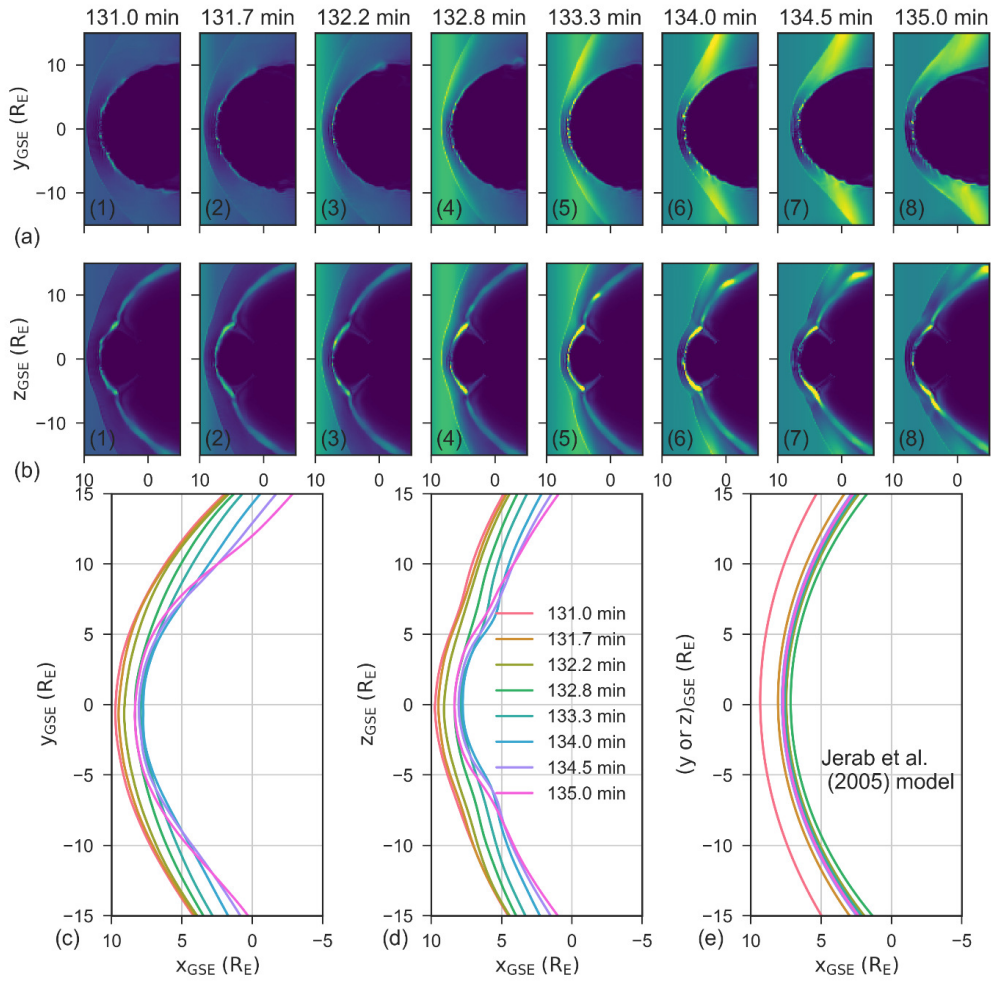


Figure 4.6: Example of bow shock motion due to changes in solar wind dynamic pressure. Panels (a) and (b) show the dynamic pressure of the plasma in the ecliptic and noon-midnight planes at successive time intervals, with the Sun off to the left (blue = low, green/yellow = high). Panels (c) and (d) show the bow shock position in the ecliptic and noon-midnight planes, at the same times as the snapshots shown in panels (a) and (b). Panel (e) shows the Jeřáb et al. (2005) model bow shock. As the solar wind pressure increases, the bow shock moves anti-Sunward, first at the Sun-Earth point, then later at the flanks. The reduction in solar wind pressure during the last two snapshots causes the sub-solar shock to moves Sunward whilst the flanks are still moving anti-Sunward. This complex motion is not captured by the empirical model.

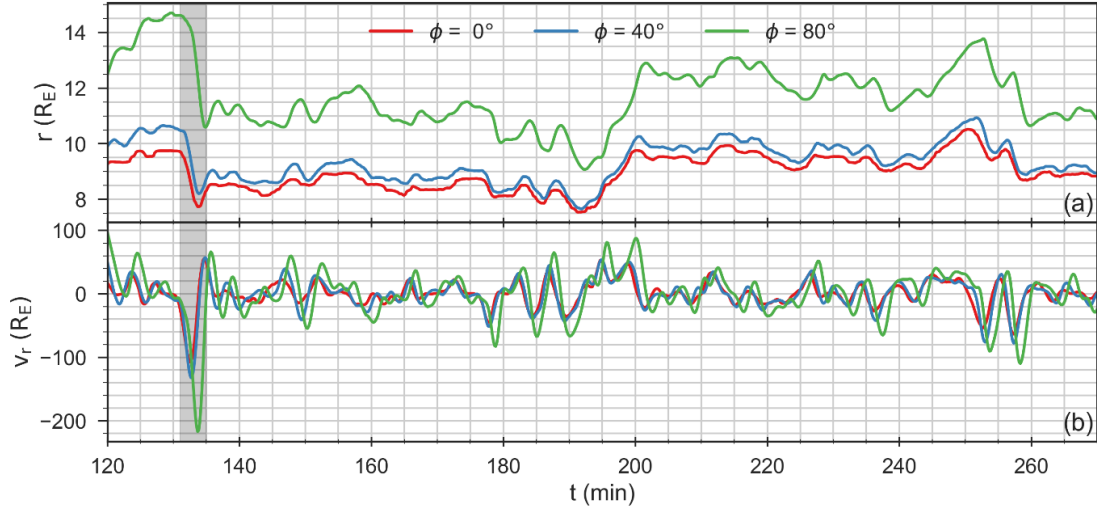


Figure 4.7: Out of plane bow shock position and radial velocity. (a) shows the radial position, (b) the radial velocity of the shock. The three lines denote the bow shock position at different angles from the Sun-Earth line in the ecliptic plane: red is the Sun-Earth line, blue is 40° duskward, green is 80° duskward. The grey shaded region shows an example where the flank shock velocity peak is much larger in magnitude than the sub-solar shock, and is delayed.

creating larger bow shock velocities. Hence it is found that the shock velocities on the flanks are in general much larger than at the Sun-Earth point. This can also be seen in figure 4.6.

Figure 4.7a also highlights the effects of the solar wind propagation delay on the flank shock location. This can best be seen in the velocity graph in figure 4.7b, comparing the 0° line (red) and 80° line (green): the 80° line trails the 0° line. One example of this is at $t \sim 130$ min, denoted as the grey shaded area in figure 4.7: this region corresponds to the same period discussed in section 3.2. The sub-solar shock initially moves Earthward due to the increase in solar wind dynamic pressure. The flanks also begin to move Earthward, a couple of minutes later than the sub-solar shock. However, by this time a new phase front with lower solar wind dynamic pressure has arrived at the sub-solar shock, which begins to move Sunward. Consequently, the sub-solar region is moving oppositely to the flanks towards the end of the shaded interval. Figure 4.6 therefore shows that the shock does not move coherently in and out at all locations in response to variations in solar wind pressure that occur on timescales less than the transit time of a phase front over the shock surface.

4.3.4 Histograms of Bow Shock Speed

By taking the radial shock velocity for all ϕ and t , one can produce a histogram showing the distribution of radial shock speeds. Figure 4.8 shows the distribution of the bow shock velocity in the ecliptic plane. The time-series data from the previous analysis from all radial lines is taken and binned in 10 km/s bins to create a histogram of the shock velocity distribution. A Gaussian profile is then fitted to the data.

In general Gaussian profiles fit the data well, especially for low velocities. At high velocities, there is a slight deviation with an enhancement in the tail of the distribution. Figure 4.7a shows that the radial shock velocity has a mean near zero, and a standard deviation of 21.4 km/s. The data were then divided into sub-solar and flank regions (figure 4.7b-c). There is a difference between velocities measured in the sub-solar region (figure 4.8b) to those in the flank regions (figure 4.8c). The flank regions have a larger standard deviation, which reflects the fact that the shock tends to move faster on the flanks than it does at the sub-solar point. These distributions are similar to those in K. Meziane, Alrefay, and Hamza (2014) who derived the bow shock velocity distribution using Cluster observations.

The mean of the shock speed in all three panels of figure 4.8 is not exactly zero, but approximately 1–2 km/s radially outward. This is likely a consequence of the non-steady solar wind input used, where the shock does not necessarily end where it starts in the analysis period, implying a mean displacement further away radially from the Earth. However, this result is not significant, as this mean is comparable to the Gaussian error on it, 1.5 km/s.

4.3.5 Virtual Spacecraft

Part of the variability between different empirical models can be ascribed to the fact that they are based on different spacecraft datasets acquired in different regions of space (Schwartz, 1998). To better understand this, virtual spacecraft are used to examine the simulation data, particularly focussing on the distribution of shock speeds. This is done with the following method.

A prototypical satellite orbit is first defined, using the apogee and perigee of Cluster's initial orbit (Escoubet, Fehringer, and Goldstein, 2001). The orbit is confined to the ecliptic plane, and its apse line orientation is chosen at random between $\phi = [-110^\circ \text{ to } 110^\circ]$. The starting position of the virtual spacecraft at $t = 120$ min of simulation time is set randomly. In order to

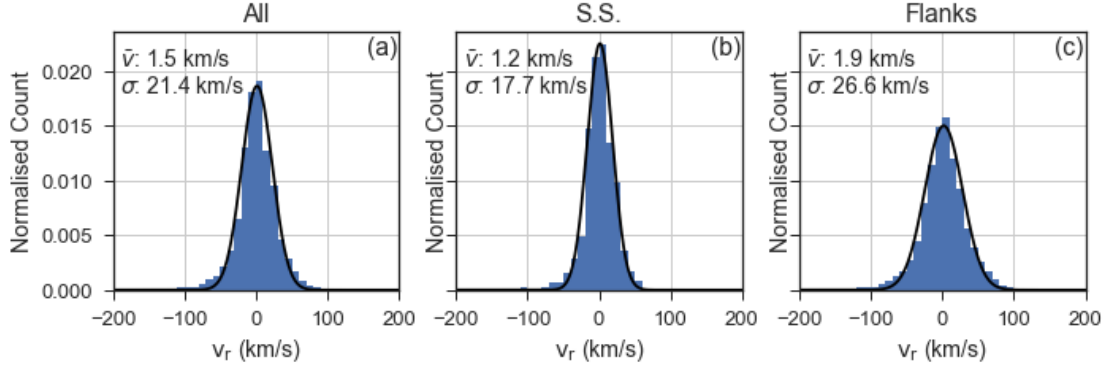


Figure 4.8: The radial bow shock velocity distribution in the ecliptic (xy) plane. Each panel shows how the distribution changes in different locations: (a) the whole shock; (b) the sub-solar region ($<45^\circ$ from x axis); (c) the flanks ($>45^\circ$ from x axis). The flank distribution has a larger standard distribution than the sub-solar distribution, since the shock is typically moving faster there.

find shock crossings, the spacecraft and bow shock positions as a function of time are compared, locating points of intersection. The shock velocity as observed by the virtual spacecraft is then found by sampling the shock velocity time-series at that position and time. This process was repeated for 201 virtual spacecraft, and resulted in a similar number of crossings to that of K. Meziane, Alrefay, and Hamza (2014). Unlike the four-spacecraft analysis, a timing analysis is not performed in this work, since the simulation provides complete spatial information of the shock and thus fewer assumptions need to be made. Figure 4.9a shows all the virtual spacecraft orbits as black dashed lines. Shock crossings are shown as red dots, with the corresponding shock surface also shown.

At the time of each shock crossing, the radial speed of the shock is also known from the analysis described in the previous sections. The histogram of radial shock speeds associated with the virtual spacecraft crossings is shown in figure 4.9b. The green curve is a Gaussian fit to the data in blue, and the black curve shows a similar fit for the velocity over the whole shock (the same Gaussian fit as in figure 4.8a). The virtual spacecraft radial shock speed distribution has a standard deviation of 43 km/s, much larger than that of the full simulation distribution. If the number of virtual spacecraft is increased (hence increasing the number of crossings) the distribution in figure 4.9 becomes smoother, but does not tend towards the global distribution (figure 4.8a). This suggests a fundamental bias to larger shock velocities in the virtual spacecraft analysis.

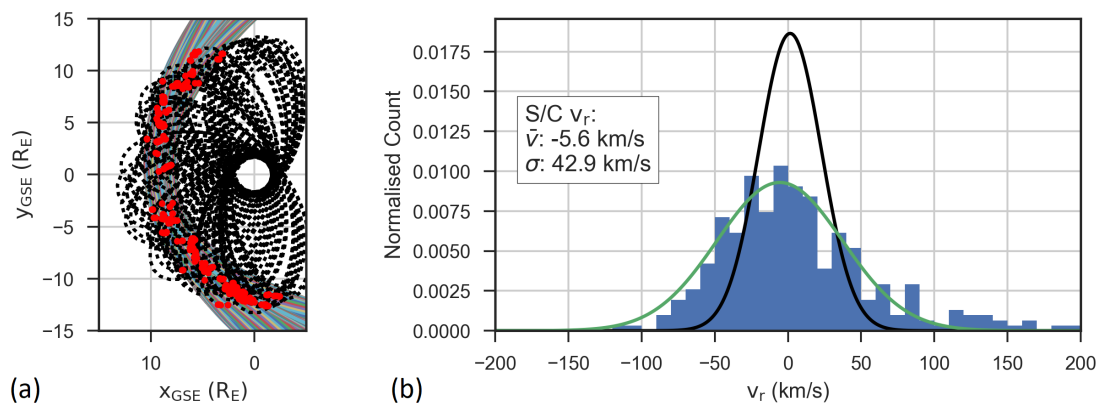


Figure 4.9: The bow shock velocity distribution as measured by virtual spacecraft. (a): shock crossings calculated using the virtual spacecraft algorithm. Black dashed lines show the randomly generated elliptical orbits, and red dots show where the virtual spacecraft crossed the shock. The corresponding shock surface is also shown. (b): The virtual spacecraft observed shock velocity distribution, together with a Gaussian fit (shown in green) and the Gaussian fit over the whole shock as in figure 4.8a. The distribution observed by the virtual spacecraft is much wider than that observed in the simulation as a whole.

4.4 Conclusion

In this chapter, Gorgon was used to simulate the magnetosphere during a real solar wind event (Maksimovic et al., 2003). The variability in the solar wind dynamic pressure causes the magnetosphere’s outer boundaries to vary in response. The bow shock position is found to be highly variable during this time: it expands and contracts largely due to changes in the solar wind dynamic pressure. There is more variability in the magnetopause location in the simulation, but this does not appear to affect the location and shape of the bow shock on a global level, suggesting a low-pass filter effect. The work in this chapter suggests two main implications on the wider magnetospheric research.

When using empirical models to derive information about the shape and motion of the bow shock, they must be used with care when used in spacecraft data analysis. In analysing the motion of the sub-solar region, it is found that the variability of the bow shock is comparable to that of the Jeřáb et al. (2005) empirical model. This offers an extra validation of the Gorgon model to that performed in chapter 2. Hence the Gorgon code can capture the expected variability of the shock, which, on average, have been found to agree with spacecraft observations of shock

positions and motion (K. Meziane, Alrefay, and Hamza, 2014; Meziane et al., 2015). However, there are a number of discrepancies when considering the motion of the shock in the flanks. Analysis of the simulation shows that for typical variability on time scales of minutes, the nose and flank can be moving differently at the same time. This is essentially due to the solar wind transit time over the shock in the anti-Sunward direction, with the sub-solar region affected by the varying solar wind several minutes before the flanks. Empirical shock models often assume a characteristic shape scaled by a single value of the upstream solar wind parameters. The results of the simulation indicate that such models must be used with care when the solar wind dynamic pressure is varying on minute timescales and when spacecraft observations are far downstream of the sub-solar point. Furthermore, this also adds to the long-standing discussion of whether the shock moves coherently in the x-direction or self-similarly (i.e. radially). However, from figure 4.6, this picture is too simplified. Since the shock is essentially a series of local phenomena, affected by local plasma parameters and waves, it can cause different parts of the shock to behave differently, due to the timing between solar wind fronts.

The results of the simulation also suggests a potential bias in spacecraft observations of the shock, which was mainly due to two factors. This was examined in more detail by flying a virtual spacecraft with a Cluster-like orbit through the simulation, where a much wider distribution in shock speed is found. This implied bias in spacecraft observations of the shock velocity is due to two factors. Firstly, majority of orbits take the spacecraft to the flanks of the shock (Merka, 2005), which has a wider distribution than at the sub-solar point (figure 4.8). Secondly, spacecraft move relatively slowly compared to the shock: this causes observations of the shock to be due to the shock moving over the spacecraft. It is therefore necessary to revisit experimental observations of shock motion and their statistics to better understand the true underlying nature of the shock motion in response to the variable solar wind.

This work can be extended in a number of ways. This work was restricted to the ecliptic and noon-midnight planes, which neglects the fact that the shock is not necessarily axisymmetric, with a more complex cross section which can deform the shock normal and the normal shock velocity. With a full three-dimensional shock surface, more general spacecraft orbits can also be considered, with the potential to compare with real spacecraft observations. Secondly, this methodology could be repeated for different solar wind conditions, which would allow for a better general understanding of the shock motion and potential spacecraft biases. This simulation was run with a relatively strong solar wind event, and hence features may be exaggerated: the large variability in solar wind parameters over small timescales (due to a high solar wind speed)

could cause faster than usual motions of the shock. By extending to three-dimensions, choosing a wider range of spacecraft orbits and applying this analysis to multiple different solar wind events, a more statistical analysis can be performed which should give an idea of the extent of the implications of this work.

In future work, one could use the output of the global MHD simulation to dynamically scale an empirical model to the local upstream solar wind conditions for position on the shock, which could account for the solar wind propagation time. MHD ‘empirical’ shock models generated from global MHD simulations do exist (e.g. I. H. Cairns and J. G. Lyon, 1996), but still do not account for the variability in the solar wind. This could then be used to piece together a more accurate representation of the shock surface. However, it may not fully account for the motion, as the bow shock moves differently at the flank than at the sub-solar region.

5 Flux Rope Dynamics on the Dayside Magnetopause

5.1 Introduction

Magnetic reconnection is an important process in space plasma physics. The process transfers magnetic energy into heat and kinetic energy of the plasma, hence plays an important role in energising plasma. It also allows solar wind entry into the magnetosphere, with global implications. When the IMF is southward (IMF $B_z < 0$), it allows for enhanced plasma entry into the magnetosphere by opening up magnetic field lines in what is called the Dungey Cycle (Dungey, 1961).

The basic picture is that reconnection initially occurs on the dayside magnetopause, where the solar wind plasma compresses the southward IMF onto the northward planetary field causing reconnection (Haerendel et al., 1978; G. Paschmann et al., 1979). This is shown in figure 5.1, field line 1: the blue solar wind field line reconnects with the green closed field line at the purple x-line. The planetary field emanating from the poles now is topologically connected into the solar wind, which transports it anti-Sunward, shown by the red field lines 2–7 in figure 5.1. Eventually, the field line enters the magnetotail, where it reconnects again (steps 7–8), closing the magnetic field line and accelerating plasma Earthward.

Dungey's original picture is approximately confined to the noon-midnight plane ($x_{\text{GSE}} - z_{\text{GSE}}$), with dayside reconnection occurring at the position of anti-parallel magnetic field, the sub-solar point. In three-dimensions, component reconnection occurs: the field on either side of the reconnection region need not be entirely anti-parallel, but rather can reconnect in the presence of a guide field (Yamada, Kulsrud, and Ji, 2010). There exists a region on the magnetopause surface where reconnection can occur, which is dependent the orientation of the IMF compared with the planetary magnetic field orientation.

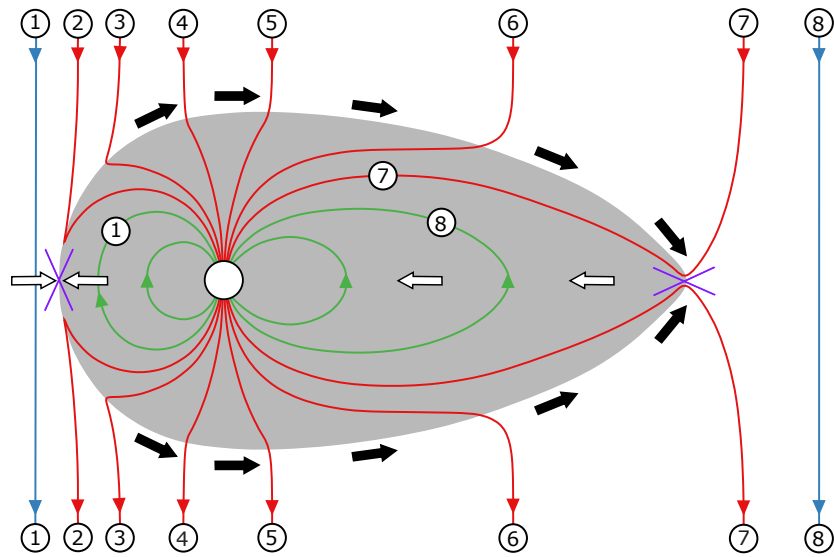


Figure 5.1: The Dungey cycle, showing the movement of field lines (grey arrows) and plasma convection direction (white arrows). Adapted from Baumjohann and Treumann (2012b)

The Dungey cycle also implies dayside reconnection is steady. However, even under steady solar wind conditions, observations and simulations have shown this not to be the case (C. Russell and Elphic, 1978; J. Raeder, 2006). When reconnection does occur on the dayside magnetopause, it can lead to the formation of Flux Transfer Events (FTEs). They are thought to transfer magnetic flux and hence energy into the magnetosphere when they are ‘open’: connected to either of the planet’s poles and to the IMF. Though the origin and evolution of FTEs are not fully understood, they have been observed to have a significant effect on magnetospheric dynamics and space weather.

FTEs are identified in spacecraft observations by their characteristic bipolar signature in magnetic field data, typically in the component normal to the magnetopause surface (C. Russell and Elphic, 1978; C. J. Farrugia et al., 2016). This signature may be indicative of rotational magnetic field structures called flux ropes: twisted magnetic field structures which can persist due to their force-free nature. Furthermore, they typically contain plasma features due to reconnection (hot and high velocity plasma populations), enhancements in the magnetic field strength (G. Paschmann et al., 1982; Fear et al., 2008), and force-free magnetic fields (C. J. Farrugia et al., 2016).

Multi-spacecraft missions such as the Magnetospheric MultiScale Mission (MMS), Time History

of Events and Macro-scale Interactions during Substorms (THEMIS), and Cluster have improved our understanding of FTE size and motion (Fear et al., 2005; Fear et al., 2008; C. J. Farrugia et al., 2016). In an FTE observed by MMS, it was found that there were two distinct electron populations within the FTE: one characteristic of originating from the magnetosphere, and the other from the magnetosheath, suggesting a single FTE can have a complex magnetic topology (Kacem et al., 2018).

With the majority of FTEs observed on the dayside, little is known about their evolution. The prevailing theory is that once reconnection forms an FTE, it is accelerated along the magnetopause surface towards the cusps by the $\vec{J} \times \vec{B}$ force (C. Russell and Elphic, 1978; Fedder et al., 2002). However, FTEs have been observed far into the magnetotail ($x \approx -67 R_E$), where their axis has been rotated from aligned in the azimuthal direction (parallel to the $x_{GSE} - y_{GSE}$ plane) to aligned with the z_{GSE} axis (J. P. Eastwood et al., 2012).

The formation of FTEs is thought to be inherent to the three-dimensional and time-dependent nature of reconnection along the magnetopause. Current theories can be generally categorised into three models (Fear et al., 2008). The first, proposed by C. Russell and Elphic (1978), describes an elbow shaped flux tube, as shown in figure 5.2a. They are formed by a reconnection patch on the magnetopause surface, creating a twisted magnetic field structure which lies on top of the magnetopause with the magnetosheath field draping over it. Southwood, C. Farrugia, and Saunders (1988) and Manfred Scholer (1988) proposes a model based on reconnection along a single x-line, where FTEs are formed by an increase in the reconnection rate, followed by a decrease. This is called bursty reconnection. The heating of the plasma due to reconnection causes expansion creating the bulge shown in figure 5.2f, which produces the typical bipolar field structure. Coupled with a shear velocity along the direction perpendicular to the reconnection region, it can create the twisted helical flux rope structure. The final model proposes FTEs are formed by reconnection occurring along multiple x-lines (Lee and Fu, 1985). In two-dimensions, this essentially creates a magnetic island, topologically connected neither to the IMF, nor the planetary field, best illustrated in figure 5.2d. In three-dimensions, there is usually a non-zero magnetic field component in the third dimension, creating a helical magnetic field structure, or flux rope, as seen in figure 5.2c.

These models are difficult to distinguish from each other using spacecraft observations, with each model predicting similar features (J. Raeder, 2006; Fear, S. E. Milan, and Oksavik, 2012). All three, for example, can produce the bipolar structure observed in spacecraft observations.

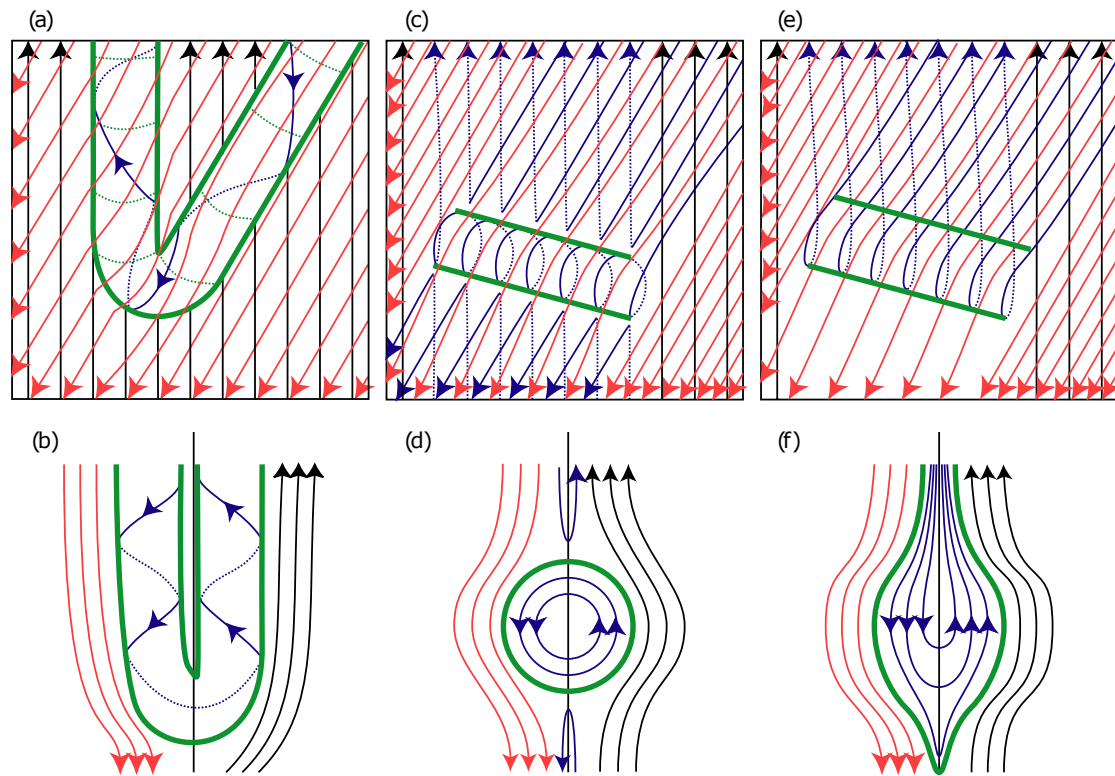


Figure 5.2: The three different FTE formation mechanisms: elbow (a, b), multiple x-line (c, d) and bursty reconnection (e, f). The top row is a view of the magnetopause surface along the sub-solar line. The bottom row is a side view, with magnetospheric field on the right, and IMF on the left. Blue lines denote solar wind field lines, red denotes closed field. The green lines show the boundaries of the flux rope. From Fear et al. (2008).

There are key differences however: the elbow shaped/patchy reconnection model does not form flux ropes which extend far into the azimuthal direction, whereas multiple x-line and bursty reconnection/single x-line do. Additionally, the multiple x-line model has the potential to create complex magnetic field topologies depending on where the x-lines are located, creating flux ropes that contain open field lines, closed closed field lines or IMF lines (Lee and Fu, 1985; Pu et al., 2013). Recently, there has been experimental evidence for multiple x-line reconnection, with MMS observing a flux rope structure in between two reconnection jets (M. Øieroset et al., 2011).

One of the main difficulties in understanding FTEs and their impact is due to the relative sparseness of spacecraft observations. Whilst auroral signatures of FTEs have been observed, they are rarely accompanied with a spacecraft in the correct position to observe the FTE structure. Though multi-spacecraft missions improve on measuring the local three-dimensional structure of FTEs, their coverage is not large enough to observe the whole flux rope at once. Hence, global simulations of the magnetosphere, which capture the magnetosphere as a whole, can be very useful in understanding a number of FTE features, giving insights into their formation, evolution, topology and impact on the magnetosphere.

A number of papers have used such simulations to study FTEs. Though it is known that MHD cannot capture the full physics required for reconnection, it has been shown that they do capture the location of reconnection in the dayside magnetopause relatively well (Komar, Fermo, and Cassak, 2015). J. Raeder (2006) uses OpenGGCM to investigate the effect of dipole tilt on the generation of FTEs. Under a strongly southward IMF, the FTEs were formed by multiple x-line reconnection, and was modulated by the dipole tilt. Cardoso et al. (2013)'s simulation of a steady, strong South- and duskward IMF, found elbow-shaped flux tubes. Other flux tubes were found to be interlinked with field lines of differing topology. Perez et al. (2018) studied a number of flux ropes spontaneously form under constant solar wind parameters. However, all of these simulations used idealised solar wind events, with parameters usually kept constant.

In this chapter, the Gorgon code is used to simulate a strongly southward IMF event as observed by ACE (Maksimovic et al., 2003), and study the FTEs formed. Using a global simulation, the full three-dimensional structure of the FTE can be obtained. This allows the study of FTE properties such as how it is generated, travels along the magnetopause, and its fate along the tail magnetopause. Furthermore, its magnetic topology can be calculated, from which the amount of flux transferred into the magnetosphere can be inferred.

5.2 Overview of the Event

The work in this chapter uses the same solar wind parameters as in Chapter 4. Whereas that chapter was focussed on the motion of the outer boundaries over the whole two hour period, but reduced to the ecliptic and noon-midnight plane, in this chapter the focus is on the formation of FTEs over a 20 minute timeframe. For this, the full three-dimensional state of the simulation is samples at 10 s intervals.

The solar wind input to the simulation is shown again in figure 5.3 for completeness (this is the same solar wind as shown in figure 4.2 on page 129). However, now the focus is on the analysis interval from $t = 190$ min to $t = 210$ min as is indicated by the grey shaded region. During this time, the solar wind number density drops from $30/\text{cm}^3$ to $10/\text{cm}^3$, as seen in panel (a). The solar wind speed stays constant at approximately 600 km/s (panel b). However, there is a significant v_y component, averaging at approximately 100 km/s. The strength of the IMF (panel c) was larger than typical IMF strength, ranging between 20 nT and 25 nT. The dominating IMF component is the B_z component, with B_z starting at approximately -30 nT, and reducing in magnitude slightly near the end. During this interval, the B_y component drops to approximately zero. With a negligibly small constant B_x , the IMF is predominantly southward, with no dawn-dusk component. The final panel (d) shows the sub-solar magnetopause and bow shock positions, as calculated by the methods in Chapter 4. As expected, the reduction of solar wind number density causes the magnetopause distance to increase, from approximately $6.5 R_E$ to $7.5 R_E$. This causes the bow shock position to also increase in distance. The magnetopause position is much more variable than the bow shock position, as was noted in Chapter 4. The speculation that this was due reconnection on the dayside will be further investigated in this chapter.

The number density profile in the ecliptic plane is shown in figure 5.4 at 1.5 minute intervals from $t = 190$ min to $t = 206.5$ min. The solar wind streaming in from the left of each subplot first passes through the bow shock, shown by the increase in mass density. It streams around the magnetopause, which shown as the sharp cut-off in mass density. The magnetopause and bow shock are approximately conical, but at an angle to the x_{GSE} line: this is due to the significant v_y component in the solar wind. At $t = 190$ min (panel 1), the magnetopause appears relatively calm: there are only a few small structures, shown by the bulges in the sharp cut-off. At this time, the solar wind number density is also at its highest, leading to the bow shock and magnetopause being most compressed, and a relatively thin magnetosheath. As time increases,

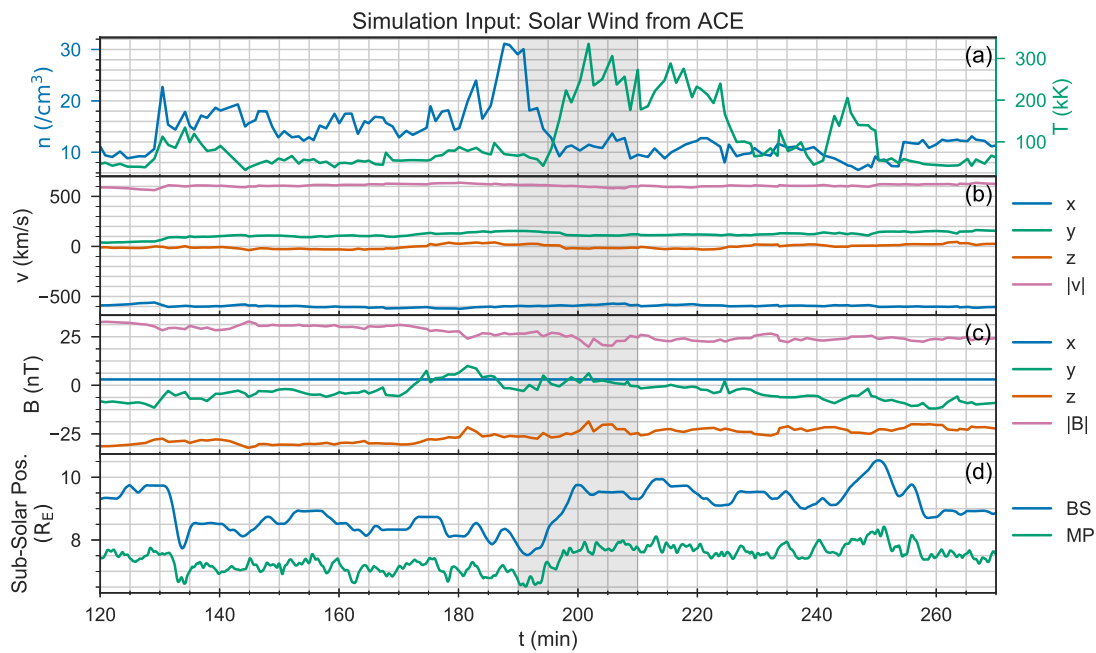


Figure 5.3: The solar wind input to the simulation, and the position of the bow shock and magnetopause in response. The shaded region denotes the region of interest during this chapter.

the solar wind number density decreases, as can be seen by the decreasing magnitude of the colour bar scales. This increases the magnetopause and bow shock standoff distances, and increases the magnetosheath width, which is most easily seen by comparing panels 1, 4, 7 and 10. As the number density decreases, more coherent structure forms on the magnetopause. On the dawn side, $y_{\text{GSE}} < 0$, there is a quasi-periodic oscillation of the magnetopause surface (indicated by the white dotted ellipse on figure 5.4.4) which persists throughout each panel in figure 5.4, with varying amplitude. On the dusk side, $y_{\text{GSE}} > 0$, distinct bulges in the plasma can be seen propagating down the flanks of the magnetopause, an example of one is shown in figure 5.4.6 by the white dotted ellipse. These bulges are much more sporadic than on the dawn side. An example of these bulges can be seen at $t = 193$ min, at position $x = 5 R_E$, $y = 10 R_E$, which propagates along the magnetopause surface before disappearing in the tail.

This same variability in the structure of the magnetosphere can also be seen in current density magnitude, as shown in figure 5.5. In general, figure 5.5, shows a number of current structures which persist throughout the time interval. As you approach the planet from the left, the first current sheet is due to the bow shock, where the shock tangential magnetic field component is compressed. Next is the magnetopause current layer, separating the magnetosheath plasma

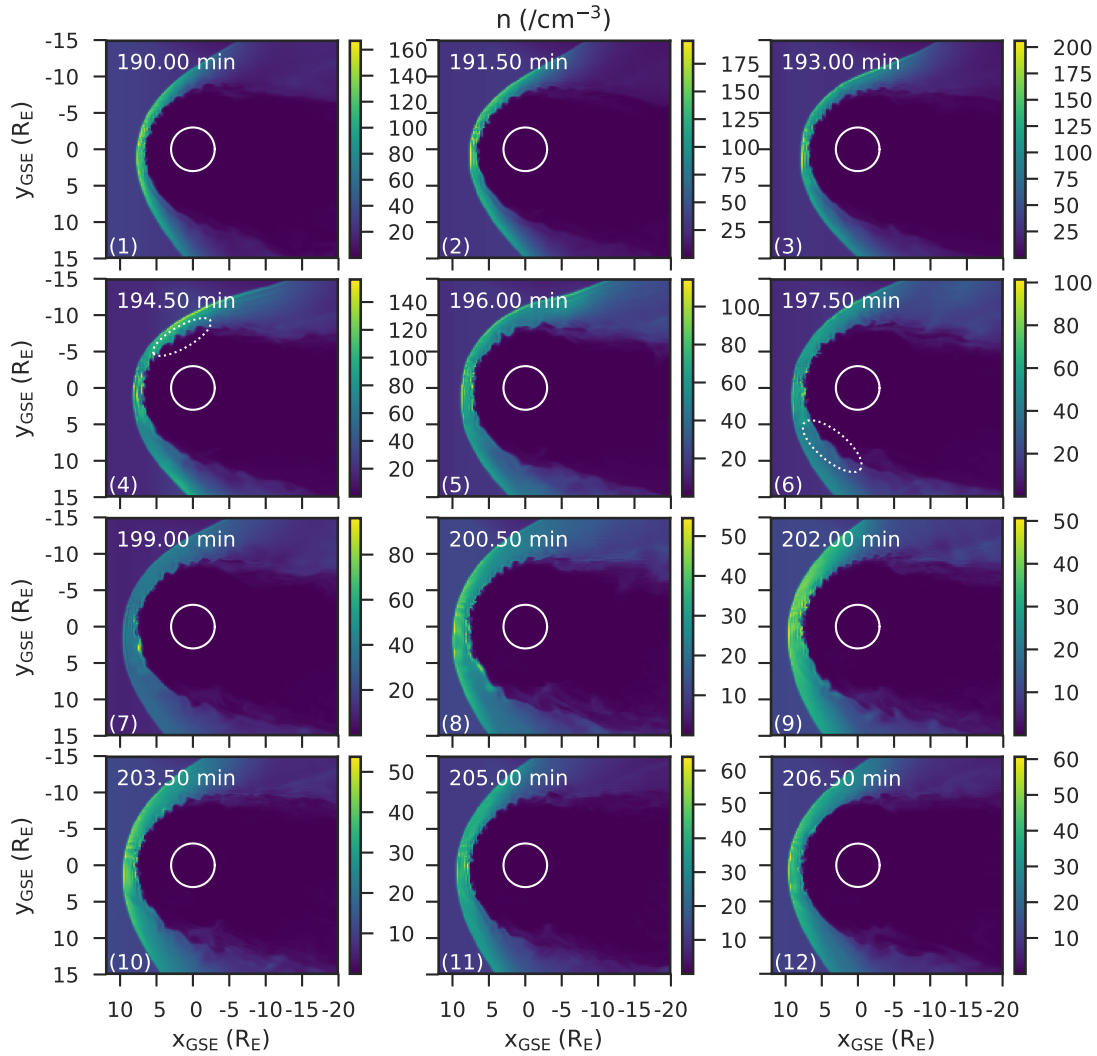


Figure 5.4: The number density in the ecliptic plane every 1.5 minutes from $t = 190$ min to $t = 206.5$ min. The number density colour bar scale changes every timestep. With the solar wind streaming from the left, the sharp increase in number density to the left denotes the bow shock. The brightly coloured region after it is the magnetosheath. The subsequent sharp decrease moving to the right is the magnetopause. The white circle is the inner boundary of the simulation, of radius $3 R_E$.

from the magnetospheric plasma. The magnetopause current density magnitude is largest in the sub-solar region, and decreases down the flanks. In the inner magnetosphere, there are small and sporadic currents right up until the inner boundary at $3 R_E$, denoted by the white circle. On the nightside, the inner magnetosphere current decreases to zero, creating an approximate semi-circle of very low current density. This is the closed field line boundary in the ecliptic plane. Further down into the tail, there are wisps of current density belonging to the plasma sheet and newly reconnected solar wind field lines.

Many of the structures seen in the number density (figure 5.4) can be seen here too. On the dawn side, the same oscillatory feature can be seen (indicated by the green circle on figure 5.44), as well as the bulges moving down the dusk side flank. More structure can be seen on the sub-solar region too. For example, at $t = 194.5$ min (panel 4), the magnetopause current sheet has local variations in current density strength.

Figure 5.4 and figure 5.5 shows that there are dynamic processes occurring on the magnetopause surface, which propagate down the flanks of the magnetosphere. These features are in fact caused by reconnection on the dayside forming flux ropes as is now described. The methodology used to study these flux ropes will be explored next.

Resistivity in MHD Simulations

Space plasmas are in general collisionless, meaning they have a very small resistivity and a broad applicability of the ‘frozen-in’ flux theorem. Under ideal MHD, a process such as reconnection should not occur, as the fields cannot disassociate from the plasma. In MHD simulations, there is a numerical resistivity which allows the magnetic field to reconnect. This is dependent on the size of the grid. Whilst Gorgon is a resistive MHD code, this resistivity has been set to the Spitzer resistivity, which is lower than the numerical resistivity. A common approach in reconnection studies in global MHD simulations is to apply an anomalous resistivity (Komar, Fermo, and Cassak, 2015), and has been found necessary when simulating substorms (J. Raeder et al., 2001). However, numerous simulations J. Raeder (2006), Cardoso et al. (2013), and Perez et al. (2018) have shown that flux ropes can form with numerical resistivity, and anomalous resistivity has no impact on dayside reconnection (J. Raeder, 2006).

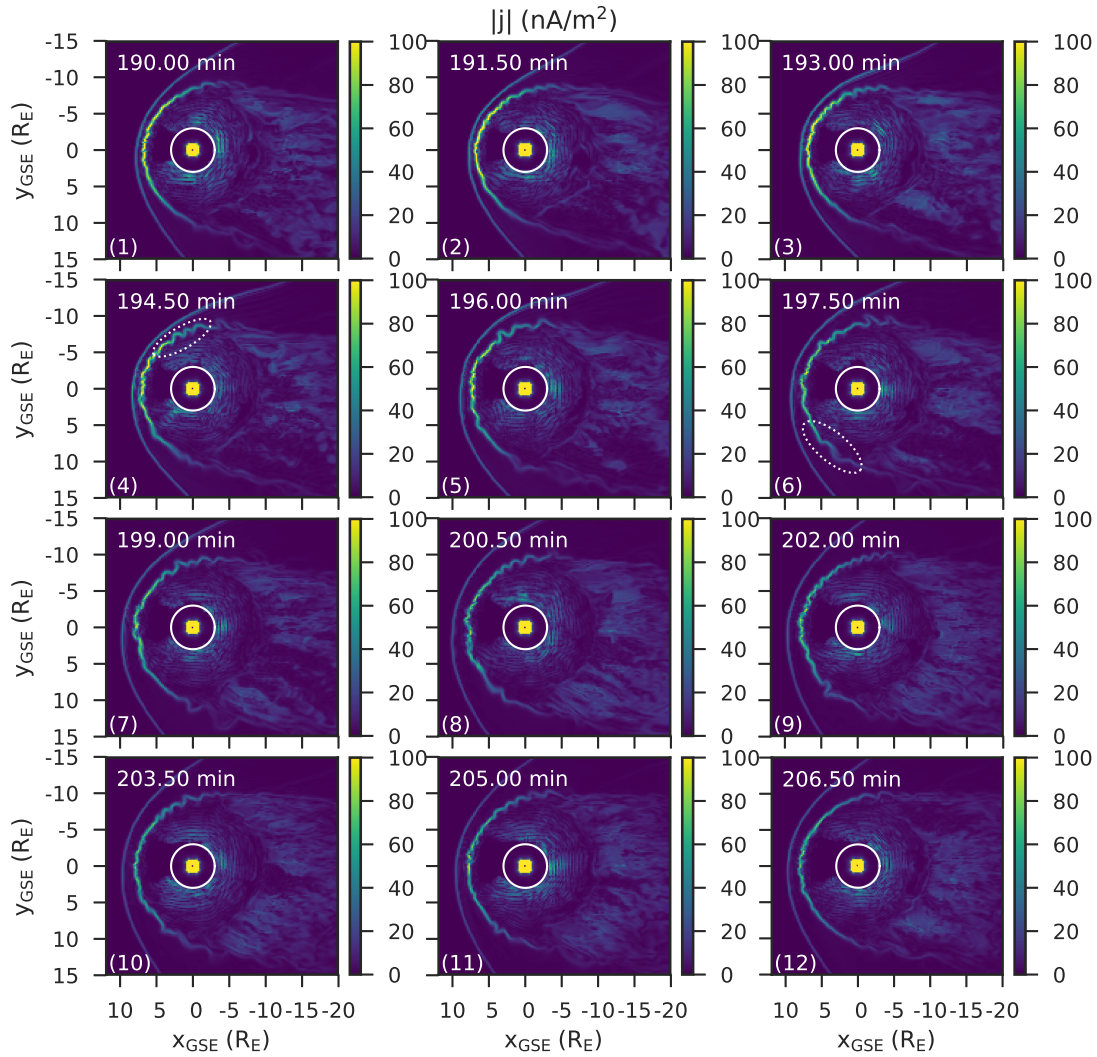


Figure 5.5: The magnitude of the current density in the ecliptic plane every 1.5 minutes from $t = 190$ min to $t = 206.5$ min. As with figure 5.4, the solar wind streams from the left. The compression of the magnetic field through the bow shock is shown as the left most current density peak. The magnetopause current sheet follows shortly after. Rather than a simple conical shape, it has oscillatory and transient structures. In the magnetotail, The white circle is the inner boundary of the simulation, of radius $3 R_E$. The subsequent sharp decrease moving to the right is the magnetopause.

5.3 Field Line Tracing and Topology Mapping

In order to visualize the magnetic field lines, the magnetic field streamlines are calculated. This allows the visualisation of the full shape and structure of each individual flux rope. For this, a new stream-tracer code has been developed. The main reason for this is malleability. Whilst visualisation software such as Paraview has in-built stream-tracers, they are difficult to manipulate. This work relies on the ability to adapt the stream-tracer code to include features such as topology mapping, filtering the streamlines, and binning them to reveal structures. An example of the latter case was in the fluopause finding algorithm in Chapter 3, where velocity streamlines were sorted and binned to estimate the position of the magnetopause. Stream-tracing is a fairly computationally intensive process, so the code makes use of parallelisation to make the computation feasible.

Stream-tracing is heavily dependent on its starting (seed) point. By under-sampling the magnetic field, the visualisation could miss vital structure and dynamics. Ideally, there would be multiple seed points for every grid cell to ensure that every flux rope of interest is captured. However, in typical simulation grid sizes, it takes a unfeasible length of time to calculate the streamlines for every point. Fortunately, the stream-tracer does not need to be run everywhere, but only at specific regions of interest. For example, regions in the solar wind will produce straight magnetic field lines, and regions close to the Earth will produce the dipole field lines. Hence the stream-tracer is run at specific regions of interest such as on the dayside or the nightside flank, as is shown in figure 5.6.

One important feature of FTEs is their topology, as this determines their impact on the magnetosphere. Given the full three-dimensional path of the field lines, their topology can be found by categorising the field line by their end point. In cases where both ends of the field line reaches the outer boundary, the field line is said to be part of the IMF, and is coloured red. When both ends of the field line ends up at the inner boundary, these are said to be closed field lines, coloured in blue. Open field lines occur when one end connects to the inner boundary, the other to the outer boundary. Open field lines are further classified by which pole they reach: North (magnetic South) is purple, and South (magnetic North) is green. The final type of field line present in the simulation belong to incomplete field lines, where one or both ends have not reached a simulation boundary. Computationally, these streamlines have a finite number of steps, so when the field line has not reached a boundary, it has likely run out of steps.

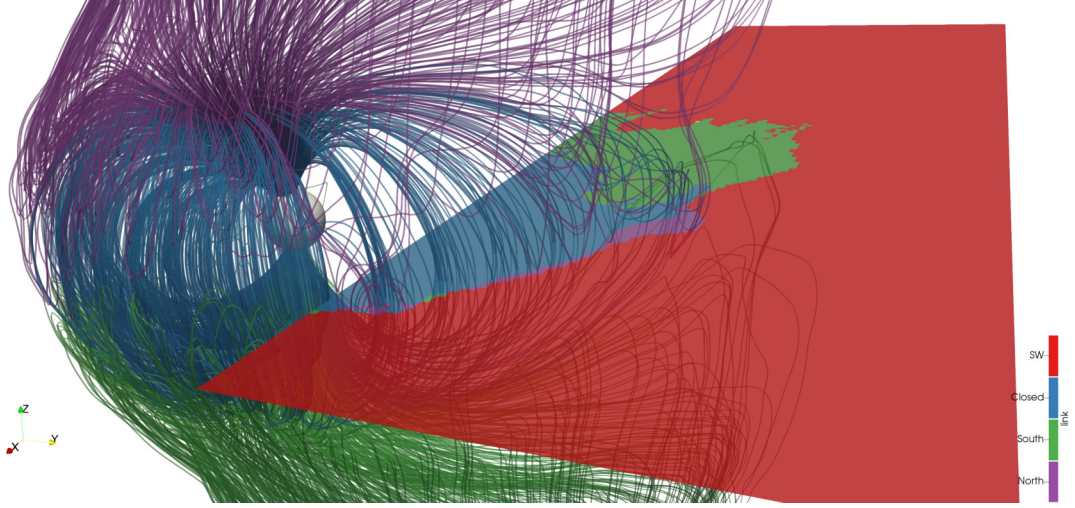


Figure 5.6: An example of field line tracing and topology mapping. The magnetic field lines are coloured by their topology: red for solar wind, blue for closed, green for open field lines connected to the South pole (magnetic North) and purple for open field lines connected to the North pole (magnetic South). The solar wind (red) field lines, which would dominate the view, have been filtered out. The slice shows the magnetic topology of cells along the ecliptic plane ($z_{\text{GSE}} = 0$): i.e. field lines connected to the red region are solar wind field lines, and map to the outer boundary of the simulation.

Another useful metric is to calculate the amount the field line twists, or rotates. This has been found to be useful in sifting out flux ropes from the background field lines. Their total rotation, Λ , is given by the angle between subsequent magnetic field directions along a field line,

$$\Lambda_s = \sum_i \cos^{-1} (\hat{\mathbf{B}}(s) \cdot \hat{\mathbf{B}}(s + \Delta s)), \quad (5.1)$$

where s is the path along the field line, Δs is the stream-tracer step size. This measure estimates the total amount a field line rotates. For example, a perfectly straight field line, such as those found in the solar wind, will have $\Lambda = 0$. Perfectly circular field lines completing one revolution will have $\Lambda = 2\pi = 360^\circ$.

The computation time can be decreased even further by removing the need to save every point in the field line and saving only the large-scale properties of each field line. A common use of this is in topology mapping, where, for every seed point of interest, the field line is calculated to determine where it is connected to. An example of this is shown by the slice shown in figure 5.6: a streamline is calculated from every position on that slice, and its topology is saved. This is then visualised on the slice.

By combining the topology mapping algorithm with calculations of the total rotation, Λ , regions where flux ropes may appear can be identified. Visualising only field lines which exhibit flux rope characteristics gives a clearer picture of the flux rope dynamics, whilst also finding all the flux ropes which meet the rotation criterion. The general method in plotting these flux ropes is as follows. For every timestep, the total rotation is calculated at cells in the simulation grid in the region of interest (e.g. the dayside). For each point where the Λ is large enough, usually $\Lambda > 8\pi$, field lines are drawn and their topology calculated. This produces the results seen in figure 5.7b.

An example of the usefulness of this filtering method is shown in figure 5.7. In panel a, the field lines are drawn using seeds spread on a uniform spherical grid. Whilst it does show some structure — for example, the magnetic separator can be inferred by the region separating the South (green) from the North (purple) field lines, it is difficult to see flux ropes, and the IMF had to be filtered out otherwise they would have dominated the view. Panel b shows the field lines filtered by the rotation method. Flux ropes can more more easily be seen, as straight field lines are not drawn. It also includes the IMF lines.

5.4 Results

5.4.1 The Dayside magnetopause

Using the flux rope drawing method described previously, magnetic field lines are examined on the dayside magnetopause. Figure 5.8 shows flux ropes coloured by their topology: red for IMF, blue for closed, purple for open and connected to the North pole, and green for open and connected to the South pole. The reconnection region, approximately denoted by the orange ellipse in figure 5.8.1, shows a complex entanglement of open, closed and IMF field lines, which can vary significantly over the 10 s cadence. Out of the tangled reconnection region, distinct flux ropes propagate out along the magnetopause surface, predominantly along the dawn and dusk directions. They are generated with the flux rope axis parallel to the azimuthal direction, as is expected. These flux ropes start off relatively small, approximately $2 R_E$ long in the azimuthal direction, and $\sim 1 R_E$ wide.

From the first timestep at $t = 192$ min (panel 5.8.1), two flux ropes have already formed (flux ropes 1 and 2), indicated by the black dashed ellipses. These are on opposite sides of the ecliptic

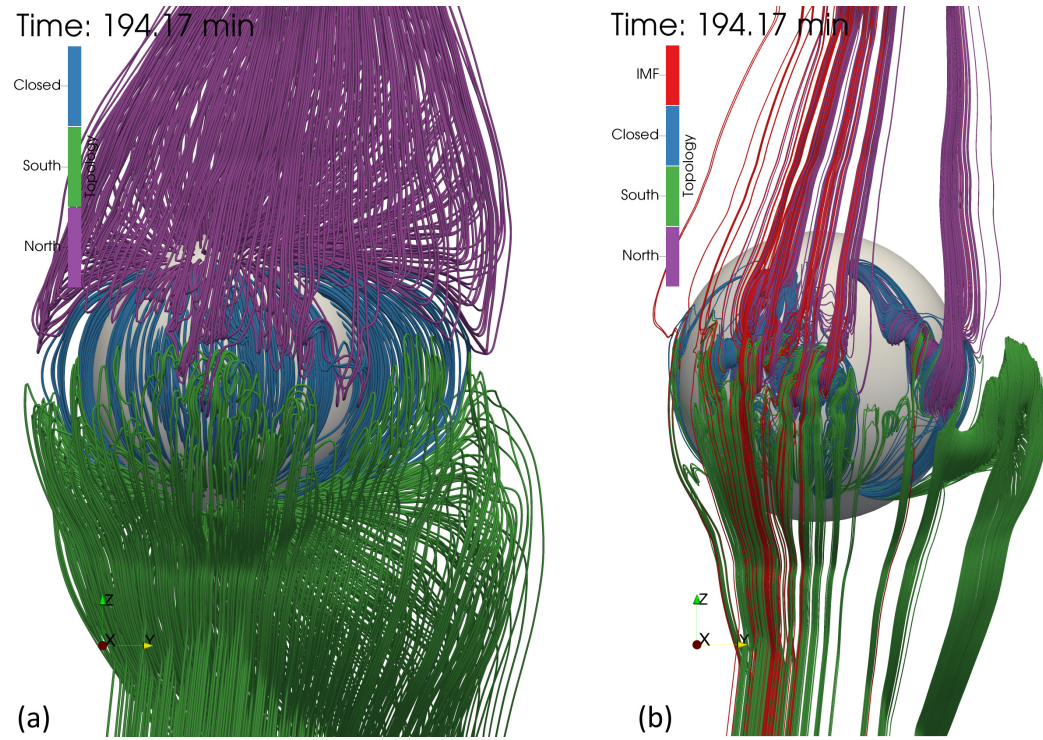


Figure 5.7: By filtering field lines by rotation, flux rope structure can more easily be seen. Panel a shows field lines drawn from a regular spherical grid. Panel b shows field lines drawn from positions on the simulation grid with a high enough rotation, $\Lambda > 8\pi$

plane: flux rope 1 is predominantly purple, indicating it is connected to the Northern pole, whilst flux rope 2 is predominantly green and is connected to the Southern pole. At $t < 192.33$ min (panels 1 to 3), both flux ropes are intertwined with closed field lines. This can best be seen in flux rope 2 in panel 2, which shows a blue strip wrapping around the green field lines. As time progresses, the closed field line region disappears and the flux ropes move around the dusk flank toward higher latitudes.

At a later time $t = 192.67$ min (panel 5) a large-scale magnetic structure emerges from the dayside reconnection region, labelled as flux rope 3. This flux rope exhibits a lot of complexity: it consists of a core of closed field lines (blue), around which open field lines are wrapped. The flux rope core magnetic field of closed magnetic field strength bends in a ‘U’ shape, which is due to a velocity shear in the z direction which kinks the flux rope. Eventually, the ‘U’ shape breaks, forming two flux ropes of different connectivity (figure 5.8.12). As it travels down the dusk flank, the South connected flux rope increases in azimuthal extent, whereas the North connected flux rope stays relatively small, and trails the South connected flux rope. Both flux ropes still have a core of closed magnetic field.

To make sense of the magnetopause reconnection region, the magnetopause surface is plotted in figure 5.9. In panel (a), a three-dimensional iso-contour at $B_z = 0$ denotes the magnetopause surface. It is coloured by the magnetic field strength normal to the iso-contour surface, with the black line showing where $v_z = 0$. The $v_z = 0$ line is highly warped. Panel (b) shows the magnetic topology, filtering out the IMF field. Regions of closed field lines (blue) denotes places where reconnection could occur, since these have direct contact with the IMF. These regions tend to be clustered together in multiple lines, suggesting the multiple x-line reconnection is occurring.

5.4.2 Flux Rope Evolution

Figure 5.10 shows a three-dimensional view on the dusk flank of the magnetopause. Two flux ropes are tracked, flux rope 3 and 4, which due to their differing topology evolve differently as they travel down the flanks.

Flux rope 3, first seen in figure 5.8.5, moves down the flank for approximately 6 min until it begins to dissipate at $t \sim 198$ min between figure 5.10.5 and 5.10.6 reaches the tail reconnection region. As it travels down the flank, it increases in size along the flux rope axis. The axis of flux

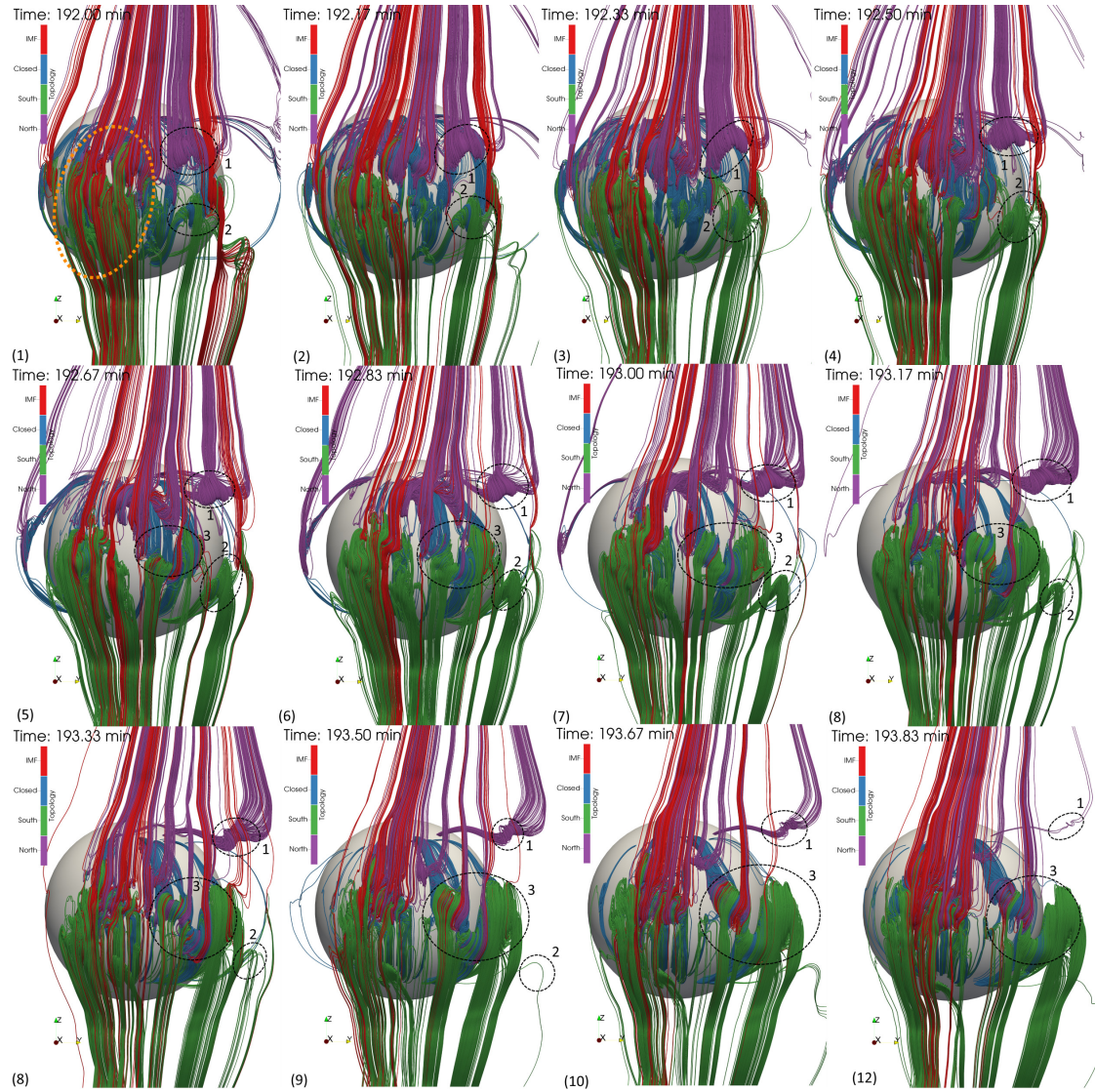


Figure 5.8: Flux Ropes generated on the dayside magnetopause. With each panel, time increases by 10 s. Magnetic field lines are filtered by their total rotation, showing only flux rope like structures on the magnetopause surface. Behind the flux ropes is a $6 R_E$ reference sphere. The dayside reconnection region (shown in orange in panel 1), creates a complicated mix of flux ropes. Out of this, a number of distinct flux rope like structures emerge, as denoted by the ellipses marked 1, 2 and 3, which travel down the dusk flank.

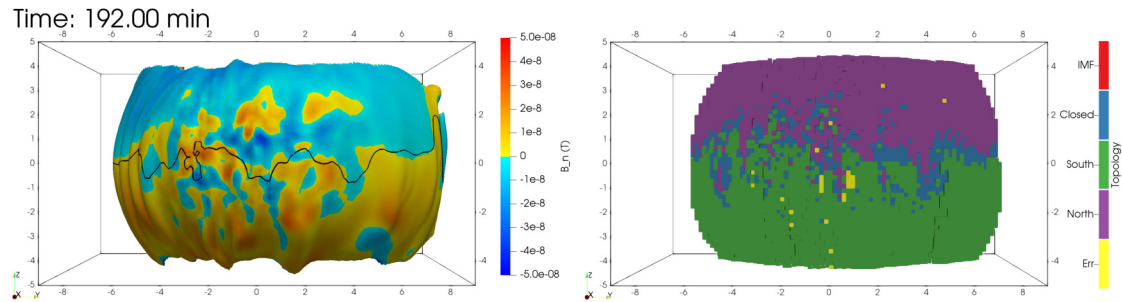


Figure 5.9: A view of the magnetopause from the magnetosheath, showing its magnetic field properties and plasma motion, and potential regions of reconnection and plasma. The left panel (a) shows the $B_z = 0$ contour, coloured by the magnetic field strength normal to the surface, B_n . The black line denotes the $v_z = 0$ line: above the line, $v_z > 0$, below $v_z < 0$. The right panel (b) shows the magnetic topology of each simulation cell. IMF field lines (red) have been filtered out.

rope 3 for both North- and South-open sections remains parallel to the ecliptic plane throughout the simulation.

Flux rope 4 can be seen being generated in figure 5.10.3, and contains a mix of all types of topology: closed, IMF, North- and South-open. Unlike flux rope 3, the two regions of opposite open topology do not separate, but remain as one flux rope. As flux rope 4 travels down the flank, its axis rotates from being in the ecliptic plane, to parallel to the z axis. It also increases in size along the flux rope axis. At the end, figure 5.10.8, flux rope 4 is seen to only consist of IMF, and becomes less twisted.

Both flux ropes 3 and 4 behave differently as they travel down the flank due to their differing magnetic topology. However, both flux ropes dissipate shortly after reaching the nightside reconnection region. This can be seen in figure 5.11, which shows slices of the simulation in the ecliptic plane as flux ropes 3 and 4 reach the tail reconnection region, with each row increasing in time by 30 s. The first column (a) shows current density magnitude, the middle (b) shows the total rotation of the magnetic field, and the right (c) shows the magnetic topology of the slice. In the topology plots (column c), the closed field line region can be seen by the blue region. The green region, connected to the South pole, essentially shows the tail reconnection region. The flank magnetopause is the region duskward of the blue and green regions, which coincides with the current sheet shown in column (a). Flux ropes 3 and 4 are shown as islands of purple, North connected, magnetic field lines, marked with circles.

Row 1 shows flux rope 4 just approaching the tail reconnection region, where blue meets green

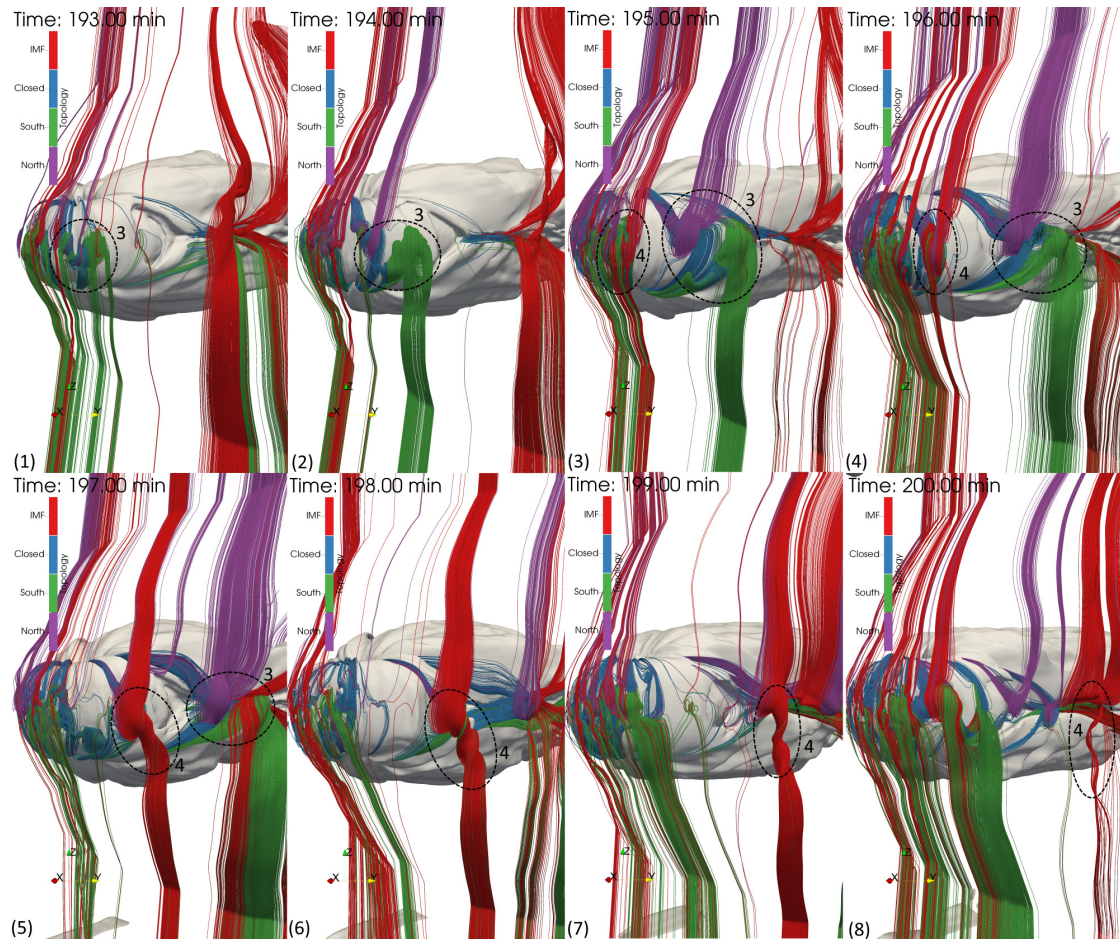


Figure 5.10: Flux ropes travelling down the dusk flank. With each panel, time increases by 1 min. Magnetic field lines are filtered by their total rotation, showing only flux rope like structures on the magnetopause surface. To filter out the inner magnetosphere, a number density contour of $1/\text{cm}^3$ is shown. The dayside reconnection region (shown in orange in panel 1), creates a complicated mix of flux ropes. Flux ropes 3 and 4 are tracked as they propagate down the dusk flank.

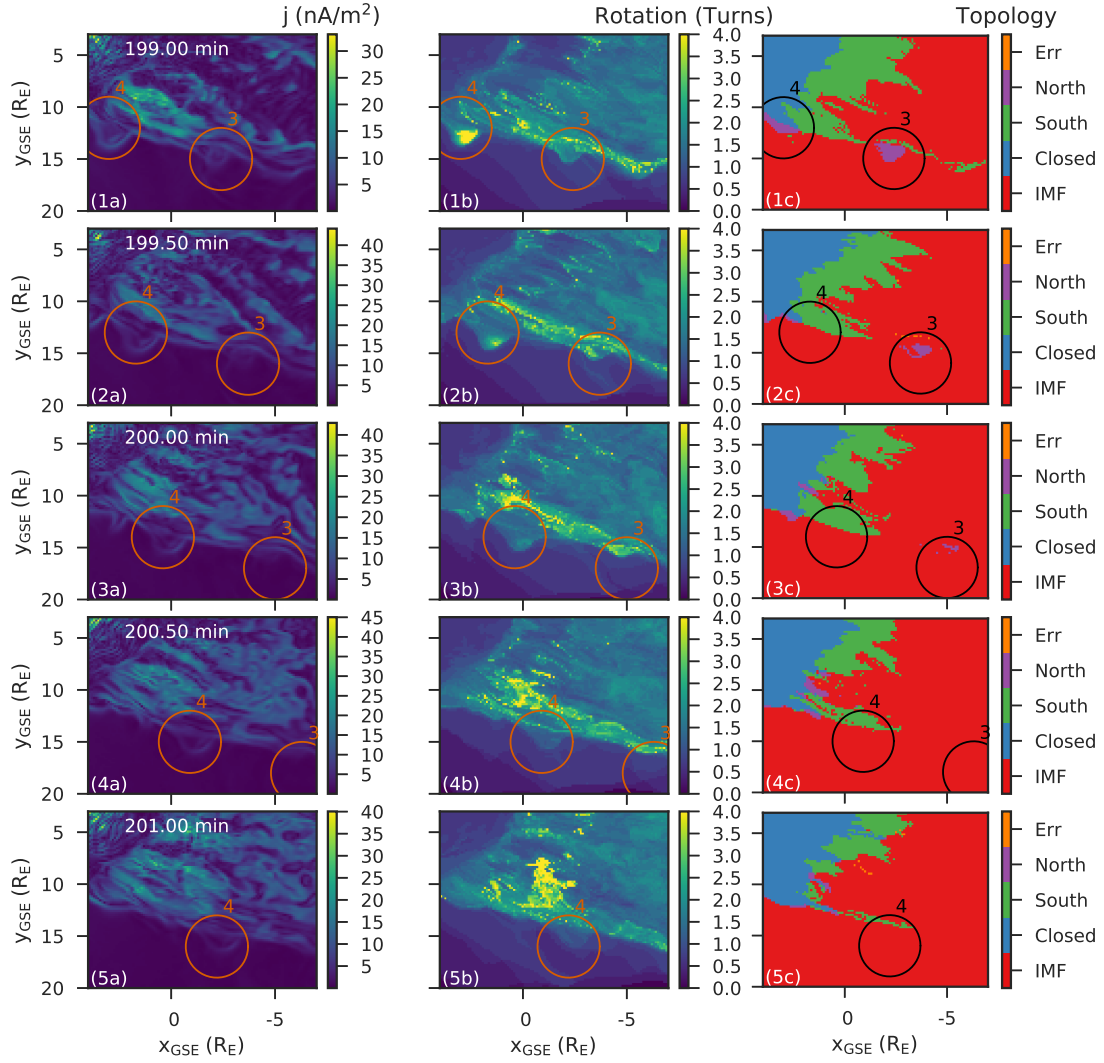


Figure 5.11: The dissipation of flux ropes past the nightside reconnection region. Each panel shows a slice in the ecliptic plane $z_{GSE} = 0$. With each row, time increases by 30 s. The left column (a) shows current density magnitude, the middle (b) the total rotation of the magnetic field, and the right (c) the magnetic topology.

(panel 1c). 30 s later, in row 2, the purple region has disappeared replaced by red. However, the flux rope is still there, as seen by the structure in the current (2a), as the high amount of rotation (2b). As time increases (rows 3 to 5), flux rope 4 remains connected to the IMF, but continues to be twisted, as shown in column b by its enhanced total rotation. Its total rotation does decrease with time, suggesting the flux rope is unravelling.

A similar process occurs with flux rope 3, except the open field line region persists long after the reconnection region. In the first timestep shown in figure 5.11, the purple, open field region has already passed the reconnection region. In subsequent time-steps, the purple region appears to dissipate, converting to red (IMF). This is due to the same reason as flux rope 4: reconnection occurs in the tail reconnection region, upstream of the flux rope rather than at the flux rope.

This process is shown in figure 5.12. Viewing the dusk magnetotail from the $+x_{\text{GSE}}, +y_{\text{GSE}}, +z_{\text{GSE}}$ direction, the field lines from flux rope 3 are shown, as well as closed field lines near Earth. These field lines have not been filtered by total rotation. Flux rope 3 is manually tracked through the simulation and seed points chosen which encompass flux rope 3, denoted by the circles labelled 3 in figure 5.11. In order to see the open field lines in the flux rope, the IMF field lines (red) start off transparent, but become opaque at $t = 200$ min when the flux rope contains no more open field lines. The closed field lines are generated from a grid in the $y_{\text{GSE}} - z_{\text{GSE}}$ plane, at $x = -5 R_E$, and only show the closed field lines (open and IMF field lines have been filtered out).

In figure 5.12, two main field line regions are shown: the flux rope and the closed field line region. In between there are elongated field lines with a large B_x component, containing both open (green and purple), closed (blue) and IMF (red) field lines. As time progresses from $t = 199.0$ min to 200.5 min, the number of field lines in this region decreases, making the region thinner. At the same time, the amount of open field lines in the flux rope is decreasing. This indicates the flux rope is changing topology, reconnecting in between the flux rope and the closed field line region. At $t = 200.5$ min, the flux rope detaches completely from the planet, becoming red. However, they are still significantly elongated, indicating the reconnection site occurred far from the actual flux rope.

Figure 5.13 shows the same field lines and sphere as figure 5.12, except they are coloured by the plasma v_x and the view is onto dusk magnetotail from the $-x_{\text{GSE}}, -y_{\text{GSE}}, +z_{\text{GSE}}$ direction. At the earlier time, $t = 199.0$ min, the flux rope is still attached to the planet. However, there is an enhanced v_x component in the $-x_{\text{GSE}}$ direction on the field lines in ellipse 1, and similarly,

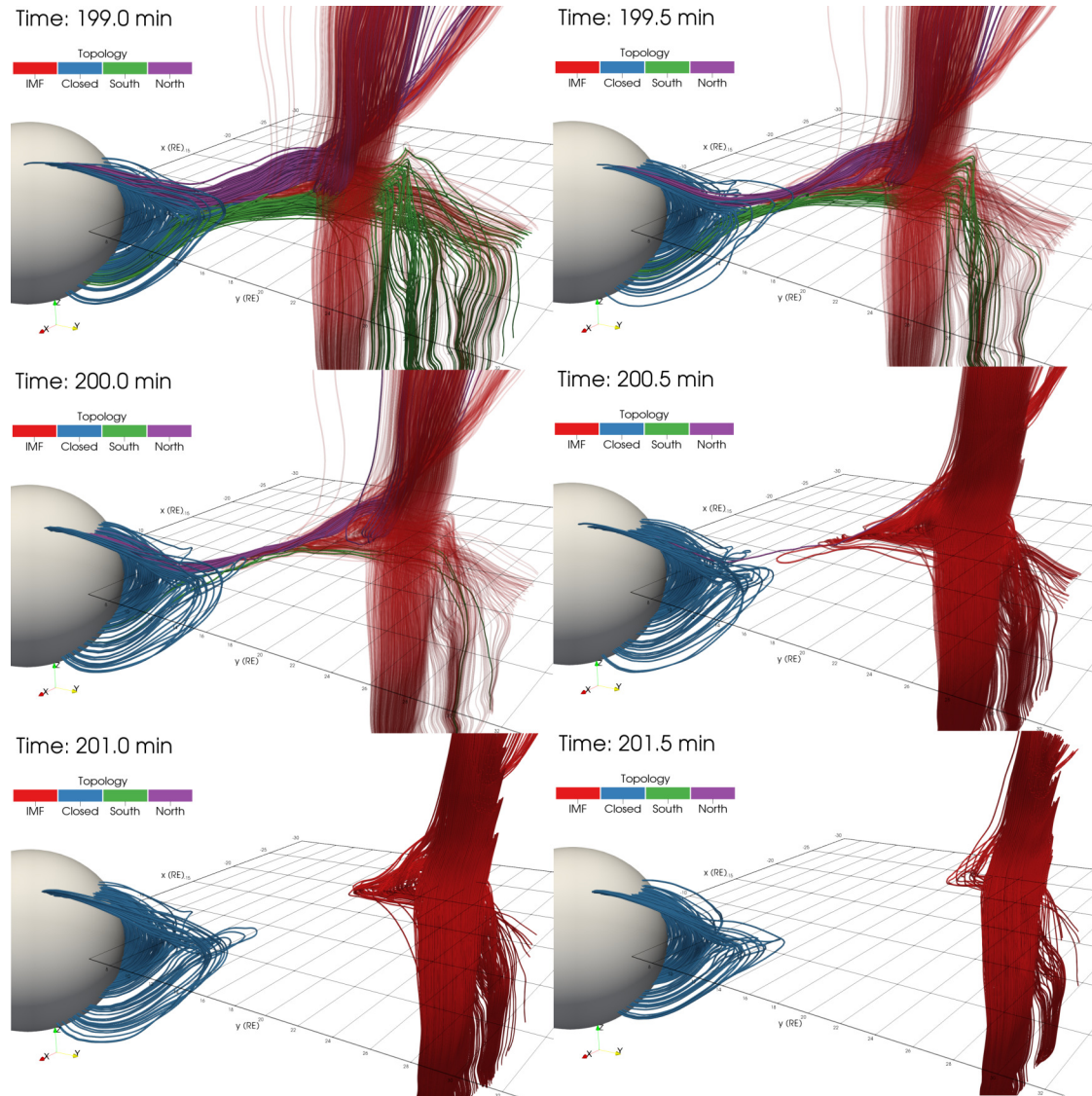


Figure 5.12: The evolution of flux rope 3 past the nightside reconnection region. The field lines have been drawn so only flux rope 3 is seen, along with the associated closed field line region. Each panel shows a grid in the ecliptic plane $z_{\text{GSE}} = 0$. With each row, time increases by 30 s. The left column (a) shows current density magnitude, the middle (b) the total rotation of the magnetic field, and the right (c) the magnetic topology.

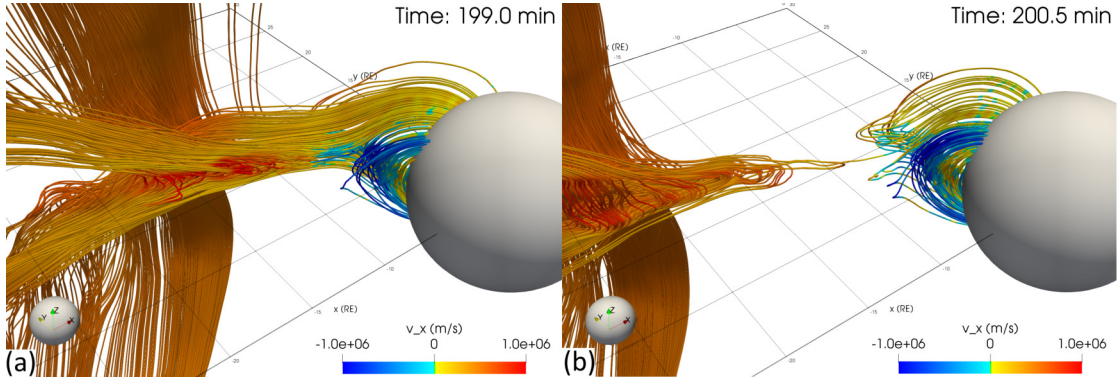


Figure 5.13: The dissipation of flux ropes past the nightside reconnection region. Each panel shows a slice in the ecliptic plane $z_{\text{GSE}} = 0$. With each row, time increases by 30 s. The left column (a) shows current density magnitude, the middle (b) the total rotation of the magnetic field, and the right (c) the magnetic topology.

enhanced positive v_x on the closed field line region in ellipse 2. This suggests that reconnection is occurring, accelerating the plasma via the $\vec{J} \times \vec{B}$ force. At the later time, $t = 200.5$ min, the flux rope has completely reconnected, losing connectivity with the planet, with the same signature in v_x as the earlier time.

5.5 Conclusion

This chapter is aimed at investigating the dynamics on the dayside magnetopause. Using the same solar wind event, and same simulation in Chapter 4, dynamic features were observed in the ecliptic plane. Large bulges of the magnetopause current sheet were observed, which were found to be flux ropes.

The flux ropes are formed on the dayside magnetopause surface, where the strongly southward IMF reconnects with the magnetospheric field. Overlapping the highly warped magnetopause current layer, there exists a complex interlaced set of field lines whose topologies vary from closed to open to completely IMF. This is due to time-dependent reconnection along multiple x-lines on the magnetopause. Out of this complex interwoven dayside field region emerge stand-alone flux ropes which propagate along the magnetopause surface. However, the latitude of their path is dependent on their topology. Two types of flux rope topologies are identified: completely open and partially open.

Flux ropes which are completely open — they are solely connected to a single pole (North/purple or South/green) — tend to move along the magnetopause at high latitudes. These flux ropes would be the equivalent to the classic flux rope evolution picture where they move towards the cusps and then dissipate. No flux ropes have been observed to go directly over the cusps. The exact reason for this requires further work, though there are two possible reasons. The two ‘cusp-like’ flux ropes, which passed over the magnetopause at high latitudes, contained some closed field line topology up until shortly after their generation. This is what causes partially open field lines to keep to the equator, so these ‘cusp-like’ flux ropes could have moved azimuthally initially, before their closed field line region disappears, allowing them to move at higher latitudes. Another reason is that these flux ropes do not exhibit a large enough total rotation, causing flux rope filtering algorithm to remove these from the results. Finally, given the high solar wind parameters, the magnetopause was significantly compressed, causing the travel time of an flux rope from the reconnection region to the cusp to be shorter than the 10 s cadence.

The second type are partially open. These flux ropes contain both closed field lines and open field lines connected to both the North and South poles. Due to this topology, they tend to stay along the ecliptic plane at low latitudes. In the classic cusp flux rope picture, the $\vec{J} \times \vec{B}$ force acts to move the flux rope over the poles. Since this flux rope is connected to both poles, the two opposite curvature forces cancel out, carrying it along ecliptic plane in the direction of the magnetosheath plasma flow. They were still found to be force free, with the thermal pressure balancing the $\vec{J} \times \vec{B}$ force.

The interwoven nature of these flux ropes suggest a number of implications on their impact on the magnetosphere and the Dungey cycle. FTEs have been shown previously to create signatures observed in the ionosphere (e.g. Sandholt et al. (1986), S E Milan et al. (1999), Wild et al. (2001), and Fear et al. (2009)). Hence, these flux ropes containing magnetic field from both poles would be topologically connected to both poles, hence there may have signatures matching signatures in both hemispheres of the ionosphere. However, flux ropes with a proportion of closed field lines complicate this picture as this magnetic flux and plasma population is already a part of the magnetosphere. This has implications of the amount of flux these flux ropes transfer into the magnetosphere. Without being able to verify the topology of FTEs in spacecraft data, estimates of the amount of open flux and plasma they contain may be an overestimation.

This leads to a number of questions that should be answered in future work. The question of how

these flux ropes form is of interest, whether they are special cases due to that particular solar wind event, or occur in more standard solar wind conditions. Secondly, the amount of magnetic flux and magnetosheath plasma these partially open flux ropes transfer into the magnetosphere needs to be determined, as they will contribute to the question of how flux ropes impact the Dungey cycle. This type of analysis should be performed for multiple different solar wind conditions, to deduce when and where different types of flux ropes are generated, and assess their impact on the magnetosphere through statistics on the amount of flux they transfer and the ionospheric impact. For estimating the ionospheric impact, the an ionosphere module such as those discussed in section 1.4.2 must be implemented in Gorgon, which allow determination of the ionospheric impact.

There is also a general question of whether these different flux ropes can be distinguished using spacecraft observations. Comparing the simulation flux ropes with multi-spacecraft observations of flux ropes (e.g. MMS, THEMIS and Cluster) can be particularly insightful. For example, in recent research, Kacem et al. (2018) was able to use MMS to determine the topology by comparing electron populations. Ideally, one would use a kinetic or hybrid code, such as Vlasiator (Alfthan et al., 2014; Pfau-Kempf et al., 2016), to simulate such events, which can capture both the global scale of an FTE including the physics required for reconnection. However, running such a code requires too many resources to render the computation feasible. Though MHD does not capture the full physics, it can shed light on the large scale structure and topology of these flux ropes and the impact such an event has on the magnetosphere.

There is also the potential for this work to be used in conjunction with the Solar-wind Magnetosphere Ionosphere Link Explorer (SMILE) (Raab et al., 2016), a spacecraft designed to observe the soft X-ray emission from solar wind charge exchange, whilst simultaneously observing the UV auroral emission. This allows it to directly image the impact of plasma entry into the magnetosphere, and its subsequent effect on the ionosphere. However, this type of imaging will only measure the volume integrated X-ray emission, calling for the need to be able to distinguish the physics of interest, solar wind entry into the magnetosphere, from the background. For this, global MHD simulations are of vital use in quantifying this background emission, as they have the full three-dimensional spatial distribution of plasma in the magnetosphere.

6 Conclusions

The solar wind – magnetosphere interaction is very complex and dynamic, spanning over large temporal and spatial scales. It consists of a set of regions with different dominating physical processes which are coupled non-linearly. Studying this interaction is important in both furthering our understanding of space plasma physics, and mitigating the effects of space weather, which has the potential to cause a vast amount of damage to vital infrastructure. A large body of work in this field uses spacecraft observations. Whilst they accurately observe the phenomena present in the magnetosphere system, they only provide local observations, and are relatively sparse compared to the size of the magnetosphere. Hence, global simulations and models play an important role in placing these observations in the global context.

This thesis is based on the adaptation of the Gorgon MHD code to simulate planetary magnetospheres. In chapter 2, the code is discussed in terms of its numerical solvers, magnetosphere implementation and verification against empirical models, and briefly compared to existing GMHD codes. Gorgon solves the resistive MHD equations on a three-dimensional uniform Cartesian staggered grid. In order to test the validity of the MHD solvers, a number of standard MHD test problems are performed. Each of these tests are designed to evaluate the MHD code’s performance in different regimes, and ultimately show that Gorgon is capable of simulating MHD and is comparable to other codes.

The magnetosphere implementation follows a similar set-up to the existing GMHD codes, in regards to the outer boundaries. Where it differs is the inner boundary: where the magnetosphere meets the dipole magnetic field of the planet. Rather than setting a fixed inner boundary, whose plasma parameters are set by an ionosphere module, Gorgon uses its vacuum solver to propagate the dipole field from a smaller radius, typically $\sim 1 R_E$, to the inner plasma boundary at $\sim 3 - 4 R_E$. This creates a resistive inner boundary, which closes the currents in the magnetosphere. Initialisation of the magnetosphere differs too: the simulation box is empty, allowing the dipole field to expand into the box. Simultaneously, the solar wind and

IMF are propagated inward from the Sunward boundary. To determine Gorgon's validity in simulating the magnetosphere, it is tested against empirical models using the Gordeev et al. (2015) methodology. In general, Gorgon performs well here too, performing similarly to other GMHD codes. Whilst performing well at the sub-solar magnetopause location, it performs less well in the lobe magnetic field and plasma sheet pressure. However, this is a well known feature, seen in other GMHD codes, and is due to the applicability of MHD in the inner magnetosphere. These results are expected to improve by implementing an ionosphere or ring current model into Gorgon.

The Gorgon code is used to study the dynamic behaviour of the outer boundaries of planetary magnetospheres, the bow shock and magnetopause. This variability of the outer boundaries due to magnetospheric drivers is not usually captured in empirical models. In simulations, however, the position of the outer boundaries are calculated self-consistently, allowing a better determination of the position and dynamics of these boundaries as they vary. Depending on the planet, there are varying drivers of the magnetosphere. In planets close to the Sun, the magnetosphere is mostly dominated by the solar wind interaction, largely due to high solar wind dynamic pressures and IMF strengths and the lack of internal drivers. In the gas giants, Jupiter and Saturn, they are strongly driven by internal drivers, namely active moons which eject plasma into the system. The Ice Giants, Uranus and Neptune, do not contain strong internal drivers, but have large dipole tilt angles, causing strong daily and seasonal variations.

In chapter 3, Neptune's magnetosphere is studied to understand how a precessing dipole affects the magnetospheric boundaries. The work is focussed on recreating the conditions as observed by Voyager 2 when it encountered the planet in 1989, and is the first published fully three-dimensional global MHD simulation. Neptune's unique combination of dipole and rotation axis tilt meant that its magnetosphere rotated through two main phases: Earth-like, when the magnetosphere resemble Earth's with the dipole axis approximately perpendicular to the solar wind flow; Pole-on, when the dipole axis is approximately parallel to the solar wind flow. As the planet precesses around, the outer boundaries responded by moving accordingly, being largest during the pole-on configuration, and smallest during the Earth-like configuration. The precession of the planetary field also modulated the reconnection region: for the last IMF configuration observed by Voyager 2, reconnection is favourable in the northern hemisphere, switching on as the planet rotated from Earth-like, and turning off shortly after Pole-on. By recreating Voyager 2's flyby, Gorgon was compared to the observations made by Voyager 2. Though there were some discrepancies in the plasma data within the magnetosphere, the

magnetic field both on the inbound and outbound legs of the flyby agrees remarkably well with the Voyager 2 data.

These new insights into the global restructuring of Neptune's magnetospheric field have implications on previous modelling and the Voyager 2 spacecraft observations. The Neptunian system is inherently three-dimensional and time-dependent, which could explain differences in the results this thesis and previous two-dimensional, steady state modelling, and affirms the need for three-dimensional simulations. For example, the location and modulation of magnetopause reconnection has implications on the global circulation of plasma and energy in the magnetosphere, and is likely to inhibit the Dungey cycle at Neptune in a similar way proposed by Cowley (2013). Furthermore, the position of reconnection in the simulation may explain the distinct lack of dynamics observed by Voyager 2. This work can be extended in a multitude of ways by varying the parameters of the system, such as the solar wind or time of year. This will allow for a more general view of the Neptunian system. Furthermore, by including Triton in the simulation, insight can be gained into its role in the magnetosphere. This information could be used to aid in planning new spacecraft missions to Neptune by providing new predictions.

In chapter 4, the focus returns to Earth by studying the position of the bow shock under a varying solar wind, as observed by the upwind spacecraft ACE during an event first analysed by Maksimovic et al. (2003). Due to rapid variations in the solar wind, the outer boundaries' position and velocity undergo rapid changes in response to the local solar wind parameters. The curved shape of the bow shock causes the same solar wind fronts to reach the shock at different times, causing the sub-solar region to respond to the new front before the flanks: this creates a shock shape that is never in a steady state. Analysis with virtual spacecraft reveals that spacecraft observations of the shock velocity may be skewed toward higher shock velocities, as the spacecraft rarely crosses a stationary shock. Furthermore, the shock is found to travel fastest in the flanks of the shock, highlighting a potential bias in spacecraft observations due to their orbit.

This work highlights the effects the variability of the solar wind has on the outer boundaries and their relation to empirical models. Many empirical models of the bow shock use a single set of upstream solar wind parameters to determine the size and shape of the bow shock. However, due to the timing between different solar wind fronts, each location of the shock is subjected to different upstream parameters, invalidating the assumption of a single set of

upstream parameters. This is important as these empirical models are often used to place spacecraft observations in a global context.

In chapter 5, the flux ropes generated by the event in chapter 4 are studied. Due to the strongly southward IMF, flux ropes are generated on the dayside magnetopause. Two types of flux ropes are primarily identified by their movement over the magnetopause. The first type have completely open topology and propagate over at high latitudes. The second type are more complex with interwoven open, closed and IMF field lines. These are approximately confined to the ecliptic plane, moving around the dusk flank. This is due to the $\vec{J} \times \vec{B}$ force of the kinked North- and South- connected open field lines acting in opposite directions, producing a net zero force in the $\hat{z}$ direction. These flux ropes dissipate by unravelling along the flank magnetotail. This unravelling only occurs when the flux rope disconnects from the Earth, consisting solely of IMF flux. However, the reconnection site is not necessarily at the flux rope, but closer to the Earth.

A general open question in magnetospheric physics is the extent to which FTEs impact the magnetosphere. Flux ropes with a proportion of closed magnetic flux, as observed in the simulation, may have a reduced of the amount of flux transferred into the magnetosphere. While MHD cannot produce as detailed results as kinetic or hybrid codes, they can infer the topology such flux ropes. This can be used in conjunction with spacecraft observations (e.g. MMS, THEMIS and Cluster) to gain further insight.

The work in this thesis has shown that Gorgon is a capable simulation, able to find novel insights into the dynamics of the magnetosphere. Many of the tools, such as magnetopause and flux rope finding algorithms, can be utilised generally in other studies. As with simulations in general, there are always ways in which simulations can be improved. The most prominent in Gorgon is the lack of an ionosphere. The boundary used is a resistive inner boundary, which closes the currents in the magnetosphere. However, it misses vital physics, such as the Hall effect and a spatially variable conductance (e.g. high conductivity on the dayside ionosphere due to solar EUV radiation). This should improve predictions in the magnetotail (Ridley, Gombosi, and Dezeewu, 2004), and be able to correlate processes in the magnetosphere to their signatures in the ionosphere.

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