

The Magnetic World of Ganymede: An Assessment of Cross-Magnetopause Energy Transport and its Effects on the Only Known Moon Magnetosphere

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Declarations

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- **Chapter 3:** Kaweeyanun, N., Masters., A., Jia, X. (2020). Favorable Conditions for Magnetic Reconnection at Ganymede's Upstream Magnetopause. *Geophys. Res. Lett.*, 47 (6), e2019GL086228. <https://doi.org/10.1029/2019GL086228>
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Abstract

Ganymede – Jupiter’s largest moon – generates a permanent magnetic field due to its Earth-like molten iron core. The internal magnetism affords Ganymede its own magnetosphere that is encased by Jupiter’s field-carrying plasma sheet at all times. However, this unique magnetic domain remains poorly understood due to insufficient observational data. This thesis provides fundamental assessments on aspects of Ganymede’s magnetosphere through analytical modelling, and offers baseline predictions for future observations of the moon.

The first work chapter considers the upstream magnetopause, where magnetic reconnection is expected to drive Dungey-like plasma convection in Ganymede’s magnetosphere. I evaluate magnetic reconnection onset to be favourable across the entire boundary, with average rates exhibiting Jovian half-synodic variation consistent with governance by Jupiter’s rotation. The second work chapter then determines that magnetopause reconnection is significantly localised as flux-transfer events, and sufficiently rapid that Ganymede’s magnetotail must also contain rapid reconnection site(s) to maintain a stable magnetosphere.

The third work chapter returns to the magnetopause and assesses Ganymede’s Kelvin-Helmholtz (K-H) instability growth. Linear K-H waves are expected to form along magnetopause flanks, with notable inter-flank asymmetry due to finite Larmor radius effects. However, nonlinear K-H vortex growth should be suppressed by concurring reconnection events, limiting the instability’s role in Ganymedeian plasma convection. Finally, the fourth work chapter estimates the magnetopause’s Chapman-Ferraro (C-F) magnetic field, which varies at Jovian half-synodic period corresponding with near-boundary conditions. The C-F signal is shown to induce a secondary magnetic response from Ganymede’s subsurface ocean of magnitude $\sim 1 - 10$ nT, which is resolvable by modern magnetometers and establishes the magnetopause as a potential probe for ocean study via magnetic induction.

The thesis concludes with introduction of the upcoming JUperiter ICy moon Explorer mission, which will not only verify predictions from the work chapters, but also exponentially advance our understanding of Ganymede as a whole.

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Chapter 1: Introduction

1.1. A Magnetic Moon

In 365 BC, a Chinese astronomer Gan De was viewing a summer night sky when he wrote down the following observation of planet Jupiter (Hughes, 1982).

“In the year of chan yan . . . , Jupiter was in Zi, it rose in the morning and went under in the evening together with the lunar mansions Xunli, Xli and Wei. It was very large and bright. Apparently, there was a small reddish star appended to its side. This is called 'an alliance'.”

Gan De, 365 BC

Gan De was a premier astronomer of his time, documenting his observations in the *Treatise on Jupiter* and the *Treatise on Astronomical Astrology* (Xida, 1368-1614 reprint). While both books would be lost in time, some of his notes, including this observation of an ‘alliance star’, would find their way into the *Treatise on Astrology of the Kaiyuan Era* (717-724 AD), and later uncovered in the modern times by a historian Xi Ze-Zong. Xi proposed that the ‘alliance star’ likely refers to one of the four larger moons of Jupiter, and showed that these moons can be visible to the naked eyes under rare good conditions despite their relatively small sizes (Xi, 1982). This claim would be collaborated by David W. Hughes, who believed that the moon in question is either Callisto or Ganymede, with the latter being more likely given its greater apparent brightness in the night sky (Hughes, 1982).

If Gan De had indeed spotted Ganymede, it would have been by far the earliest recorded observation of the largest Jovian moon, predating the more famous discovery from Galileo Galilei by almost two thousand years. The Italian astronomer was aided by a recent invention of the telescope, which extricated astronomy from the power of the naked eye. Galileo announced his remarkable findings on the title page of *Sidereus Nuncius*, first published in 1610 (Galilei, 1610)

“... with the help of a spyglass lately devised by [Galileo], about the face of the Moon, countless fixed stars, the Milky Way, nebulous stars, but especially about four planets flying

around the star of Jupiter at unequal intervals and periods of wonderful swiftness; which, unknown by anyone until this day, the first author detected recently and decided to name Medicean Stars.

Galileo Galilei, 1610

Galileo proceeded to carefully illustrate the four moons' daily positions relative to Jupiter over the course of two months between January and March of 1610 (Figure 1). He collectively called these satellites the Medicean Stars, after his patron Cosimo (II) de Medici, the Grand Duke of Tuscany. The Italian's claim of first discovery was somewhat disputed at the time by a German astronomer Simon Marius, who claimed to have started cataloguing the same Jovian moons independently a few months prior (Pasachoff, 2015). However, Marius did not circulate his findings locally until 1611, or internationally until 1614 when he published *Mundus Iovialis*, describing (without diagram) the motions of the moons (Marius, 1614). Galileo would accuse Marius of plagiarism, a charge that remained uncleared until a jury cross examination in the early 20th century (Oudemans & Bosscha, 1903; Bosscha, 1907). The investigation indicated that Marius observed the moons independently, but did not start his recordings until one day after Galileo's. Irrespective of the discovery priority, *Mundus Iovialis* would provide the modern names for the Galilean moons, whose ideas Marius credited to a fellow German astronomer Johannes Kepler.

"Jupiter is much blamed by the poets on account of his irregular loves. Three maidens are especially mentioned as having been clandestinely courted by Jupiter with success. Io, daughter of the River Inachus, Callisto of Lycaon, Europa of Agenor. Then there was Ganymede, the handsome son of King Tros, whom Jupiter, having taken the form of an eagle, transported to heaven on his back, as poets fabulously tell I think, therefore, that I shall not have done amiss if the First is called by me Io, the Second Europa, the Third, on account of its majesty of light, Ganymede, the Fourth Callisto"

Simon Marius, 1614

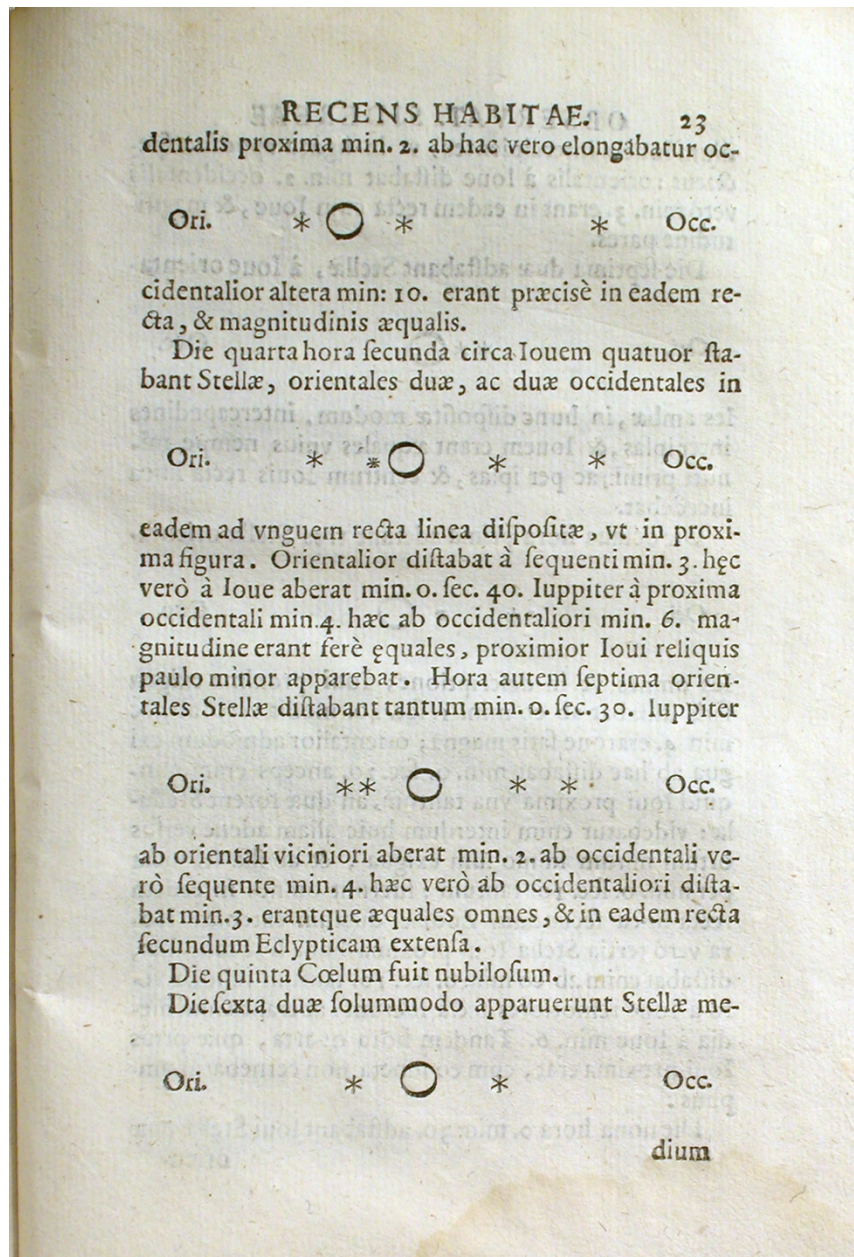


Figure 1: A page of Galileo's notes on the positions of the four largest moons of Jupiter. Image courtesy of History of Science Collections, University of Oklahoma Libraries (2001).

Over the next three centuries, our knowledge of Ganymede developed through Earth-based observations with increasingly powerful and numerous telescopes. Ganymede's positional evolution relative to other Galilean moons and Jupiter were carefully catalogued, and then compared to complex ephemeris tables built on advanced orbital mechanics (Sampson, 1921; de Sitter, 1931). This combination of mathematics and observation allowed some notable insights into Ganymede including estimates for the moon's mass and size (Brouwer & Clemence, 1961; Morrison et al., 1977). During this time, Ganymede was referred to as Jupiter

III, a designation adopted by Galileo. It was not until the mid-20th century, with rapid discovery of outer planet moons, that the name Ganymede entered common parlance.

The advent of the Space Race brought an explosion of investment into space travel and research. Development of unmanned space probes sparked new ambitious frontiers for space exploration. It was then that the Jet Propulsion Laboratory (JPL), under the National Aeronautics and Space Administration (NASA), forecasted a rare alignment of all outer Solar System planets (Jupiter, Saturn, Uranus, and Neptune) in the late 1970's (Newburn, 1970). During this window of opportunity, which occurs only once every 175 years, a single spacecraft would be able to complete a 'Grand Tour' of every outer planet with help from gravity assists¹, providing a great wealth of scientific outputs for comparatively efficient monetary input (e.g., Kohlase & Penzo, 1977).

The Grand Tour objective produced four spacecraft that would visit Jupiter and consequentially Ganymede. Pioneer 10 and Pioneer 11, launched for capability testing, were the first to arrive in 1973 and 1974 respectively (e.g., Siddiqi, 2018). They were followed by Voyager 1 and Voyager 2, the Grand Tour spacecraft, both of whom arrived in 1979 (Kohlase & Penzo, 1977). Each spacecraft encountered Ganymede once, flying past the moon at high velocity. Pioneer and Voyager provided the first in-situ observations of Ganymede that verified (and improved upon) previously remote estimates of the moon's mass and radius (Anderson et al., 1974; Null, 1976; Hollingsworth Smith, 1978; Smith, Soderblom, Johnson, et al., 1979). But more notably, the four spacecraft gave us the first high-resolution photographs of Ganymede that captured topological details of the moon's surface, which are shown in Figure 2. The images, alongside measurements of high reflectance (albedo), suggested that Ganymede possessed a surface of ice which could make up a significant portion of the moon's internal structure (Smith, Soderblom, Johnson, et al., 1979). Despite all the information obtained, neither Pioneer nor Voyager spacecraft ventured very close to Ganymede. The closest approach was achieved by Voyager 2 at 62,130 km (Mueller, 2011), which is over 23 times the radius of Ganymede (using modern estimate of $R_G = 2,634$ km).

¹ A mechanism in which a spacecraft 'slingshots' around a planet, using the planet's gravitational force to boost the spacecraft's velocity via transfer of momentum.

Consequently, much of Ganymede's nearby environment remained undiscovered until the arrival of Galileo. Built by JPL/NASA and named after the eminent Italian astronomer, the Galileo spacecraft was the first outer planet mission to specifically orbit around a planet – Jupiter (NASA, 1995). Between orbital insertion in 1995 and decommission in 2003, Galileo made 33 full orbits around Jupiter, each including a flyby of a Galilean moon (Meltzer, 2007). Ganymede received six flybys from Galileo (G1, G2, G7, G8, G28, and G29). These encounters were not only more in number than all previous missions combined, but they also occurred

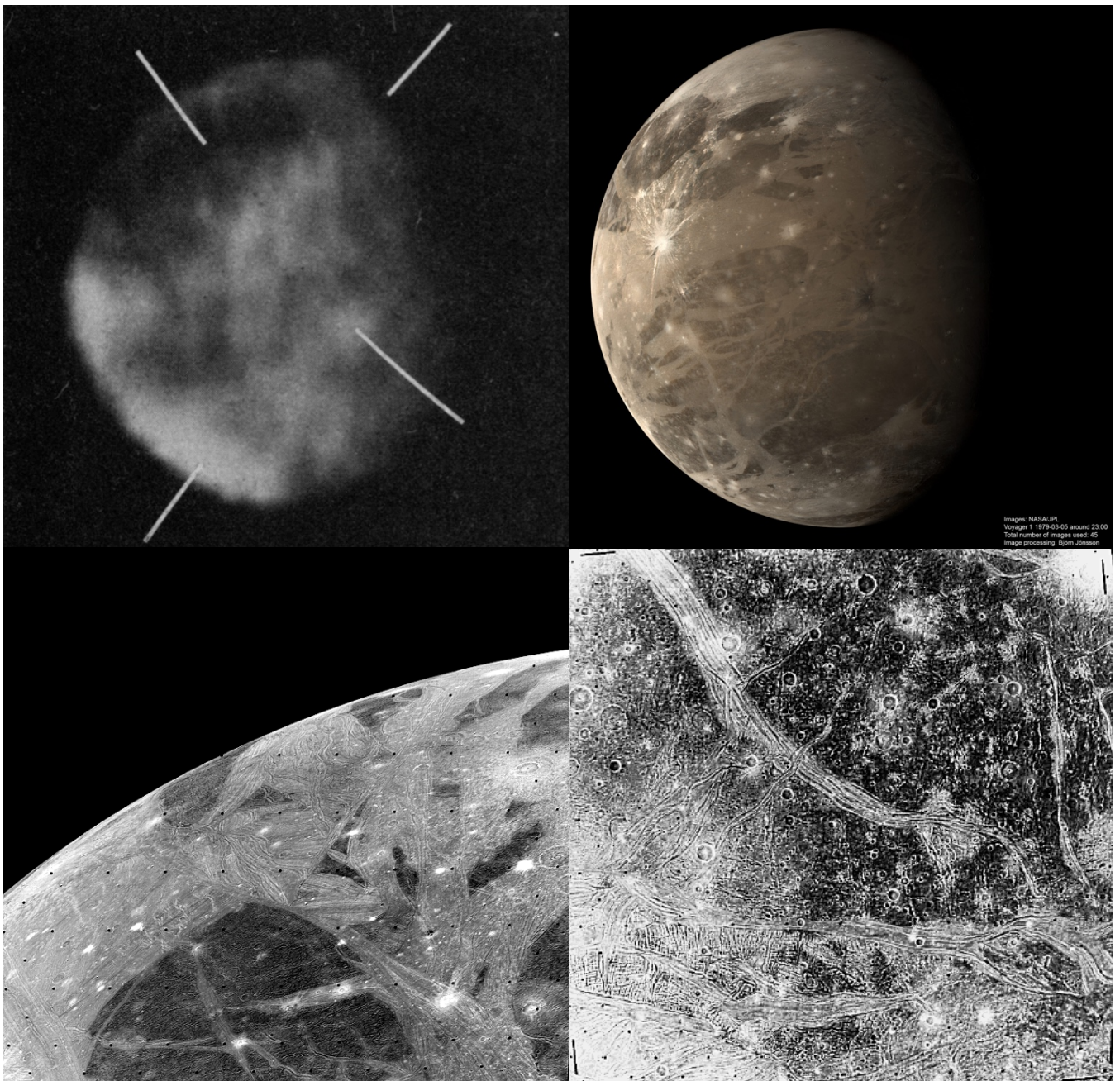


Figure 2: (Top Row) Images of Ganymede from (a) Pioneer 10 and (b) Voyager 1. (Bottom Row) Voyager 2 images of Ganymede's surface at (c) Nipper Sulcus and (d) Lagash Sulcus. Images courtesy of (a) Frieden & Swindell (1976), *Science*, (b) Jónsson (1979) (c-d) Williams & Friedlander (2015a, 2015b), NASA.

at much closer distances ranging between 260 km (G2) and 3,102 km (G7) at closest approaches, the upper limit of which was less than 2 R_G away from Ganymede's surface (Meltzer, 2007).

Similar to previous Pioneer and Voyager probes, Galileo carried aboard a number of scientific instruments which are shown in Figure 3. The bottom, non-spinning section contained an imaging camera with 20-1,000 times higher resolution than Voyager's (NASA, 1995), alongside spectrometers and a polarimeter for measuring radiation across the electromagnetic spectrum. Meanwhile, the top, spinning section contained subsystems designed to detect intensity, temperature, and other qualities of dust and energetic charged particles that come into contact with the spacecraft. Finally, on the antenna extending away from Galileo's main body was the magnetometer, which measured the strength and direction of magnetic field experienced by the spacecraft.

It was during the first flyby of Ganymede (G1) when Galileo's magnetometer detected a large spike in magnetic field strength, alongside a significant change in field direction, as the spacecraft ventured in the vicinity of the moon in Figure 4 (Kivelson et al., 1996). At the same time, the energetic particle detector observed electrons trapped under influence of a strong

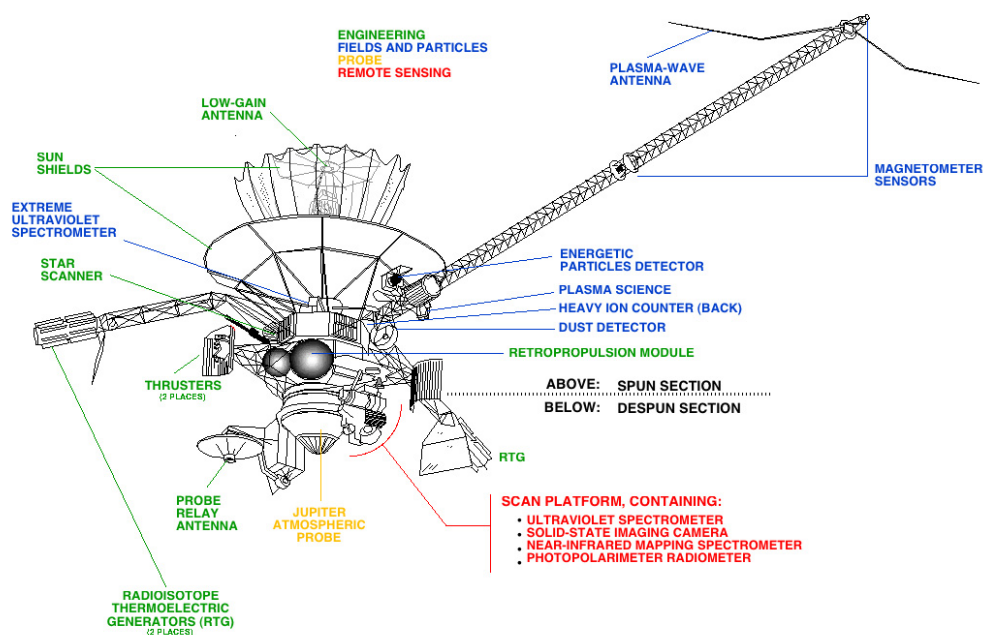


Figure 3: Schematic diagram of scientific instruments aboard Galileo. Image courtesy of NASA (2004).

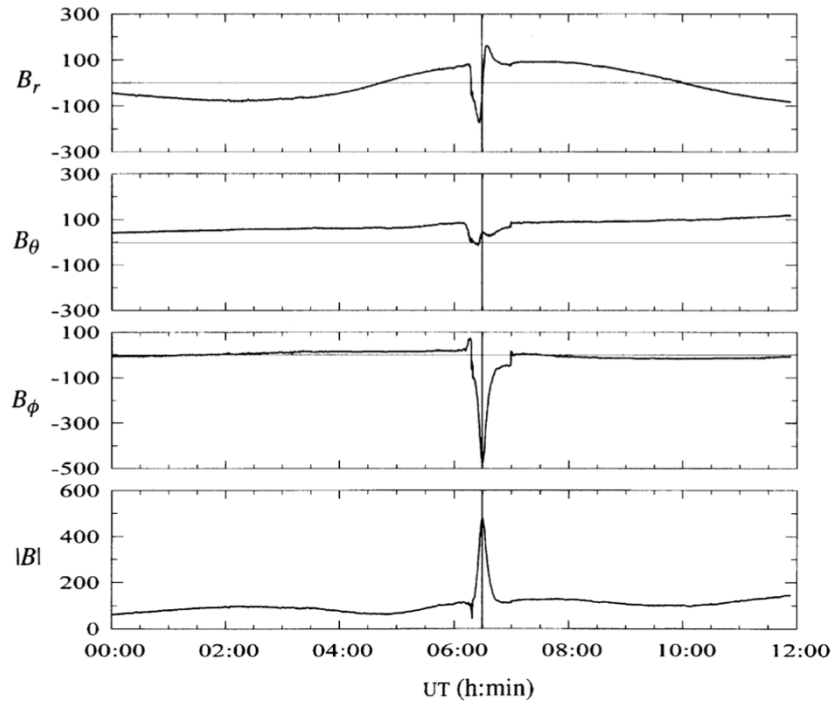


Figure 4: A notable spike in magnetic field strength as Galileo flew past Ganymede. Image courtesy of Kivelson et al., (1996), *Nature*.

localised magnetic field (Gurnett et al., 1996). The two unusual findings together lead to one conclusion, that Ganymede must generate its own magnetic field whose source is permanent and internal to the moon.

The discovery puts Ganymede into a category distinct from all other Solar System moons, given that internal magnetism is a trait previously thought exclusive to planets and stars. In the following decades, numerous investigations have been made into the structure of Ganymede's magnetic field, and how it shapes the plasma environment around the moon. This avenue of research produces further significant discoveries, including a potential subsurface ocean whose liquid water may be a harbour for life. However, before we can review latest research on Ganymede's magnetic environment, we must familiarise ourselves with relevant concepts on magnetic and plasma behaviours in space, whose introductions shall cover the rest of this chapter.

1.2. Internal Magnetism of Celestial Bodies

When a celestial body is said to be permanently magnetic, it generates a large-scale magnetic field that remains stable over astronomical timescales i.e., significant relative to the body's

own lifespan. The idea was first proposed in 1600 with the publication of *De Magnete* by William Gilbert (Gilbert, 1600). According to the author, Earth behaves like a bar magnet due to internal magnetic materials, particularly lodestones. The modern explanation, however, did not emerge until 1919, when Joseph Larmor proposed the concept of a dynamo (Larmor, 1919).

An internal dynamo theory describes magnetic field production and maintenance by a rotating, convecting electrically conductive fluid, which typically resides in the celestial body's outer core. For telluric bodies like Earth and Ganymede, the fluid is liquid iron mantle (e.g., Jeanloz, 1990; Schubert et al., 1996). For gaseous bodies like Jupiter and Saturn, the fluid is thought to be metallic hydrogen (e.g., Fortney & Nettelmann, 2010; Helled & Guillot, 2013; Militzer et al., 2016). Kinetic energy from the body's rotation and other internal energy sources (e.g., energy released by the inner cores) drives fluid convection, which then creates magnetic field via the principle of electromagnetic induction (e.g., Roberts & Soward, 1992)

$$\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{U} \times \mathbf{B}) + \eta \nabla^2 \mathbf{B} \quad (1.1)$$

where $\mathbf{U}$ = fluid velocity vector, $\mathbf{B}$ = magnetic field vector, and η = magnetic diffusivity². The first and second terms on the equation's right-hand side describes magnetic field diffusion and advection of magnetic field by the fluid's motion respectively. Fluid convection is essential for long-term field stability as it counters magnetic decay via ohmic dissipation (e.g., Robert & Soward, 1992). In the next decades, the dynamo theory would become the consensus explanation for large-scale magnetic field generation for Earth and beyond.

Many attempts have been made to directly simulate magnetic dynamo models starting from numerically solving the induction equation. Simpler models are known as kinematic dynamos, in which the convective fluid is prescribed a fixed flow velocity and remains unimpacted by the magnetic field it produces (e.g., Roberts, 1972). More complex models are nonlinear dynamos, where the convective fluid velocity is a dynamic variable and influenced by the

² Diffusivity can be derived by $\eta = 1/(\sigma\mu)$ where σ is conductivity and μ is permeability of (free) space.

generated magnetic field (e.g., Parker, 1955). As of this writing, there are many self-consistent nonlinear dynamos, between which discrepancies remain due to different assumptions used in the models (e.g., Kono & Roberts, 2002; Stanley & Glatzmaier, 2009).

When considering a dynamo-generated magnetic field at large scale, it is reasonable to approximate the magnetic field mathematically as the negative gradient of a scalar potential, which is expressed as a series of spherical harmonic functions (e.g., Connerney, 1993).

$$\mathbf{B} = -\nabla\Phi_{\text{pot}} \quad (1.2)$$

$$\Phi_{\text{pot}} = R_P \sum_{l=1}^{\infty} \sum_{m=0}^l \left\{ \left(\frac{R_P}{r} \right)^{l+1} (g_l^m \cos m\phi + h_l^m \sin m\phi) P_l^m(\cos \theta) \right\} \quad (1.3)$$

The above equations utilise a spherical coordinate system (r, ϕ, θ) where r is radial distance from the origin, ϕ is the azimuthal angle, and θ is the co-latitude angle. R_P is the radius of the magnetic body. g_l^m and h_l^m are Gauss coefficients, and $P_l^m(\cos \theta)$ are Schmidt quasi-normalised associated Legendre polynomials. Degrees and orders of spherical harmonics are described by subscript l and superscript m respectively.

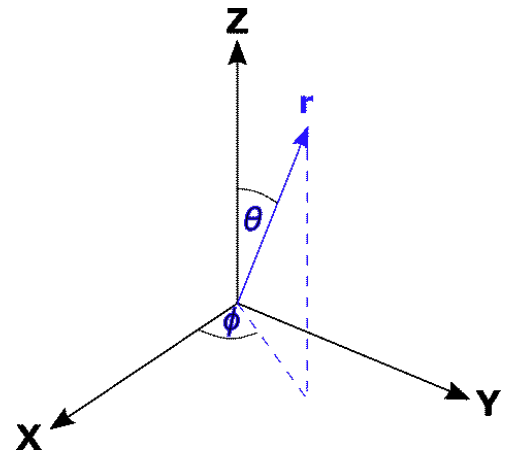


Figure 5: Overlay of Cartesian and spherical coordinate systems.

The degree parameter (l) describes the general structure of the magnetic field. Note that $l = 0$ produces a monopole field, which cannot exist according to electromagnetic theory³. Hence, the series starts from $l = 1$, which produces a dipole field equivalent to that of a bar magnet. In this case, the dipole axis orientation is determined by the three $l = 1$ Gauss coefficients, which translate to magnetic dipole axis vector (nT) in Cartesian coordinates where $g_1^0 = M_z$, $g_1^1 = M_y$, and $h_1^1 = -M_x$. The overlay of Cartesian and spherical coordinates is shown in Figure 5. Going up in degree, $l = 2$ produces a quadrupole field with five Gauss coefficients, $l = 3$ produces an octupole field, and so forth. Hence, any dynamo-

³ $\nabla \cdot \mathbf{B} = 0$ is the magnetic Gauss's law, which is part of the Maxwell's equations.

generated magnetic field can be broken down into summation of spherical harmonic fields starting from the simplest dipole. For Solar System bodies except Uranus and Neptune, the dipole field contribution is dominant, and higher degree contributions decrease with increasing l value (Connerney, 1993; Olsen et al., 2000; Kivelson et al., 2002).

So far, we have considered convection in a liquid outer core as the only source of large-scale magnetic field. However, certain celestial bodies contain more than one shell of electrically conductive fluid. Ganymede, in particular, is thought to have a large subsurface ocean beneath its surface (Kivelson et al., 2002). Under influence of the primary dynamo field, the saltwater will convect and induce a secondary magnetic field, which is then added to the final magnetic field. Details of the secondary induction process will be revisited in Chapter 6.

Figure 6 provides a simple illustration of Earth's magnetic field, which is here assumed to be dipole-like and in empty vacuum. Magnetic field lines flow away from (into) the northern (southern) magnetic pole, indicating a dipole axis that is almost perfectly antiparallel to Earth's rotation axis (with a 10° tilt). The field lines that have two ends at northern and southern magnetic poles are considered to be closed. Meanwhile, field lines that only have one end at the planet are considered to be open. From Eq. 1.2–1.3, the dipole field of degree $l = 1$ in spherical coordinates follows

$$B_r = \left(\frac{2R_p}{r}\right)^3 \{g_1^0 \cos \theta + [g_1^1 \cos \phi + h_1^1 \sin \phi] \sin \theta\} \quad (1.4)$$

$$B_\phi = \left(\frac{R_p}{r}\right)^3 \{g_1^1 \sin \phi - h_1^1 \cos \phi\} \quad (1.5)$$

$$B_\theta = \left(\frac{R_p}{r}\right)^3 \{g_1^0 \sin \theta - [g_1^1 \cos \phi + h_1^1 \sin \phi] \cos \theta\} \quad (1.6)$$

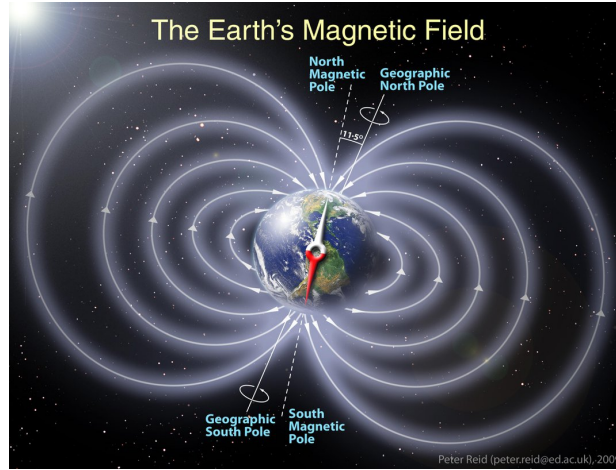


Figure 6: A simplified illustration of Earth's magnetic field lines. Image courtesy of Reid, (2009).

For this point, we can consider a simpler case where the dipole axis is aligned with the Cartesian Z-axis. In this case, the $m = 1$ order gauss coefficients are zero i.e., $g_1^1 = h_1^1 = 0$, therefore the dipole expression is reduced to

$$\mathbf{B} = g_1^0 \left(\frac{R_P}{r} \right)^3 (2 \cos \theta \hat{\mathbf{r}} + \sin \theta \hat{\boldsymbol{\theta}}) \quad (1.7)$$

At a particular radial distance and latitude, the dipole field strength then follows

$$B = |g_1^0| \left(\frac{R_P}{r} \right)^3 (1 + 3 \cos^2 \theta)^{1/2} \quad (1.8)$$

Eq. 1.8 illustrates that (a) the dipole field strength decreases cubically with radial distance, and (b) at particular distance r , a dipole is twice as strong at the pole ($\theta = 0^\circ$) as at the equator ($\theta = 90^\circ$) with a monotonic decrease in-between. Magnetic field strength then determines the pressure a magnetic field exerts on its surrounding by the formula.

$$P_B = \frac{B^2}{2\mu_0} \quad (1.9)$$

where $\mu_0 = 4\pi \times 10^{-7} \text{ NA}^{-2}$ is the permeability of free space.

The dipole representation of a celestial magnetic field is overly simplistic not only due to higher-degree contributions and potential secondary internal sources, but also because the field is not in complete vacuum and therefore must interact with its environment. Surrounding a celestial magnetic field are other large-scale magnetic fields and charged particles in the form of plasma, both of which will significantly shape the structure of a celestial body's magnetic domain. In the next section, we will introduce the fundamental properties of plasma, and discuss the intimate relationship that allows plasmas to exert vital influences on magnetic fields.

1.3. Plasma in Space

Plasma is commonly considered to be the fourth fundamental state of matter, first named and systematically studied by Irving Langmuir in the 1920's (Langmuir, 1928). A plasma is produced when a substance becomes significantly ionised, either by heat or potential difference, to the point where motions of its positive ions and negative electrons are strongly governed by long-ranged electromagnetic force between the particles. Plasma can be found throughout the universe, from distant nebulae to small colourful lamps (Figure 7).

On Earth, the human application of plasma is diverse and numerous, including electronics, industrial processing, measurement instruments, and potentially energy generation through fusion. In space, natural plasmas drive wind-like transport of particles around large planetary systems (e.g., the Solar System) and individual celestial bodies (e.g., Earth and Jupiter) (e.g., Parker 1958a; Dungey, 1961, Vasyliūnas, 1983). Hence, plasma is responsible for many notable phenomena in space, including solar storms, radiation belts, and auroras (e.g., Parker, 1958b; Arnoldy et al., 1962), some of which are exemplified in Figure 7. The study of plasma influence on the immediate environment around a body is known as space weather, which is rapidly gaining traction given our ever-increasing satellite capabilities around Earth and beyond (e.g., Pulkkinen, 2007).

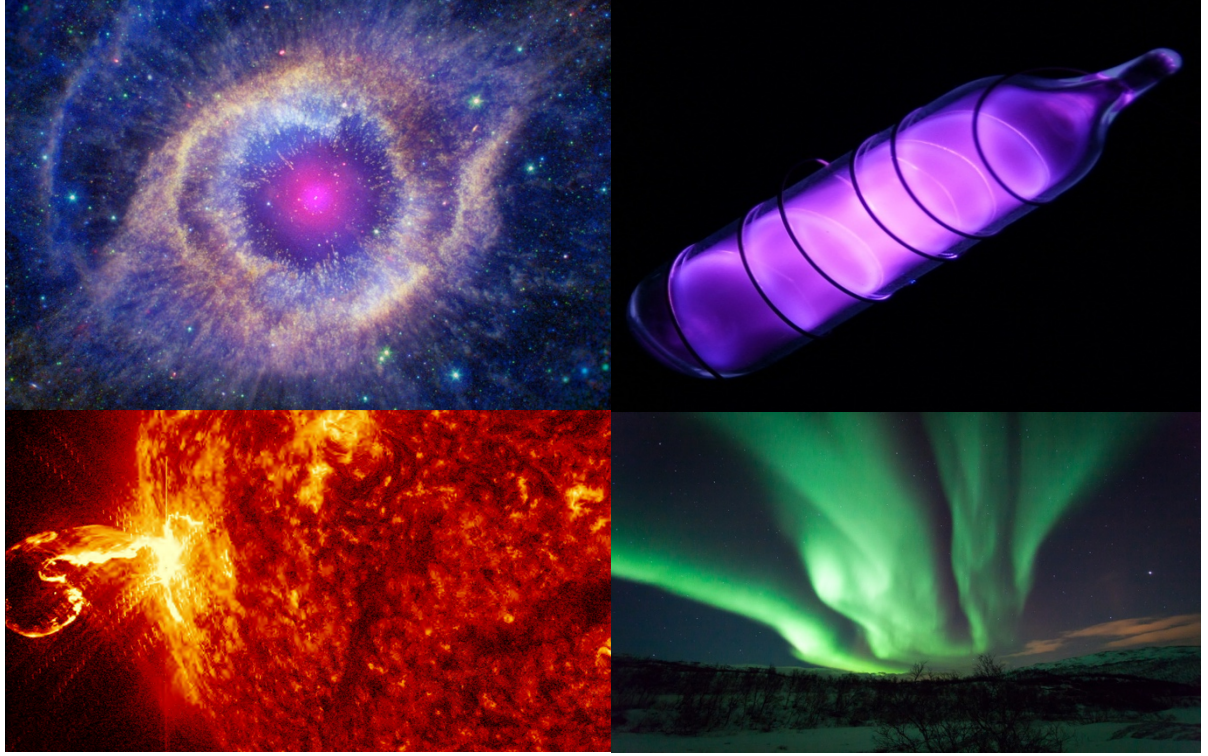


Figure 7: (Top-left) The Helix nebula. (Top-right) An argon plasma lamp. (Bottom-left) Coronal mass ejection from the Sun. (Bottom-right) Aurora display seen from Earth. Images courtesy of (top-left) NASA, STScI, ESA, NRAO (2020), (top-right) *Kummer, (2016)*, (bottom-left) *Gleber, NASA, (2014)*, (bottom-right) *Tunli, 2015*.

A plasma can be considered as a mixture of positive ion and negative electron populations, each with its own number density (n) and temperature (T). At large scale, a plasma is quasi-neutral. This means that its overall net charge is zero, and that ion and electron densities (n_i and n_e) are related by

$$n_i = \langle Q \rangle n_e \quad (1.10)$$

where $\langle Q \rangle$ is the average charge of ions in the plasma. The extent of ionisation in plasma is thus described by the fraction

$$\alpha_i = \frac{n_i}{N} \quad (1.11)$$

where N is the total density of ions and unionised neutrals. A plasma is fully ionised if $\alpha_i = 1$ and weakly ionised if $\alpha_i \ll 1$. The ionisation threshold for plasma creation can be as low as $\alpha_i \sim 10^{-6}$ (Conde, 2018). Particle ionisation in plasma typically occurs via impact collision,

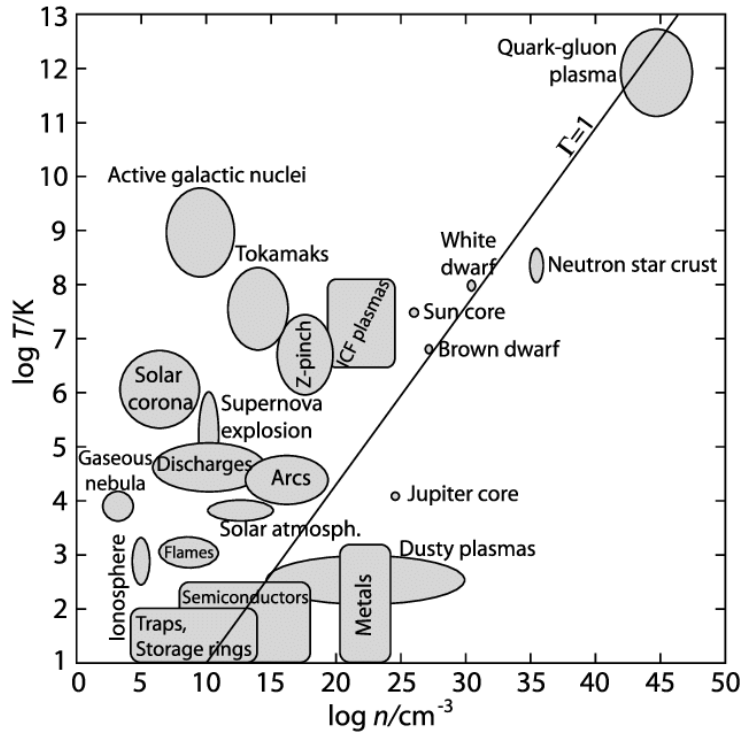


Figure 8: Plasma classification by number density and temperature. Solid line indicates $\Gamma = 1$ gradient where Γ is ratio between potential and kinetic energy. Image courtesy of Donkó et al., (2007).

hence the particles must possess considerable kinetic energy to maintain electric conductivity (Conde, 2018). As the result, temperatures of ions and electrons in plasmas typically exceed 10^{5-6} K, and they are often expressed in the electron-volts where $1 \text{ eV} \approx 11,600 \text{ K}$.

Figure 8 shows classification of various plasmas by their number density and temperature. Densest plasmas are found in cores of stars and planets ($n_i > 10^{20} \text{ m}^{-3}$), with very high temperatures ($T_i > 10^7 \text{ K}$) required to maintain ionisation in the plasma. Medium density plasma ($n_i > 10^{15} \text{ m}^{-3}$) can be found in many environments, from magnetic fusion confinements to semiconductors, with more relaxed constraints on viable temperature. Space plasmas, meanwhile, possess lowest relative densities ($n_i < 10^{10} \text{ m}^{-3}$) compared to other plasmas. Their temperatures are dependent on those of their sources, as seen from higher temperatures in nebulae ($T_i \sim 10^6 \text{ K}$) compared to stellar winds and planetary plasmas ($T_i < 10^5 \text{ K}$).

While a plasma is quasi-neutral and behaves collectively at large scales, effects of individual ions and electrons become more important at local levels. The sphere in which the electric

field produced by individual charges cannot be neglected is known as the Debye sphere, whose radius – the Debye length – is defined by (Debye & Hückel, 1923)

$$\lambda_D = \sqrt{\frac{\epsilon_0 k_B / q_e^2}{n_e / T_e + (\sum_j Z_j^2 n_j) / T_i}} \approx \sqrt{\frac{\epsilon_0 k_B T_e}{n_e q_e^2}} \quad (1.12)$$

where $\epsilon_0 = 8.85 \times 10^{-12}$ F/m is the permittivity of free space, $k_B = 1.38 \times 10^{-23}$ JK⁻¹ is the Boltzmann constant, $q_e = 1.6 \times 10^{-19}$ C is the electron charge, n_e and T_e are electron density and temperature, Z_j and n_j are charge (in unit of q_e) and density of individual ion species, and T_i is ion temperature in the plasma. In most cases, ions are assumed to be stationary in the plasma given their large masses relative to electrons, so the ion term is dropped from Eq. 1.12 leading to a much simpler expression. If a plasma's Debye length is much smaller than its physical size, then the electromagnetic interactions within can be considered as if they occur inside a bulk volume, and boundary effects can be ignored.

The number of electrons that can fit into a single Debye sphere is called the Debye number (Gibbon, 2014)

$$N_D = \frac{4}{3} \pi n \lambda_D^3 = \frac{\Lambda}{3} \quad (1.13)$$

where Λ is defined as the plasma parameter. If $N_D \gg 1$, the Debye sphere is densely populated and any change in local charge (ions) will be quickly crowded by redistributed opposite charge particles (electrons), screening its effects from the rest of the plasma. Consequently, charged particles will interact collectively with external electromagnetic fields rather than in pairwise electromagnetic (Coulomb) scattering with other individual particles. If $N_D < 1$, local particle charges are ineffectively screened, and Coulomb scatterings become significant.

When a plasma experiences an external electric field, its ions and electrons become displaced from their original positions. This creates two plate-like positive and negative layers, between which exists an internal electric field that exerts restoring forces on the charged particles

(Tonks & Langmuir, 1929). Electrons and ions thus oscillate in undamped simple harmonic motion with frequency

$$\omega_e = \sqrt{\frac{n_e e^2}{m_e \epsilon_0}} \quad (1.14)$$

$$\omega_i = \sqrt{\frac{n_i Z^2 e^2}{m_i \epsilon_0}} \quad (1.15)$$

As plasma ions are largely stationary due to their large relative mass, the electron oscillation frequency is usually taken as the plasma frequency $\omega_p = \omega_e$. From Eq. 1.14—1.15, we can define the inertial lengths of electrons and ions

$$d_e = \frac{c}{\omega_p} \quad (1.16)$$

$$d_i = \frac{c}{\omega_i} \quad (1.17)$$

where c = speed of light. The electron inertial length is also known as the skin depth $\delta = d_e$, which quantifies the extent of background field penetration into the plasma. Meanwhile, the ion inertial length defines a scale below which ion behaviours cannot be dictated by the background fields.

As electrons travel through a plasma, it is said to undergo a collision if a series (one or more) of Coulomb scatterings alter its direction by at least 90° . Without going into further detail, the electron-ion collision frequency can be approximately expressed as a function of Debye number, plasma parameter, and plasma frequency (Gibbon, 2014)

$$\frac{\nu_{ei}}{\omega_p} \simeq \frac{Z \ln \Lambda}{10 N_D} \quad (1.18)$$

If the electron-ion collision frequency far exceeds the electron-neutral collision frequency, then pairwise collisions are governed by electrostatic interactions over kinetic gas laws, and

the plasma is said to be collisionless. A plasma is also collisionless if the mean free path i.e., the average particle travel distance between two Coulomb scatterings, is much greater than the length scale of the system considered. The mean free path is defined as (Livadiotis, 2019)

$$\lambda_{\text{mfp}} \sim \Lambda \lambda_D \quad (1.19)$$

An effectively screened, collisionless plasma with negligible boundary effects is said to be *ideal*. In this condition, the ideal gas equation of state becomes applicable for each particle species in the plasma (Cravens, 1997). Ideal plasmas are generally hot and diffuse, while its non-ideal counterparts are cold and dense. One example of ideal plasma is the solar wind, which refers to plasma streams continually emitted by the Sun. From Figure 8, the solar wind has typical density $n_i \approx n_e \sim 10^7 \text{ m}^{-3}$ and temperature $T_i \approx T_e \sim 10^5 \text{ K}$ at 1 AU. Assuming fixed ions, the plasma has Debye length $\lambda_D \sim 10 \text{ m}$ and thus plasma parameter $\Lambda \sim 10^{10} \gg 1$. The solar wind travels millions of kilometres across the Solar System i.e., $L \sim 10^6 \text{ m}$, which is much greater than the Debye length but much less than the mean free path ($\lambda_{\text{mfp}} \sim 10^{11} \text{ m}$). Hence, the solar wind can be considered a bulk collisionless plasma with weakly-coupling particles – suitable for the ideal approximation. The same principle holds for other plasmas in space, including planetary plasmas which will be later revisited.

Plasma particles, be it slow-moving ions or fast-moving electrons, are in constant motion with varying velocities. As it is not practical to catalogue particles individually, their velocities are almost always described statistically through velocity distributions. For an ideal plasma in equilibrium, particle velocity distributions are most often Maxwellian i.e., they follow the Maxwell-Boltzmann distribution (e.g., Maxwell, 1860a)

$$f(v) d^3v = \left(\frac{m}{2\pi k_B T} \right)^{\frac{3}{2}} \exp \left(-\frac{mv^2}{k_B T} \right) d^3v \quad (1.20)$$

where $f(v)$ prescribes relative number of particles with velocity v , mass m , and collective temperature T in velocity space with volume element d^3v . The distribution can be broken

down into one-dimensional components (e.g., $v^2 = v_x^2 + v_y^2 + v_z^2$). It is then possible to derive the mode, mean, and root-mean-squared (rms) velocities of the particles

$$v_{i,\text{mode}} = \sqrt{\frac{2k_B T}{m}} \quad (1.21)$$

$$v_{i,\text{mean}} = \sqrt{\frac{8k_B T}{\pi m}} \quad (1.22)$$

$$v_{i,\text{rms}} = \sqrt{\frac{3k_B T}{m}} \quad (1.23)$$

where subscript i refers to a velocity component. The distribution is asymptotic toward zero as $v \rightarrow 0$ and $v \rightarrow \infty$, however this does not mean that extreme particle velocities are impossible, only unlikely. A particle population whose members follow the Maxwellian velocity distribution is said to be thermal. Thus, ideal plasmas contain thermal ion and electron populations. But unlike ideal gases where the vast majority of particles are always Maxwellian, ideal plasmas can contain significant non-Maxwellian, out-of-equilibrium particle populations, usually in the form of a superthermal tail of energetic particles that possess much higher temperature but comparable pressure compared to their thermal counterparts (e.g., Van Dam, 2006). In plasma, the superthermal tail has been described mathematically by the generalised Lorentzian (κ) distribution (Summers & Thorne, 1991).

For a thermal ideal plasma, average kinetic energy of a particle is linearly proportional to its temperature

$$\langle E_{i,e} \rangle = \frac{1}{2} k_B T_{i,e} \quad (1.24)$$

Then, through the ideal gas equation of state, the thermal pressure exerted by a plasma (assuming equal ion and electron temperatures) is

$$P_{\text{th}} = nk_B T \quad (1.25)$$

The sum of pressures from thermal and non-thermal populations are known as plasma pressure (P_p). Consequently, the ratio between plasma and magnetic pressure (Eq. 1.9) can be defined as the plasma beta

$$\beta = \frac{P_{th}}{P_B} \quad (1.26)$$

Although a plasma is defined by collective behaviours of its particles, at very low density individual charge behaviours can have significant large-scale impacts. Hence, in the case of ideal space plasma, it is useful to understand how one single charged particle reacts to external electromagnetic fields. Consider a simple configuration with one particle and a uniform magnetic field (and no electric field). This particle will experience a Lorentz force

$$\mathbf{F} = q(\mathbf{v} \times \mathbf{B}) \quad (1.27)$$

where q is the particle charge, $\mathbf{v}$ is the particle velocity vector, and $\mathbf{B}$ is the magnetic field vector. Since force is perpendicular to displacement at all times, the particle does no work as it circularly gyrates around the magnetic field line. The radius of gyration, also known as gyroradius or Larmor radius, is therefore

$$R = \frac{mv_{\perp}}{|q|B} \quad (1.28)$$

where m is the particle mass, and $v_{\perp}$ is the particle velocity component perpendicular to the magnetic field of strength B . The frequency of gyration, or gyrofrequency, is then

$$\Omega = \frac{|q|B}{m} \quad (1.29)$$

The centre point of gyration is known as the guiding centre, which remains stationary in case of uniform magnetic field (Figure 9a). In case of non-uniform field, however, the guiding centre can experience drifts due to gradient and curvature of the magnetic field. Eq. 1.28

shows that a particle's gyroradius decreases (increases) with stronger (weaker) magnetic field. Hence under a magnetic field gradient, the gyroradius will decrease and increase alternately as the particle experience both positive and negative changes in magnetic field, tracing out a cycloid path during which the guiding centre drifts in direction perpendicular to both the magnetic field and its gradient (Figure 9d)

$$\mathbf{v}_{\nabla B} = \frac{mv_{\perp}^2}{2qB} \frac{\mathbf{B} \times \nabla B}{B^2} \quad (1.30)$$

If the background magnetic field line is curved, the gyrating particle will experience a centrifugal force in addition to the Lorentz force. This results in a guiding centre drift out of the plane of magnetic field curvature (Figure 9e)

$$\mathbf{v}_R = \frac{mv_{\parallel}^2}{qB} \frac{\mathbf{R}_c \times \mathbf{B}}{R_c^2 B} = \frac{mv_{\parallel}^2}{qB} \frac{\mathbf{B} \times \nabla B}{B^2} \quad (1.31)$$

where $v_{\parallel}$ is the particle velocity component parallel to the magnetic field and $\mathbf{R}_c$ is the field's curvature vector. The relationship between radius of curvature and magnetic field is achieved through the electrostatic Ampère's Law⁴.

In reality, the external magnetic field is also accompanied by a corresponding electric field. A uniform electric field adds to the Lorentz force on the particle as follows

$$\mathbf{F} = q(\mathbf{v} \times \mathbf{B} + \mathbf{E}) \quad (1.32)$$

Like non-uniform magnetic field, this introduction adds a drift to the guiding centre known as the $\mathbf{E} \times \mathbf{B}$ drift

$$\mathbf{v}_E = \frac{\mathbf{E} \times \mathbf{B}}{B^2} \quad (1.33)$$

⁴ $\nabla \times \mathbf{B} = \mu_0 \mathbf{J}$

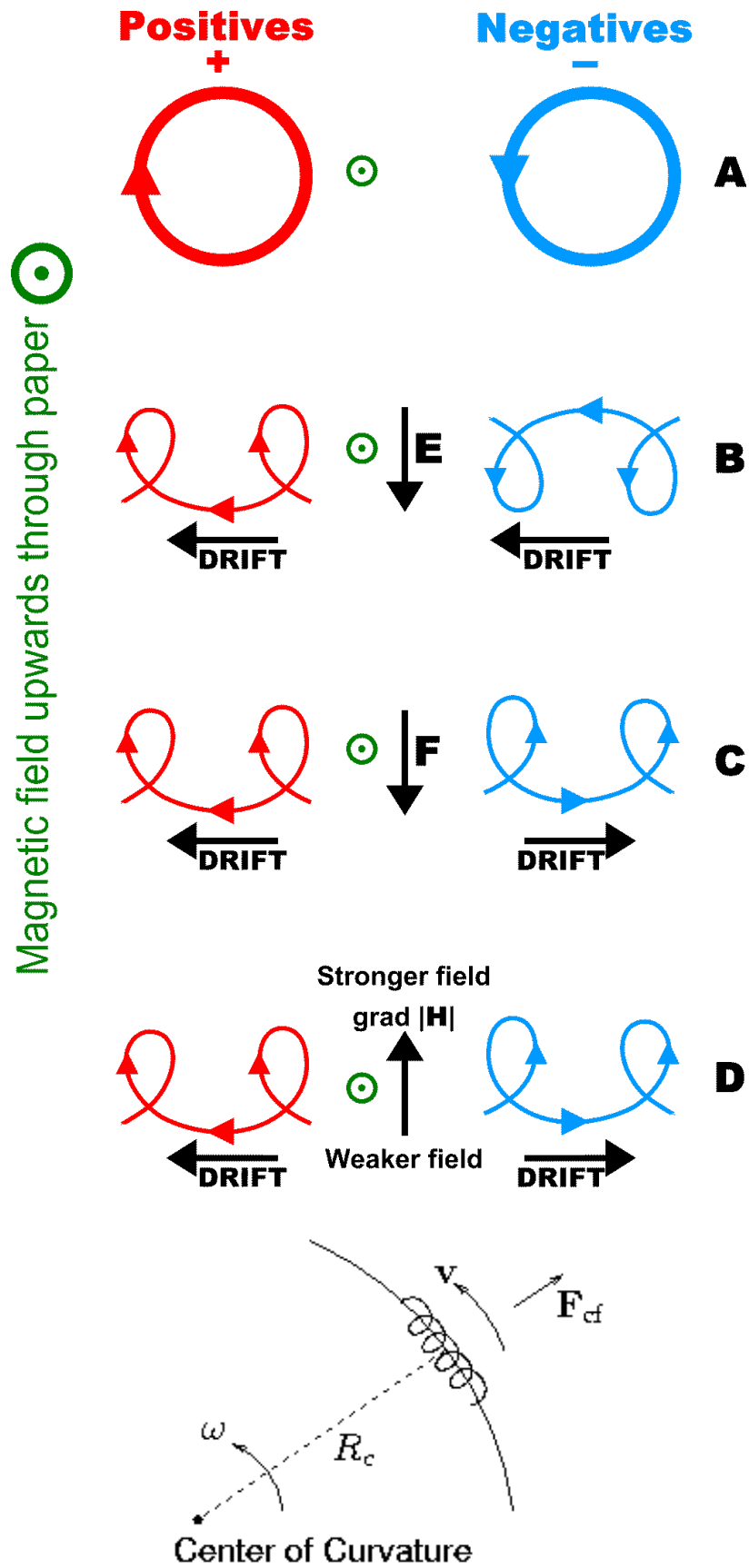


Figure 9: Illustrations of guiding centre drifts under influence of external force fields. (a) Larmor gyration under uniform magnetic field, (b) $\mathbf{E} \times \mathbf{B}$ drift, (c) gravitational (independent force) drift, (d) gradient drift under non-uniform magnetic field $\mathbf{B}/\mathbf{H}$, (Bottom) curvature drift under bent magnetic field. Images courtesy of (a-d) *Wikipedia*, (2007), (bottom) Hutchinson, (2001).

Since the electric addition to the Lorentz force is also charge-dependent, the $\mathbf{E} \times \mathbf{B}$ drift affects ions and electrons in the same way (Figure 9b), which is reflected by the lack of q parameter in Eq. 1.33. If the electric field is non-uniform, Eq. 1.33 is modified to account for electric field gradient.

Beyond electromagnetic force, particle motion can also be affected by the gravitational force field on the plasma, resulting in the following guiding centre drift

$$\mathbf{v}_g = \frac{m \mathbf{g} \times \mathbf{B}}{q B^2} \quad (1.34)$$

where $\mathbf{g}$ is gravitational acceleration vector⁵. Since gravitational field is not charge dependent, the subsequent guiding centre drift affects ions and electrons oppositely (Figure 9c). The dependence on mass also means the drift is largely negligible for low-mass electrons.

Highly diffuse space plasma sometimes can be treated as a group of individual particles whose motions translate to macroscopic fluid behaviour. A common example is magnetic mirrors, where gradient drifts and magnetic moment conservation together trap plasmas inside a magnetic field (Cravens, 1997). Otherwise, single particle motion can create anisotropy in plasma conditions and introduce complications to plasma-magnetic interactions (see Chapter 5).

Nevertheless, plasma analysis still requires theoretical frameworks that characterise complex collective behaviours. One example for this need is the diamagnetic drift. Under a pressure gradient not parallel to the background magnetic field, ions and electrons are affected differently following (Swisdak et al., 2003)

$$\mathbf{v}_d = \frac{\nabla P_{th} \times \mathbf{B}}{qnB^2} \quad (1.35)$$

⁵ Eq. 1.34 also applies for any arbitrary non-electromagnetic force field.

where P_{th} is thermal pressure for each particle species (Eq. 1.25). This divergence produces an electric field between two charged particle populations, whose magnetic effects weakens magnetic field inside the plasma (hence the diamagnetic name). However, there is no translation in guiding centre of any individual particle involved. The diamagnetic drift appears only when ions and electrons are viewed as collective populations inside a resting reference frame.

Taking a further step back, we can consider the entire plasma (ions and electrons) as a single fluid and use fluid dynamics knowledge to predict plasma behaviour. This framework is known as magnetohydrodynamics (MHD), which was first named by Hannes Alfvén (Alfvén, 1946), and underwent rapid development in the 1960's (Davidson, 2001).

The MHD theory builds on a number of important assumptions that must be kept in mind. First, the plasma is quasi-neutral with isotropic pressure and equal temperature between electron and ion populations. Second, the length scale of plasma considered is sufficiently large that plasma density and velocity do not vary within a fluid volume element i.e., length scale is larger than particle gyroradii, inertial lengths, and mean free paths. Third, the time scale characteristic of bulk plasma dynamics is much greater than that of particle phenomena within the plasma i.e., time scale is longer than particle gyrofrequency, oscillation frequency, and collision rates. Fourth, plasma is in equilibrium with Maxwellian particle velocity distribution (although it can be non-ideal). Finally, the MHD theory can only consider electrostatic situations i.e., negligible displacement currents.

Under MHD theory, a plasma fluid with velocity $\mathbf{U}$, mass density ρ , plasma pressure $P = P_i + P_e$ is governed by the continuity equation

$$\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{U}) = 0 \quad (1.36)$$

and the momentum equation

$$\rho \frac{\partial \mathbf{U}}{\partial t} + (\mathbf{U} \cdot \nabla) \rho \mathbf{U} = -\nabla P + \mathbf{J} \times \mathbf{B} + \mathbf{F} \quad (1.37)$$

where $\mathbf{J}$ is the current density and $\mathbf{B}$ is the background magnetic field (e.g., Cravens, 1997). Eq. 1.37 illustrates that change in plasma momentum is driven by negative pressure gradient, Lorentz force, and external force which is zero for ideal plasma. The MHD momentum equation is similar to the Cauchy fluid momentum equation, except the intermolecular force is electromagnetic. The Lorentz force term in particular can be expanded via the electrostatic version of Ampère's law to obtain

$$\mathbf{J} \times \mathbf{B} = -\nabla \left(\frac{\mathbf{B}^2}{2\mu_0} \right) + \frac{(\mathbf{B} \cdot \nabla) \mathbf{B}}{\mu_0} \quad (1.38)$$

The right-hand side terms show Lorentz force contributions from the magnetic pressure (Eq. 1.9) and magnetic tension respectively. The pressure force arises from inherent energy density of the magnetic field, while the tension force arises from curvature in the magnetic field lines and acts to straighten out the field.

Another important MHD equation is the generalised Ohm's law (e.g., Cravens, 1997)

$$\mathbf{E} + \mathbf{U} \times \mathbf{B} = \eta \mathbf{J} + \dots \quad (1.39)$$

where η is resistivity of the plasma. In case of small resistivity, higher order terms in Eq. 1.39's right-hand side can be neglected. If the plasma is ideal, then the fluid is perfectly conducting and $\eta = 0$. Consequently, the non-ideal MHD formulation is often referred to as resistive MHD, while the ideal formulation is referred to as ideal MHD.

Starting from Faraday's Law⁶, and then substituting in Eq. 1.39, it is possible to derive the MHD induction equation

⁶ $\nabla \times \mathbf{E} = -\frac{\partial \mathbf{B}}{\partial t}$

$$\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{U} \times \mathbf{B}) + \nabla \cdot \left(\frac{\eta}{\mu_0} \nabla \mathbf{B} \right) - \frac{1}{\mu_0} \nabla \mathbf{B} \cdot \nabla \eta \quad (1.40)$$

On the right-hand side, the third term is typically neglected in case of small resistivity. The second term describes plasma's ability to move across magnetic field lines due to its resistivity. If the plasma is ideal, however, then the second right-hand term is zero and the infinitely conducting plasma is 'frozen into' the magnetic field, inseparable from each other. This is known as Alfvén's theorem and thus the first right-hand term is often called the frozen-in term (Cravens, 1997). Since plasma in space is strongly ideal, the frozen-in condition is a vital tool for evaluation of plasma behaviour.

Finally, MHD plasma also obeys the adiabatic energy equation (e.g., Cravens, 1997)

$$\left(\frac{\rho^\gamma}{\gamma - 1} \right) \left(\frac{\partial}{\partial t} + \mathbf{U} \cdot \nabla \right) \left(\frac{P}{\rho^\gamma} \right) = -L \quad (1.41)$$

where $\gamma = \frac{5}{3}$ is the ratio of specific heat capacities and L is net energy loss-gain function. If the plasma is ideal, then Eq. 1.41 is simplified by setting $\left(\frac{\rho^\gamma}{\gamma - 1} \right) = 1$ and $L = 0$. The four equations of continuity (Eq. 1.36), momentum (Eq. 1.37), induction (Eq. 1.40), and energy (Eq. 1.41) are together considered the foundation of the MHD plasma theory.

If we consider an ideal plasma fluid with small linear departures from equilibrium

$$\mathbf{B} = \mathbf{B}_0 + \mathbf{B}_1(\mathbf{r}, t) \quad (1.42)$$

$$\mathbf{U} = \mathbf{U}_0 + \mathbf{U}_1(\mathbf{r}, t) \quad (1.43)$$

$$P = P_0 + P_1(\mathbf{r}, t) \quad (1.44)$$

$$\rho = \rho_0 + \rho_1(\mathbf{r}, t) \quad (1.45)$$

where $\mathbf{B}_0$, $\mathbf{U}_0$, P_0 , and ρ_0 are constant equilibrium conditions and $\mathbf{B}_1$, $\mathbf{U}_1$, P_1 , and ρ_1 are spatially dependent perturbations. If the perturbations are part of a plane wave, the linear terms are proportional to $\exp(i\mathbf{k} \cdot \mathbf{r} - i\omega t)$ and insertion of Eq. 1.42–1.45 into the (ideal) MHD equations yield two sets of dispersion relations (e.g., Cravens, 1997)

$$\omega^2 - \frac{B_0^2}{\mu_0 \rho_0} k^2 \cos^2 \theta = 0 \quad (1.46)$$

$$\omega^4 - \omega^2 k^2 \left(\frac{B_0^2}{\mu_0 \rho_0} + \frac{\gamma P_0}{\rho_0} \right) + k^4 \left(\frac{B_0^2}{\mu_0 \rho_0} \right) \left(\frac{\gamma P_0}{\rho_0} \right) \cos^2 \theta = 0 \quad (1.47)$$

where θ is the acute angle between equilibrium magnetic field $\mathbf{B}_0$ and the wave's propagation direction $\mathbf{k}$. Each dispersion relation represents a unique type of wave in MHD plasma. The first dispersion relation describes the Alfvén wave, a transverse incompressive oscillation in response to restoring force from magnetic tension (Eq. 1.38) which travels at phase speed

$$C_A = \sqrt{\frac{B_0^2}{\mu_0 \rho_0}} \quad (1.48)$$

The second dispersion relation describes the magnetosonic wave, which is a longitudinal compression of the plasma. Eq. 1.47 also has two ω^2 roots, this mean that the dispersion relation can be further broken down into solution corresponding to slow and fast magnetosonic waves. The slow wave propagates along the magnetic field at the sound speed in plasma

$$C_S = \sqrt{\frac{\gamma P_0}{\rho_0}} \quad (1.49)$$

Meanwhile, the fast wave propagates orthogonally to the field at a faster phase velocity

$$C_F = \sqrt{C_A^2 + C_S^2} \quad (1.50)$$

Note that if the plane wave propagation is parallel to the magnetic field ($\mathbf{k} \parallel \mathbf{B}_0$), a fast magnetosonic wave will degenerate into an Alfvén wave.

A bulk plasma flow can be described as super-Alfvénic (sub-Alfvénic) if its speed is faster (slower) than the Alfvén phase speed, and supersonic (sub-sonic) if its speed is faster (slower) than the sound speed. C_A and C_s are often used as significant benchmarks for dividing plasma into behavioural regimes. For instance, super-Alfvénic plasma flow is heavily influenced by its thermal pressure, while sub-Alfvénic plasma is more impacted by the surrounding magnetic pressure (e.g., Neubauer, 1998).

On the subject of pressure, a flow of plasma also generates a dynamic ram pressure which is equivalent to energy density of the flow

$$P_r = \frac{1}{2} \rho_0 U^2 \quad (1.51)$$

where U is in the rest frame of the pressured obstacle.

Given the ideal nature of plasmas in space, their bulk flows should be governed by Alfvén's frozen-in theorem where plasma and magnetic field are inseparable. Since magnetic field lines from two distinct origins cannot cross, any two populations of plasmas can never mix under the ideal MHD theory, a conclusion dissonant to known observations. Consequently, alternative plasma theories have been proposed to allow plasma mixing. Their improvements focus on capturing plasma microphysics under electron/ion inertial length scales, which is previously assumed invariant under the ideal MHD theory.

One solution is achieved by modifying the MHD equations to include the Hall electric field, which arises locally within the plasma due to differing electron and ion responses to the external magnetic field. Mathematically, the generalised Ohm's law acquires an additional Hall term (Huba, 2003)

$$\mathbf{E} + \mathbf{U} \times \mathbf{B} = \frac{1}{n_e e} \mathbf{J} \times \mathbf{B} \quad (1.52)$$

where n_e is the electron number density, e is the electron charge, and magnetic diffusivity is $\eta = 0$. This term allows decoupling of electrons and ions from the magnetic field below the ion inertial length (Eq. 1.17). As the result, the MHD induction equation becomes

$$\frac{\partial \mathbf{B}}{\partial t} = \nabla \times [(\mathbf{U} + \mathbf{U}_H) \times \mathbf{B}] \quad (1.53)$$

where $\mathbf{U}_H = -\frac{\mathbf{J}}{n_e e}$ is the Hall drift velocity. Expanding Eq. 1.53 using the $\mathbf{U}_H$ definition gives

$$\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{U} \times \mathbf{B}) - \frac{1}{n_e e} \nabla \times (\mathbf{J} \times \mathbf{B}) + \frac{1}{n_e^2 e} \nabla n \times (\mathbf{J} \times \mathbf{B}) \quad (1.54)$$

Again, we take $\eta = 0$. The Hall modification thus gives two additional terms alongside the frozen-in term. When the plasma is linearly perturbed, the first extra term corresponds to whistler-mode waves, and the second extra term produces Hall drift waves (Huba, 2003). Note that the Hall modification also introduces changes to the MHD energy equation, but this will not be covered in detail.

Another solution to the plasma mixture dilemma is to describe plasma by a completely new theoretical framework where plasma is no longer seen a fluid. The kinetic theory of plasma considers a population of charged particles (electron/ion) as a six-dimensional distribution function in spatial and velocity phase space $f(\mathbf{r}, \mathbf{v})$ where

$$N = \int f(\mathbf{r}, \mathbf{v}) d^3x d^3v \quad (1.55)$$

where N is total population number. The distribution function then obeys the Vlasov equation of motion (e.g., Cravens, 1997; Chaudhary et al., 2017)

$$\frac{\partial f}{\partial t} + \mathbf{v} \cdot \nabla f + \frac{\mathbf{F}}{m} \cdot \nabla_v f = \frac{dS}{dt} \quad (1.56)$$

where ∇ and ∇_v are gradients in spatial and velocity components respectively, and $\mathbf{F}$ is the vector field of net external forces (electric, magnetic, and gravitational) acting on the particles. For collisionless ideal plasma, the entropy function S is unchanged with time therefore $\frac{dS}{dt} = 0$. From Eq. 1.56, macroscopic plasma properties can be derived via statistical physics or moment integrations.

The kinetic theory allows direct capture of plasma processes below the scales of MHD, or when the plasma is out of equilibrium and its particle distributions are non-Maxwellian. In the case of ideal space plasma, kinetic theory allows plasma-field decoupling at length scale below the ion gyroradius, and thus plasma mixture across magnetic fields. The kinetic theory is generally considered as a more fundamental plasma framework compared to the MHD fluid description.

1.4. Energy and Particle Transport through Plasma-Magnetic Interactions

The previous section highlights the importance of understanding the interaction between plasma and magnetic field. The intimate relationship, exemplified by the imperfect Alfvén's frozen-in theorem, indicates that changes in plasma flows will have significant implications on the magnetic topology and vice versa. The next step of plasma-magnetic analysis is to understand the interaction between two different plasma-magnetic regimes across a boundary, which can result in significant energy and material transport. This section will focus on two most integral transport mechanisms in plasma – magnetic reconnection and Kelvin-Helmholtz (K-H) instability.

The process of magnetic reconnection is illustrated in Figure 10. Two magnetic fields, each with plasma frozen in, meet in an antiparallel configuration across a common boundary. The closest opposing magnetic field lines then break and reconnect as cross-boundary lines, which then depart the site via magnetic tension force. In the equivalent plasma picture, the field line breakage allows plasma detachment and diffusion across the boundary, and the reconnection produces bidirectional plasma jets that release energy from the diffusion site (Eastwood, Phan, Øieroset, et al., 2013). In Figure 10a, the antiparallel magnetic fields are

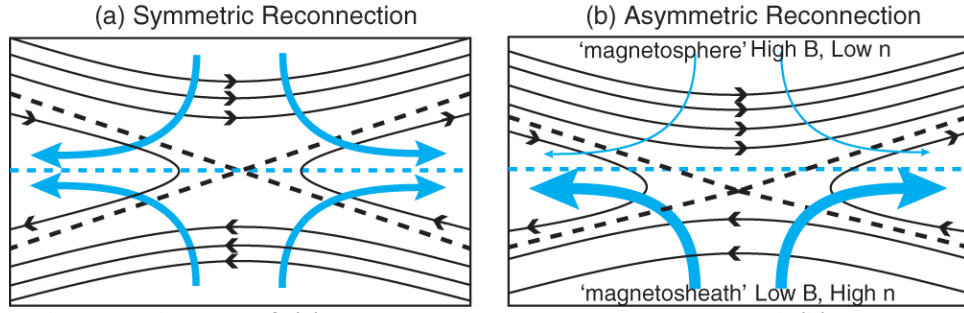


Figure 10: Schematic diagram of (a) symmetric magnetic reconnection and (b) asymmetric magnetic reconnection. Solid black lines indicate magnetic fields. Dashed black lines indicate reconnection ‘separatrix’ dividing open and closed field lines. Dashed blue line indicates the boundary between two plasma-magnetic regimes. Finally, solid blue lines show plasma diffusion paths. Image courtesy of Eastwood, Phan, Øieroset, et al., (2013), *Institute of Physics Publishing*.

equal in strength, so their reconnection is symmetric and the departing plasma jets (blue arrows) contain equal contributions from the two mixing populations. In Figure 10b, one magnetic field is significantly stronger than the other. Reconnection in this case is asymmetric and the departing jets contain highly unequal plasma contributions i.e., there is net plasma diffusion across the boundary.

In the more general resistive MHD picture, magnetic reconnection can occur as plasma diffusion is permitted in the MHD induction equation (second term of Eq. 1.40). However, reconnection is still abundant in ideal plasma, whose negligible resistivity means the MHD diffusion term is always effectively zero. Ideal (or collisionless) reconnection becomes viable when the boundary between the two plasma-magnetic regimes become thinner than one ion inertial length (Eq. 1.17) (e.g., Cassak et al., 2006; Phan et al., 2011). At this point, plasma density and velocity become varied within a MHD volume element, violating the locality assumption that underpins the fluid theory. Alfvén’s theorem is no longer applicable and collisionless plasmas can cross the boundary through a small ‘diffusion region’ (e.g., Vasylūnas, 1975).

The first model of magnetic reconnection was conceived by Sweet and Parker, in which single bulk plasma fluid flows through a long and narrow diffusion region (Sweet, 1956; Parker, 1957). However, such reconnection process occurs too slowly to explain solar flare phenomena observed at the time without an anomalous resistivity appearing in the plasma (Parker, 1963). Around similar time period, an alternative model of reconnection was proposed by Petschek, in which the diffusion region is square-like and fluid plasma crosses

the boundary via four slow-mode shocks surrounding the diffusion region (Petschek, 1964). Both adjustments decrease the reconnection timescale enough for the process to be physically useful, but the Petschek model still relies on large anomalous resistivity that remains poorly understood (Cassak et al., 2006).

A fast magnetic reconnection mechanism with narrow diffusion region is achieved by the modern Hall model, in which plasma electrons and ions are considered as separate fluids. Due to the electron's small mass, it has a much shorter inertial length and hence a smaller diffusion region (Birn et al., 2001). This allows electrons to remain a frozen-in fluid as the ions decouple from the magnetic field, pulling the field lines away from the boundary to form a quadrupolar Hall electric field structure (Birn et al., 2001; Cassak et al. 2006), which then rapidly accelerates reconnection outflow through whistler wave mechanism (Mandt et al., 1994). Consequently, the outflow velocity is independent of plasma resistivity and the size of ion diffusion region. The Hall reconnection signatures have since been observed empirically near Earth and seen in numerical simulations (Øieroset et al., 2001; Mozer et al. 2002; Cassak et al., 2006).

The Kelvin-Helmholtz (K-H) instability describes turbulent motions that can arise when a boundary, across which exists a velocity shear between two fluids, is perturbed by an external force (e.g., Kelvin, 1910). As illustrated in Figure 11, the instability is divided into two distinct phases – the linear phase in which local perturbations produce steepened boundary waves, followed by the nonlinear phase in which the boundary rolled up into vortices. Originally a hydrodynamic phenomenon, K-H instability is later adopted for early descriptions of near-Earth plasma behaviours (Dungey, 1955; Southwood, 1968), and then theoretically developed as an MHD interaction over the latter half of the 20th century (e.g., Southwood & Hughes, 1983).

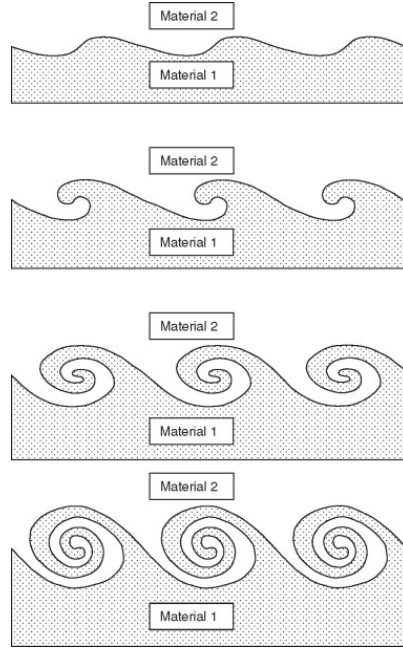


Figure 11: An illustrative diagram of K-H instability growth, showing steepening linear wave (top two) followed by development of nonlinear vortices (bottom two). Image courtesy of Grujicic, (2007).

While linear K-H waves can impart energy-momentum across plasma-magnetic boundary via their motions (e.g., Axford, 1964; Pu & Kivelson, 1983b, Southwood & Hughes, 1983), plasma particle transport occurs more prevalently through nonlinear K-H vortices, which have been studied extensively through numerical MHD simulations (e.g., Thomas & Winske, 1993; Fujimoto & Terasawa, 1994; Nykyri & Otto, 2001; Matsumoto & Hoshino, 2004). A K-H vortex can facilitate plasma and energy transport in multiple ways including particle mixing from turbulent vortex decay (Nakamura et al., 2004; Matsumoto & Hoshino, 2004, 2006), local magnetic reconnection from wounding motion of the vortex (Nykyri & Otto, 2001; Nakamura et al., 2006, 2008), and coupling between the vortex and local Alfvén waves (Chaston et al., 2007). As nonlinear K-H instability studies are largely based on simulation of MHD equations, the point of transition between linear and nonlinear instability phases remains unclear (and likely system dependent).

While K-H instability arises from fluid nature of plasma, the phenomenon cannot be entirely captured within the MHD framework. Local particle motions can have significant influences on macroscopic properties of the instability including its growth rate. The kinetic aspect of K-H instability is particularly exemplified by the finite Larmor radius (FLR) effect, which will be discussed in detail in Chapter 5 of this thesis.

In this chapter, we traverse millennia's worth of observations on Ganymede, until the moon was found to generate a permanent internal magnetic field. The discovery opens a new avenue of research on magnetism and plasma, whose fundamental properties and relationships with each other are outlined in subsequent sections. Given limited space, it is not possible to expand every concept in detail. However, the brief explanations should provide sufficient overview, upon which we can build an understanding of Ganymede's unique magnetic environment.

Chapter 2: Magnetospheres of the Solar System

In the previous chapter, I have laid groundwork on large-scale magnetism, plasma in space, and the intimate relationship between the two. Now, the discussion will move onto real space environment, and how plasma-magnetic interactions significantly shape conditions around celestial bodies through the concept of magnetosphere. Our journey will first travel through Solar System planets, whose studies have developed the bulk of current magnetospheric knowledge, before ending at Ganymede, whose magnetic environment shall be studied in the upcoming work chapters.

2.1. Magnetic Domain

A magnetic field, alone in empty vacuum, can exert its effects over infinite distance. However, this is not true in real space, where numerous magnetic fields compete for influence, meeting along pressure boundary lines without crossing due to Alfvén's frozen-in theorem. Consequently, space can be divided into magnetic domains. In each domain, charged particle motions are dominantly governed by one specific magnetic field. A domain belonging to a magnetic celestial body is thus defined as the body's *magnetosphere*.

A magnetosphere is typically confined by an external magnetic field carried by a flow of plasma around the magnetic body. As plasma flows in one direction, the magnetosphere is compressed on the upstream side and elongated on the downstream side in what is known as the magnetotail. On the upstream side, the magnetospheric field meets the external field along a boundary known as the magnetopause, at which exists a current sheet arising from gradient between the magnetic fields (per Ampère's Law). The current sheet generates its own Chapman-Ferraro (C-F) magnetic field, whose magnitude cancels out the magnetospheric field outside the magnetopause, resulting in domain confinement (Chapman & Ferraro, 1940). Simultaneously, the C-F field adds to the magnetospheric field just inside the boundary, resulting in magnitude compression by a factor of ~ 2 . Effects of the C-F magnetic field is schematically shown in Figure 12.

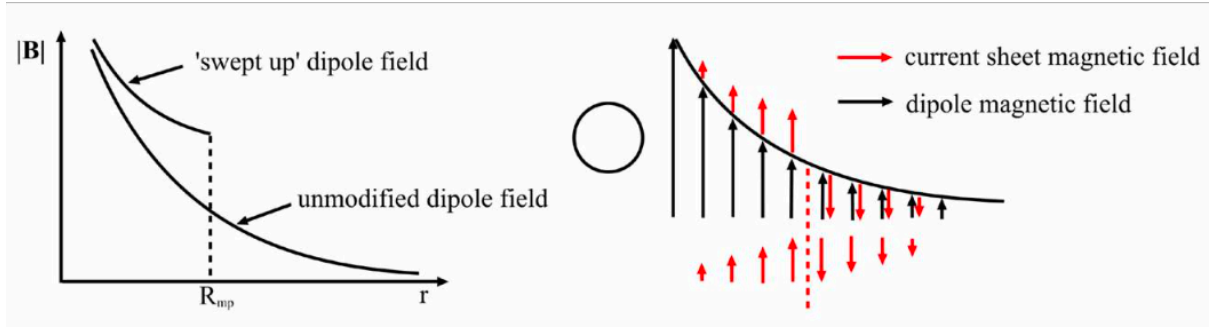


Figure 12: (Right) Chapman-Ferraro (C-F) magnetic field confines a magnetosphere by removing its field beyond the magnetopause (red dashed line) and enhancing its field just before it, resulting in compressed (or 'swept up') magnetospheric dipole on the upstream side (left). Image courtesy of Eastwood, (2018).

At the magnetopause, the total pressure exerted by the two adjacent magnetic fields, and by their frozen-in plasmas, must be balanced i.e.,

$$P_{B,1} + P_{p,1} + P_{r,1} = P_{B,2} + P_{p,2} + P_{r,2} \quad (2.1)$$

where P_B , P_p , and P_r are magnetic, plasma, and dynamic ram pressures on the external (1) and magnetospheric (2) sides of the magnetopause. From Eq. 2.1, it is possible to calculate the standoff distance, or the maximum extent that the magnetosphere protrudes into the upstream. To illustrate, if the magnetospheric field is assumed to carry cold and stationary plasma (so exerting no plasma and dynamic pressures), then we can write

$$P_{\text{ext}} = P_{B,2} = \frac{B_{\text{MP}}^2}{2\mu_0} \quad (2.2)$$

where $P_{\text{ext}} = P_{B,1} + P_{p,1} + P_{r,1}$ is the total external acting pressure, B_{MP} is the magnetospheric field strength at the magnetopause, and $\mu_0 = 4\pi \times 10^{-7} \text{ NA}^{-2}$ is the permeability of free space. Then, if the magnetospheric field is assumed to be dipolar with rotation-aligned dipole axis, then the magnetopause field strength at the body's equator is

$$B_{\text{MP}} = 2|g_1^0| \left(\frac{R_P}{r_{\text{MP}}} \right)^3 \quad (2.3)$$

where the pre-factor 2 accounts for compression due to the C-F field. The remainder of the expression is simplified from Eq. 1.8 (with $\theta = 0^\circ$). Therefore, the magnetopause standoff distance is

$$r_{\text{MP}} = R_p \left(\frac{2|g_1^0|^2}{\mu_0 P_{\text{ext}}} \right)^{\frac{1}{6}} \quad (2.4)$$

The location on the magnetopause with distance r_{MP} from the body's centre is defined as the subflow point or the local noon of the magnetosphere. The magnetopause extends equatorially around the magnetic body toward the flank regions where the magnetosphere transitions from upstream to downstream.

The magnetopause also extends latitudinally away from the equator. If the external plasma flow is super-Alfvénic (i.e., faster than local Alfvén waves), then its dynamic ram pressure becomes more important and the magnetopause quickly wraps around the body's polar (lobe) regions, resulting in a bullet-shaped magnetosphere (e.g., Neubauer, 1990). However, if the external flow is sub-Alfvénic, then its magnetic pressure becomes more important and the magnetopause extends almost orthogonal to the incoming plasma flow, resulting in a cylindrical (or Alfvén wing) magnetosphere (Neubauer, 1998). Otherwise, if the external flow is supersonic, then the magnetopause is preceded externally by the bow shock, where inflow plasma is slowed below the sound speed by the approaching boundary and forms a 'magnetosheath' layer. Since the magnetopause is established upon a delicate pressure balance, its location is dynamic under temporal variations of competing pressures, and vulnerable to transient disruptions from energy transport processes.

2.2. Planetary Magnetospheres

In the Solar System, every planet outside of Venus and Mars possesses internally generated magnetism, and thus produces their own magnetospheres. Planetary magnetospheres are confined externally by the solar wind, a continual stream of plasma particles released by interactions on the Sun's surface (Cravens, 1997). As mentioned in Chapter 1, the solar wind is strongly diffuse ($n_i \sim 1 - 10 \text{ cm}^{-3}$), possesses high temperature ($T_i \sim 10^{5-6} \text{ K}$), and radiates

outward from the Sun at super-Alfvénic/supersonic velocity (200 – 800 km/s) across the Solar System. The plasma stream contains primarily single-charged protons H^+ and electrons. It also carries the interplanetary magnetic field (IMF), whose orientation becomes increasingly tightly wound with heliocentric distance (the Parker spiral (Parker, 1958a)). The solar wind's deflections around magnetised planets produce bullet-shaped magnetospheres, whose diverse structures will be explored in this section.

The magnetosphere of our own Earth is shown in Figure 13 with the solar wind flowing rightward from the upstream dayside. The external compression produces a magnetopause with the subflow (or subsolar) standoff distance of $r_{MP} \sim 10 R_E$, where $R_E = 6,371$ km is the Earth's radius. The typical IMF near Earth has the strength of ~ 5 nT (Wilson III, et al., 2018), compared to the terrestrial field strength of ~ 60 nT (from Eq. 2.3 with $g_0^1 \sim 31,000$ nT). Inside the magnetopause, closed magnetospheric field lines are found near the equator, converging at northern and southern cusp points. Beyond the cusps, magnetic field lines become open as they wrap over the poles and extend into the nightside magnetotail, which can extend up to $\sim 100 R_E$ (Walker et al., 1999).

The dayside magnetopause, under frozen-in ideal MHD description, is a closed shield that protects the Earth from external particle streams. However, in-situ observations have demonstrated an abundance of solar wind particles within Earth's magnetosphere, thus indicating particle injections across the magnetopause, primarily through magnetic reconnection and K-H instability (e.g., Welling et al., 2015). The prevalence of each transport

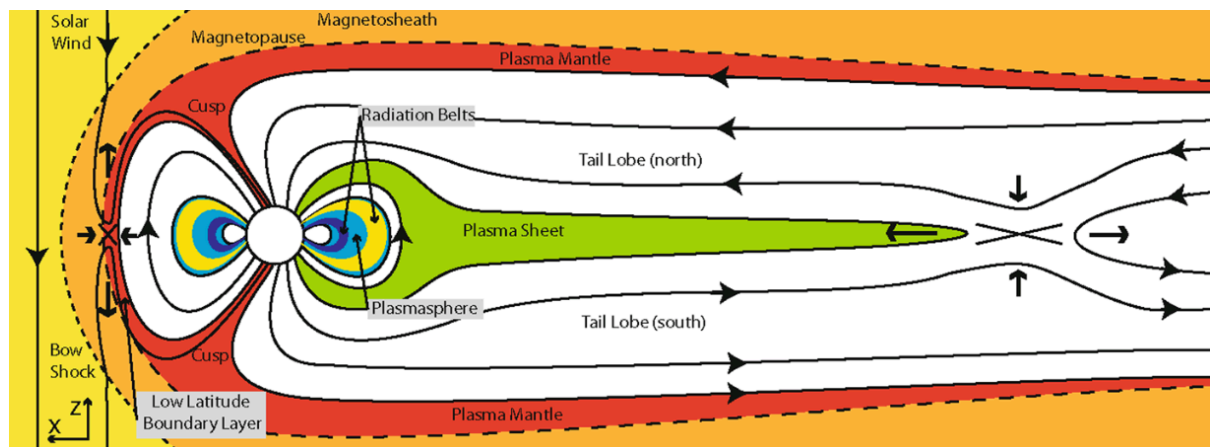


Figure 13: A schematic diagram of Earth's magnetosphere. Image courtesy of Eastwood et al., (2015), *Space Science Review*.

mechanism is determined by the IMF orientation. When the IMF is southward (as in Figure 13), the magnetic configuration at the dayside magnetopause becomes strongly antiparallel⁷, facilitating magnetic reconnection events (e.g., Trattner et al., 2007a, 2007b) that disrupt growth of K-H instability (Nakamura et al., 2020). When the IMF is northward, the magnetopause magnetic configuration becomes parallel, reducing the likelihood of reconnection. Pressure pulses from the solar wind also results in stronger undulation of the boundary, which generate more favourable conditions for K-H instability as a competing energy-particle transport interaction (Miura, 1995).

When the two mechanisms are compared, magnetic reconnection is found to have much greater impact on Earth's magnetospheric dynamics (Phan et al., 2000; Milan et al., 2007). Given the Earth's relatively slow rotation compared to plasma flow speeds, and the planet's lack of internal plasma sources, plasma convection around the planet follows the Dungey cycle (Dungey, 1961), which starts from reconnection on the dayside magnetopause (marked X on the left side of Figure 13). Open magnetic field lines depart from the reconnection site, carrying (formerly solar wind) plasma into the magnetotail. This builds up two plasma populations in northern and southern tail lobes with antiparallel magnetic fields, which induces another reconnection site in the magnetotail (marked X on the right side of Figure 13). The reconnection will send some plasmas out of the magnetosphere, and return others toward Earth along an equatorial plasma sheet (green in Figure 13) that connects to the Earth's polar regions. Subsequent near-Earth plasma convection set up electric currents aligned with local magnetic field lines, which can generate visible auroras in the polar night skies (Akasofu, 1964).

As part of our home planet, Earth's magnetosphere has been extensively explored by numerous space missions. A list of notable spacecraft with explicit magnetospheric research objectives include

1. International Earth-Sun Explorer (ISEE) 1 and 2 (NASA–ESA, launch year 1977)

⁷ The internal magnetospheric field is northward at the dayside magnetopause due to the southward-pointing terrestrial dipole axis

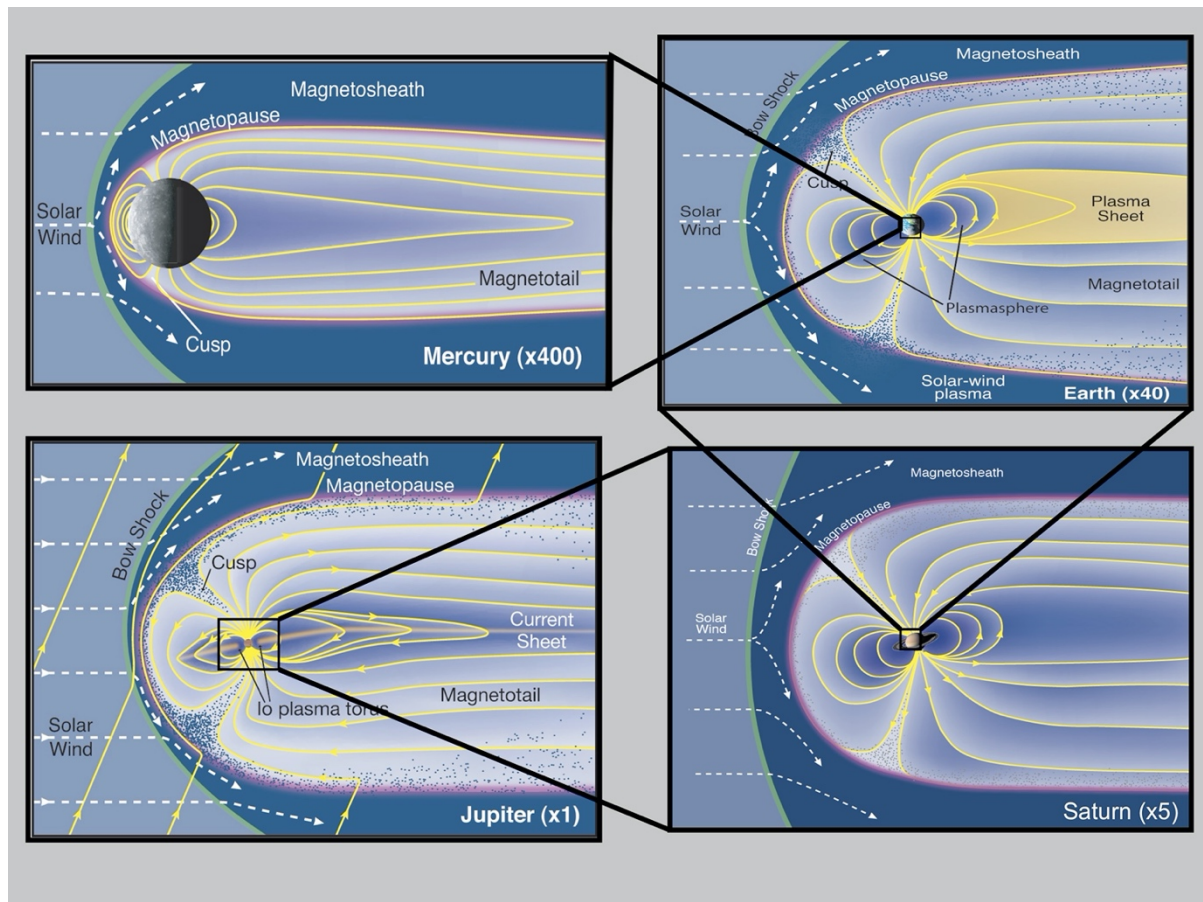


Figure 14: Relative illustrations for magnetospheres of Mercury (top-left), Earth (top-right), Jupiter (bottom-left), and Saturn (bottom-right). Image courtesy of Bagenal & Bartlett, *Laboratory for Atmospheric and Space Physics, University of Colorado, Boulder*.

2. Geotail (ISAS⁸ – NASA, 1992)
3. WIND/Polar (NASA, 1996 – 2008)
4. Cluster II (ESA, 2000)
5. Time History of Events and Macroscale Interactions during Substorms (THEMIS) (NASA, 2007)
6. Van Allen (NASA, 2012)
7. Magnetospheric MultiScale (MMS) Mission (NASA, 2015)

While ISEE 1/2 and Geotail, and Polar are individual spacecraft (Ogilvie et al., 1977; Acuna et al., 1993; Nishida, 1994), subsequent missions operate through synchronised groups of spacecraft – four for Cluster II, five for THEMIS, two for Van Allen, and four for MMS Mission (Escoubet et al., 2001; Angelopoulos, 2009; Fox & Burch, 2013; Burch et al., 2016). Synchronised observations allow pinpoint triangulation and thus detailed mapping of smaller

⁸ Institute of Space and Aeronautical Science (Japan)

physical structures. The ongoing MMS mission, in particular, seeks to understand the physics of magnetic reconnection below the MHD length scale (Burch et al., 2016).

Figure 14 illustrates the other notable planetary magnetospheres in the Solar System and their relative sizes compared to Earth's. The magnetosphere that is closest to the Sun is that of Mercury (top-left). Similar to Earth's, the Hermean magnetosphere is bullet-shaped and the plasma convection within is driven by the Dungey cycle (Slavin, Imber, et al., 2012). The two systems have some notable differences, however. The dynamo-generated dipole magnetic field of Mercury is much weaker, possessing only 1.1% the strength of Earth's (Russell, 1993). This, alongside one magnitude stronger solar wind pressure, means the Hermean magnetosphere is much smaller, and highly compressed on the dayside with standoff distance $r_{MP} \sim 1.5 R_M$, where $R_M = 2,440$ km is Mercury's radius (Slavin, 2004; Winslow et al., 2013). In fact, an Earth-like dynamo field alone cannot hold the magnetopause at its observed location (e.g., Christensen, 2006). More recent studies suggest that the Chapman-Ferraro magnetic field from the Hermean magnetopause current sheet induced a secondary field from within Mercury's electrically conducting interior, which strengthens the magnetospheric field (e.g., Saur et al., 2010; Heyner et al., 2011).

The Hermean magnetosphere was first observed during brief 1974-1975 flybys of NASA's Mariner 10 (Dunne & Burgess, 1978). However, the bulk of in-situ observations arrived a few decades later during NASA's MErcury Surface, Space ENvironment, GEochemistry, and Ranging (MESSENGER) mission, which performed 4,105 orbits around the planet between 2011-2015 (Bedini, 2017). Later this decade, another spacecraft BepiColombo (ESA-JAXA⁹) will arrive at Mercury and perform more comprehensive studies of the planet's magnetic environment (Rothery et al., 2020).

The bottom diagrams of Figure 14 illustrate voluminous magnetospheres of the Gas Giants. Of the two, Saturn (bottom-right) has a smaller magnetosphere at eight times the size of Earth's. The Saturnian internal magnetic field is strongly dipolar, its axis deviating less than 1° from the planet's rotation axis, with an equatorial field strength $g_0^1 \sim 20,000$ nT (Russell,

⁹ Japanese Aerospace eXploration Agency

1993; Burton et al., 2008). Although Saturn possesses a weaker magnetic field than Earth's, the solar wind pressure is two orders of magnitude weaker at 9.5 AU, resulting in larger standoff distance $r_{MP} \sim 21 R_S$, where $R_S = 60,330$ km is the Saturnian radius (Russell, 1993). Since Saturn's dipole axis is parallel to its rotation axis, the magnetospheric field at the magnetopause points southward (opposite from Earth's case), so magnetic reconnection is more prevalent when the IMF is northward.

Unlike Earth and Mercury, Saturn's magnetosphere contains a significant internal plasma source – the moon Enceladus ($\sim 4 R_S$ from Saturn), which continuously emits water group ions at rate ~ 100 kg/s (e.g., Tokar et al., 2005). Furthermore, Saturn rotates much faster at 10.6 hours per revolution, hence rotational effects become significant when considering plasma motion. Enceladus' plasma is carried radially outward and around Saturn, forming a dominant plasma sheet (torus) centred by the planet's centrifugal equator (Gombosi et al., 2009; Mauk et al., 2009). The plasma sheet causes significant distensions of the equatorial magnetic field (via Alfvén's theorem) from its quasi-dipolar configuration, leading to a disc-like magnetic structure in the outer magnetosphere ($\sim 15 - 20 R_S$) known as the magnetodisc. Plasma convection in Saturn's magnetosphere is heavily influenced by the rotational Vasyliūnas cycle especially in low-latitude regions (e.g., Vasyliūnas et al., 1983). However, Dungey-like circulation is also expected at high-latitudes and the outer magnetosphere, where magnetic reconnection continues to inject external solar wind plasma (Cowley et al., 2004; Masters et al., 2012; Southwood & Chané, 2016).

Four unmanned spacecraft have previously visited Saturn's magnetosphere. Three individual flybys were made by Pioneer 11 (1979), Voyager 1 (1980), and Voyager 2 (1981) (Kohlhase & Penzo, 1977; Siddiqi, 2018). But like Mercury, most extensive observations of Saturn were made by one orbiting spacecraft. The Cassini-Huygens probes arrived at Saturn in 2004 as part of the collaboration between NASA, ESA, and the Italian Space Agency (ASI). While Huygens immediately landed on the moon Titan, Cassini would make 294 orbits around Saturn until its mission termination in 2017 (Ingersoll, 2020).

The largest gas giant magnetosphere – Jupiter's (bottom-left, Figure 14) – is the most voluminous planetary magnetic domain in the Solar System. The Jovian dynamo magnetic

field has equatorial strength $g_0^1 \sim 400,000$ nT, about 13 times stronger than Earth's (Russell, 1993). The solar wind pressure at 5.2 AU is an order weaker than at Earth's. Even so, the Jovian magnetopause standoff distance $r_{MP} \sim 70 - 100 R_J$, where $R_J = 71,492$ km is the Jovian radius, is beyond the pressure capability of the internal field alone. The distant magnetopause is due to pressure provided by a large, rotating plasma population inside Jupiter's magnetosphere.

Like Saturn, the Jovian plasma has an internal origin. Volcanoes on the inner moon Io emit heavy ions at rate ~ 300 kg/s (Crary et al., 1998; Schneider & Bagenal, 2007). The volcanic plasma consists mostly of O^+ ions due to longer oxygen ionisation lifetime, with smaller concentrations of S^{++} , O^{++} , and S^{+++} . The plasma is then convected following the Vasyliūnas cycle, in which magnetic tension force accelerates the plasma such that it co-rotates with Jupiter i.e.,

$$U = \frac{2\pi r}{T} \quad (2.5)$$

where U is the plasma's bulk flow speed, r is the plasma's radial distance from Jupiter, and $T = 9.6$ hours is the planet's rotation period. This magnetic forcing is imperfect, however, due to non-idealised factors including finite conductivity of Jupiter's ionosphere. Consequently, as the plasma moves toward the planet's middle and outer magnetosphere, its flow speed drops below the values required for co-rotation (McNutt et al., 1979; Russell & Dougherty, 2010).

Similar to the case of Saturn, convection of Io's emitted plasma produces a $\sim 3 R_J$ thick sheet (torus) in Jupiter's centrifugal equator (Khurana et al., 2004; Kivelson et al., 2004). Net outward diffusion in the sheet stretches Jupiter's mostly dipolar magnetic field lines, forming a magnetodisc structure at medium and large distances away from the planet ($r > \sim 20 R_J$) (Khurana et al., 2004; Russell & Dougherty, 2010). Magnetospheric flattening brings together opposing magnetic field lines in the centrifugal plane, forming a very strong (>160 MA) current sheet of $\sim 1 R_J$ thickness embedded at the centre of the plasma sheet (Khurana et al., 2004). Plasma and magnetic conditions are vertically differentiated with scale height $\sim 1 R_J$

from the sheet's centre (Kivelson et al., 2004). As absolute distance increases, the Jovian plasma becomes less dense and exerts weaker pressure, but the magnetic field it carries becomes stronger. As Ganymede never orbits beyond $3 R_J$ from the Jovian centrifugal equator, the moon's magnetosphere is always enveloped and influenced by conditions in the plasma sheet. Properties of Jovian plasma and magnetic field near Ganymede will be covered more specifically in Chapter 2, Section 2.3.

Among planetary magnetospheres, Jupiter's magnetic domain is the most visited after Earth's. The list of missions that have provided in-situ observations includes

1. Pioneer 10/11 (NASA, arrival year 1973/1974)
2. Voyager 1/2 (NASA, 1980/1981)
3. Ulysses (NASA, 1992)
4. Galileo (NASA, 1995)
5. Cassini-Huygens (NASA-ESA-ISA, 2000)
6. New Horizons (NASA, 2007)
7. Juno (NASA, 2016)

Most of the missions are individual flybys, with two orbiters in Galileo and Juno. Galileo made 33 orbits around the planet between 1995-2003, as its instruments conducted detailed observation of the equatorial magnetosphere (Meltzer, 2007). Meanwhile, Juno focuses on the magnetosphere's polar regions, and is still in operation with 76 planned orbits and increasing (NASA, 2009; Siddiqi, 2018). Some of the listed missions have also visited Ganymede and its magnetosphere, as discussed in Section 1.1.

Beyond the gas giants, Uranus and Neptune also generate their own magnetospheres which are much less well-known. With only one flyby encounter from Voyager 2 in 1986 for Uranus and 1989 for Neptune respectively (Siddiqi, 2018), knowledge of the ice giants' magnetospheres is derived primarily through simulations (e.g., Voigt & Ness, 1990; Tóth et al., 2004; Mejnertsen et al., 2016; Cao & Paty, 2020), alongside a number of Earth-based remote observations (Lamy, 2020).

The magnetic domains of Uranus and Neptune are complex compared to those of gas giants due to (a) significant tilts of the dipole axis orientations from the rotation axes and (b) large contributions from higher-degree magnetic fields i.e., quadrupoles and beyond (Holme & Bloxham, 1996). The two magnetospheres thus possess many peculiar features including (but not limited to) trapped plasmas on the dayside, cusp indentations, magnetic reconnection modulated by planetary dipole axis oscillations, and aurora ovals far away from geographical poles (Ness et al., 1986; Mejnertsen et al., 2016; Cao & Paty, 2020; Lamy, 2020). The Uranian and Neptunian magnetopauses have standoff distances of $\sim 18 R_U$ and $\sim 4 R_N$ respectively where $R_U = 25,362$ km and $R_N = 24,622$ km (Cao & Paty, 2020; Mejnertsen et al., 2016). Little is known regarding plasma sources internal to these magnetospheres, although Triton might be a significant contributor in the Neptunian system (Hill & Dessler, 1990; Masters et al., 2014). Further details regarding Uranus' and Neptune's magnetic domains will not be covered in this thesis.

With the tour of planetary magnetospheres completed, there is only one magnetic domain we have not discussed in the Solar System. The next section will discuss the detailed structure magnetosphere of Ganymede, how it uniquely interacts with the Jovian plasma environment, and how it drives the particle and field environment around the Solar System's largest moon.

2.3. Ganymede's Magnetosphere

At radius $R_G = 2,643$ km, Ganymede is slightly larger than the planet Mercury. The largest Solar System moon is tidally locked with Jupiter, moving around the planet with period of 7.16 days, in an almost circular orbit with semi-major axis distance $15 R_J$ in Jupiter's ecliptic plane (Jacobson, 2003; Bills, 2005). As discussed in Chapter 1, Galileo flybys of Ganymede provide empirical evidence of a permanent magnetic field (Gurnett et al., 1996; Kivelson et al., 1996). The moon's internal field is close to dipolar with a magnetic axis that tilts 176° latitudinally from the rotation axis, and 24° longitudinally from the Jupiter-facing meridian (Kivelson et al., 2002).

A cross section of Ganymede's internal structure is illustrated in Figure 15. Note that thickness values shown are rough estimates. Ganymede's primary dipole magnetic field is likely

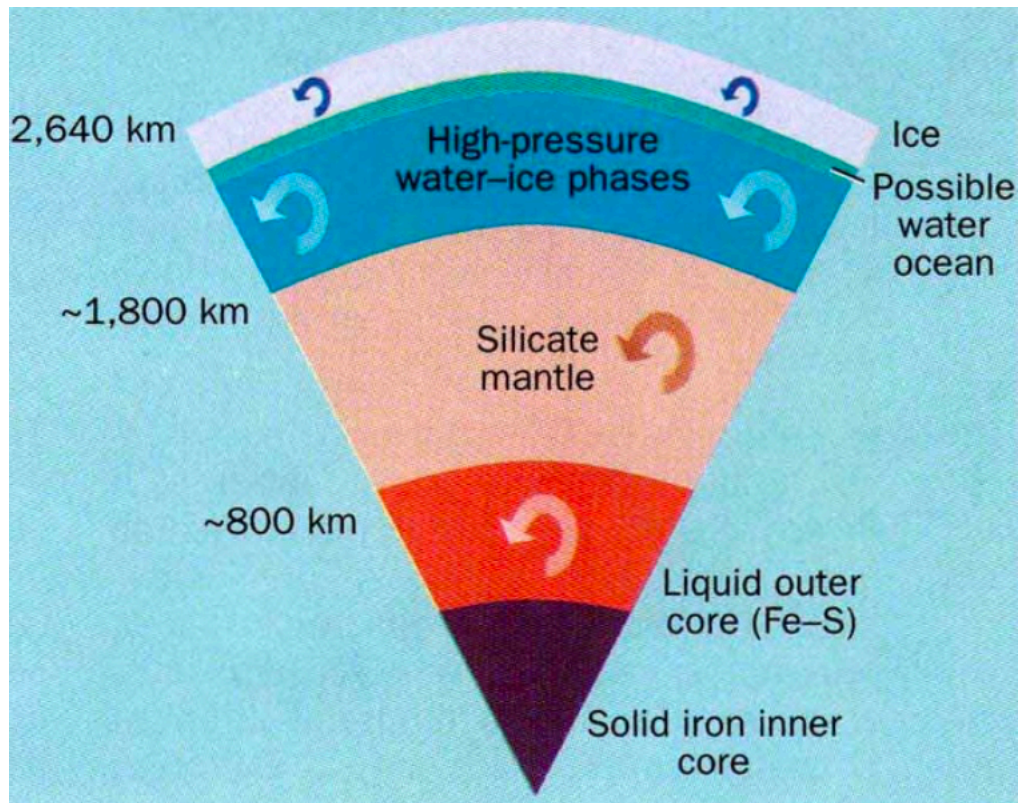


Figure 15: Cross-section of Ganymede's interior. Depth values are rough estimates. Image courtesy of Stevenson, (1996), *Nature*.

produced via dynamo action in a molten iron-sulphide outer core, which encases a solid iron inner core and is surrounded by a silicate mantle (Anderson et al., 1996; Schubert et al., 1996). Earth-based radio observations suggest a combined core thickness between 400 – 1,300 km and a mantle thickness of about 1,000 km (Anderson et al., 1996). The silicate is then encased by an 800 – 900 km thick frozen ice layer, which is partially melted into saltwater at greater depths due to high internal pressure and thermal conduction from the core, forming a subsurface ocean (Schubert et al., 1996; Stevenson, 1996; McKinnon, 1997).

Between Galileo flybys, Ganymede's dipole axis was found to wobble by $\sim 3^\circ$. This fairly rapid change may be produced by non-negligible higher degree moments in the dynamo field (e.g., quadrupole). Or more likely, a subsurface conducting ocean under the dynamo field induces a secondary magnetic field whose dipole is orthogonal to the primary dipole (Kivelson et al., 2002). The depth of this subsurface ocean is currently uncertain. Kivelson et al., (2002) propose a fairly thin ~ 1 km water layer sandwiched between two thick ice shells locating ~ 150 km from Ganymede's surface. However, the liquid layer can be significantly thicker, penetrating into the high-pressure water-ice phase layer in Figure 15. At Ganymede's surface,

the primary dynamo-generated dipole is much stronger than the secondary induced dipole, with equatorial field strengths of 719 nT and ~ 50 nT respectively (Kivelson et al., 1998; Kivelson et al., 2002). Consequently, the secondary dipole effects can be largely neglected for global-scale and steady-state study of Ganymede's magnetosphere.

During Ganymede's orbit of Jupiter, the moon maximally deviates from the planet's centrifugal equator by a distance of

$$d_{\max} = (15 R_J) \sin(6.3^\circ) \sim 1.6 R_J \quad (2.6)$$

where 6.3° is the approximate tilt between Jupiter's centrifugal and spin (rotational) equators. Given that the Jovian plasma sheet is $\sim 3 R_J$ thick with continuous boundaries, Ganymede's magnetosphere should always be submerged in the Jovian plasma, whose composition is dominated by heavy oxygen (mostly O^+) and sulphur (mostly S^{++}) ions with $<10\%$ protons (Bagenal et al., 2016). The Jovian plasma near Ganymede is diffuse ($n_i \sim 4 \text{ cm}^{-3}$) and warm ($T_i \sim 60 \text{ eV}$) (Kivelson et al., 2004), flowing at 80% the speed of Jupiter co-rotation i.e., $U \approx 150 \text{ km/s}$ (Williams, Mauk, McEntire, Roelof, et al., 1997). Subtracting Ganymede's Keplerian velocity, the flow speed is reduced $U \approx 140 \text{ km/s}$ in Ganymede's rest frame, which is sub-Alfvénic and sub-sonic (Jia et al., 2008). The external environment is a sharp contrast from the proton-dominated, super-Alfvénic, and super-sonic solar wind around planetary magnetospheres, making Ganymede a unique plasma laboratory in the Solar System.

Until most recently, close in-situ observations on Ganymede are limited to six Galileo flybys. Hence, our knowledge of the moon's magnetosphere is obtained through modelling simulations. Figure 16 shows a model of the magnetosphere in Ganymede-centred Cartesian coordinate system (GphiO), where the X axis points along the incoming Jovian flow and is parallel to Jovian equatorial plane, the Y axis points from Ganymede to Jupiter, and the Z axis is parallel to Jupiter's (and approximately Ganymede's) rotation axis. The magnetosphere is $\sim 6 R_G$ in width and $\sim 8\text{--}10 R_G$ in length (Jia, Kivelson, et al., 2010). Despite its slow flow, the Jovian plasma carries a strong magnetic field $\sim 100 \text{ nT}$ in magnitude that is always southward (Kivelson et al., 1997). Thus, the significant magnetic pressure compresses the upstream

magnetopause to standoff distance $1.65 - 2 R_G$ (Kivelson et al., 1998; Jia et al., 2008, Jia et al., 2009).

Compression by sub-Alfvénic plasma results in a cylindrical magnetosphere with an extended tail structure known as the Alfvén wing (Neubauer, 1998). The bending angle of Ganymede’s magnetic field in response to the Jovian plasma flow can be quantified by (Kivelson et al., 2004)

$$\theta_A = \tan^{-1} \left(\frac{U}{C_A} \right) \quad (2.7)$$

where U is the Jovian plasma flow speed at Ganymede and C_A is the local Alfvén velocity (Eq. 1.48). Taking the Jovian magnetic field strength $B_0 \sim 100$ nT and plasma density $n_i \sim 4$ cm⁻³, the Alfvén velocity is $C_A \sim 290$ km/s, and thus the Alfvén wing bends by $\theta_A \sim 25^\circ$. This angle assumes a straight Alfvén wing that follows the orange dashed line in Figure 16. In reality, the Alfvén wing bending is small near the upstream magnetopause and accelerates beyond the moon’s polar regions.

Ganymede’s magnetosphere can be sectioned into closed and open magnetic field regions. Near the equator, magnetospheric field lines are closed, connecting the moon’s magnetic

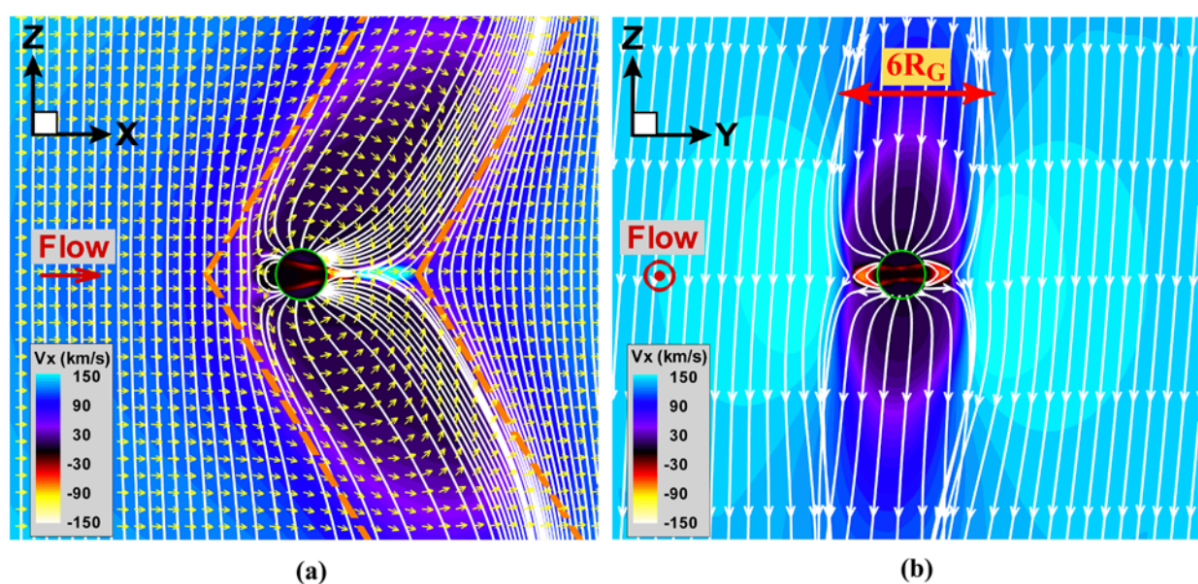


Figure 16: Ganymede’s magnetosphere (a) the sideview on X-Z plane and (b) the inflow view on Y-Z plane of GphiO coordinates. Image courtesy of Jia, Kivelson, et al., (2010), *Journal of Geophysical Research: Space Physics*.

poles and pointing northward ($Z > 0$) through the equatorial plane. Near the poles and in the magnetotail, however, fields lines are open. They connect Ganymede to Jupiter and points southward ($Z < 0$) into or away from the moon. The open-closed field boundaries (OCFB) form two rings encircling Ganymede's magnetic poles at $Z \sim \pm 0.6 R_G$. Under MHD theory, open magnetic field lines can carry plasma into and out of Ganymede's magnetosphere (Frank et al., 1997; Williams, Mauk, McEntire, 1997, Williams, Mauk, McEntire, Roelof, et al., 1997). However, plasma movement around Ganymede likely follows a Dungey-like circulation (Collinson et al., 2018), where convection is primarily driven by plasma injection across the closed upstream magnetopause through magnetic reconnection and/or K-H instability. The circulation is reflected through the plasma velocity colour map in Figure 16a, which shows downstream-ward flows over Ganymede's polar caps, a flow speed peak near the expected magnetotail reconnection site, and upstream-ward flow back to Ganymede along the downstream equatorial closed magnetic field region.

In the years since Ganymede's magnetic field was discovered, several attempts have been made to model the moon's magnetosphere. Kivelson et al., (1998) first considers the cylindrical magnetosphere as an ovoid column with shifting centre points, which account for the bending Alfvén wing structure. However, the analytical description poorly describes the downstream magnetosphere, whose elongation distorts the domain from the ovoid shape. In subsequent years, our modelling of Ganymede's magnetosphere turns toward numerical simulations. The first code schemes for Ganymede were adopted from plasma flow study around the Saturnian moon Titan (Kopp & Ip, 2001, 2002; Ip & Kopp, 2002). The model considers Jovian plasma as a single fluid and magnetospheric responses to the external environment, but does not explicitly evaluate boundary effects from cross-magnetopause energy transport or processes in the moon's interior. A few years later, multi-fluid resistive MHD models, in which electrons and ions are treated as separate fluids, were adopted for Ganymede (Paty & Winglee, 2004, 2006; Paty et al., 2008). These studies focus on ion flows and energisation in the magnetosphere, but with limitations in common with the early single-fluid studies above.

The first Ganymede model that captures boundary interactions is provided by Jia et al., (2008, 2009), whose outputs have been shown in Figure 16. The three-dimensional global

magnetospheric model considers plasma as a single *resistive* MHD fluid despite its ideal condition. This allows diffusion-related terms in the MHD equations to predict, with adequate accuracy, locations of magnetopause interactions that actually occur through kinetic microphysics. Resistive MHD models are generally favoured over full kinetic simulations, which are very computationally intensive for three-dimensional global-scale modelling that compute whole magnetospheres. Meanwhile, the inclusion of interior effects allows consideration of magnetic diffusion through the moon's surface, magnetic coupling with the ionosphere¹⁰, and the induced magnetic field from the subsurface ocean (Kivelson et al., 2002; Jia et al., 2008; Jia et al., 2009; Jia, Kivelson, et al., 2010). Outputs from the resistive MHD model suggest frequent magnetic reconnection events on Ganymede's upstream magnetopause, each lasting a short time between 1-10 seconds (Jia, Walker, et al., 2010).

The next major advance in Ganymede's numerical modelling is the Hall-MHD model (Tóth et al., 2008; Dorelli et al., 2015). The Hall modifications allow more direct capture of local plasma demagnetisation, reducing reliance on model-produced numerical resistivity. The Hall-MHD model predicts asymmetric Jovian plasma deflection around Ganymede, leading to higher plasma flow shear on the sub-Jovian magnetopause flank, and thus cross-boundary plasma injection through K-H instability (Dorelli et al., 2015). Magnetic reconnection on the upstream magnetopause is also shown to circulate ions throughout Ganymede's magnetosphere, consistent with expectation from a Dungey convection cycle. The Hall-MHD has since been further modified to capture more kinetic plasma effects near Ganymede (Fatemi et al., 2016; Poppe et al., 2018).

An alternative attempt to address kinetic effects is to embed the Hall-MHD model with particle-in-cell (PIC) codes, which directly solve the collisionless Vlasov equation to capture plasma demagnetisation and diffusion processes (Daldorff et al., 2014; Tóth et al., 2016; Zhou et al., 2019). The MHD-EPIC code solves Hall-MHD equations at global scale, while restricting its kinetic component to small parts of the domain i.e., the magnetopause, in order to keep computational cost reasonable. Results between the two models are similar, albeit with some quantitative differences in plasma properties in the magnetosphere (Tóth et al., 2016). The

¹⁰ A conducting, ionised layer at the top of Ganymede's atmosphere

MHD-EPIC model also shows magnetic reconnection signals consistent with measurements from a Galileo flyby. Further Ganymede kinetic model development besides MHD-EPIC include high-moment multi-fluid codes, in which the Vlasov equation is integrated in velocity space to obtain moment equations of motion for multiple particle species (Wang et al., 2018).

Global numerical simulations of Ganymede typically take temporally constant, spatially uniform inputs of upstream Jovian plasma-magnetic parameters. The models are then left to run over a period of time, during which the magnetosphere settles into approximately steady-state equilibrium with fluctuations that potentially denote interactions. Hence, numerical simulations do not account for real-time variations in the upstream Jovian plasma sheet. Fortunately, this shortcoming often does not impact general conclusions drawn from the models, as upstream conditions are relatively steady at Ganymede's length and time scales (Jia, Walker, et al., 2010).

Outside of simulations, our knowledge of Ganymede's magnetosphere is supplemented by remote observations. The 1990 launch of Hubble Space Telescope (HST) represents a major advancement in our ability to image distant objects. Unhindered by atmospheric attenuation, HST delivered the first view of Ganymede's aurora (Hall et al., 1998), illustrating evidence of magnetospheric plasma convection. The kinetic physics of auroral excitation at Ganymede is not fully understood, but it is likely related to plasma acceleration by its random collective behaviour or surrounding current systems (Saur et al., 2015). The moon's emissions form polar rings with significant upstream-downstream asymmetry ($>20^\circ$ difference in co-latitude), which is a unique feature due to sub-Alfvénic magnetospheric interactions (Saur et al., 2015; Musacchio et al., 2017).

A few examples of Ganymede's auroral rings are shown in Figure 17. Composite images taken over multiple days show the rings to be fairly stable in location, although some movements are possible at shorter time scales (McGrath et al., 2013). One well-studied ring motion is its rocking oscillation under periodic variation of the Ganymede magnetic field. HST measurements show that the rocking agrees with magnetospheric field variation if the field contains a secondary ocean-induced dipole, rather than higher-degree contributions to the primary dynamo-generated dipole (Saur et al., 2015). Aurora rings may also shift latitudinally due to their positional correlation to the open-closed field line boundary (OCFB). This point will be further discussed in Chapter 4, when I consider dynamics inside Ganymede's magnetosphere in more detail.

Ganymede's magnetosphere can also be studied indirectly through its effects on Jupiter's magnetosphere. The two-body connection, facilitated by open magnetic field lines, means particle streams from Ganymede induce footprint emission in Jupiter's polar regions, which is then observed by visiting spacecraft and radio telescopes on Earth (e.g., Louis et al., 2017; Zarka et al., 2018; Louis et al., 2020). Time distribution of Ganymede-Jupiter footprint emissions suggest sporadic particle/energy releases by magnetic reconnection at Ganymede's upstream magnetopause, that is modulated in some ways by Jupiter's rotation (Zarka et al., 2018). Otherwise, HST measurements have observed motions in Ganymede-Jupiter footprint emissions, which can indicate changes to Ganymede's magnetospheric field structure (Promfu et al., 2020).

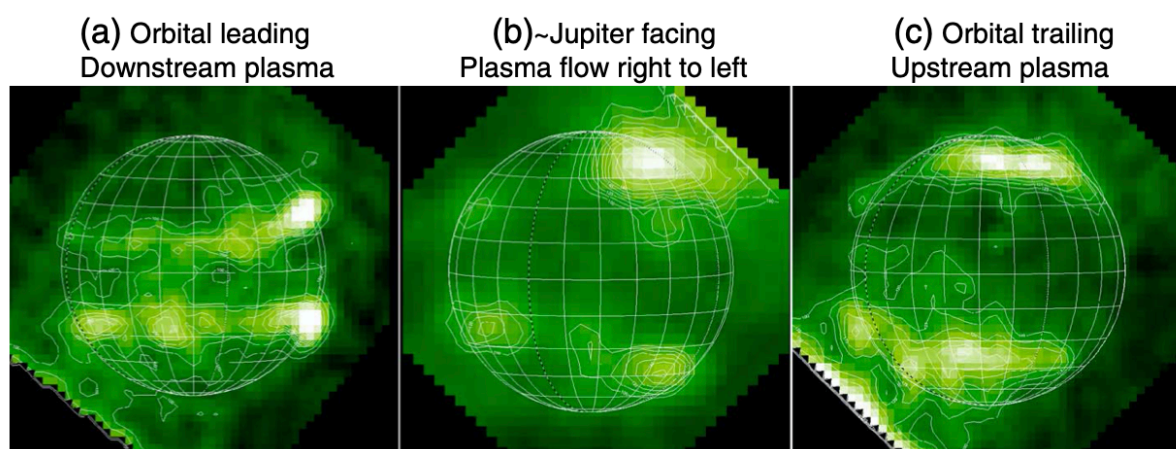


Figure 17: Images of Ganymede's aurora taken by the Hubble Space Telescope from (a) downstream, (b) sub-Jovian terminator, and (c) upstream of the moon. Image courtesy of McGrath et al., (2013), *Journal of Geophysical Research: Space Physics*.

Our knowledge of the Ganymedean magnetosphere is in its infancy and will grow with added observations. As of this writing, Juno's extended mission has conducted at least one flyby of Ganymede on 7 June 2021 (Buccino et al., 2022; Kurth et al., 2022; Paranicas et al., 2022). Given Juno's polar orbits, the spinning spacecraft, and its extensive instrumentation, the flyby's still-to-be-published data should add insights into plasma-magnetic conditions in the high-latitude magnetosphere, including the Alfvén wing. After Juno, the next major spacecraft to visit Ganymede will be the JUpiter ICy moon Explorer (JUICE) (Grasset et al., 2013). The ESA mission is currently slated to launch in 2023 and arrive at Jupiter around 2030. After a few years of Jovian orbits, the spacecraft will be inserted into orbit around Ganymede, becoming the first spacecraft to orbit a planetary satellite beyond our own Moon. JUICE will extensively study Ganymede's magnetic environment and allow unprecedented observations of the subsurface ocean, significantly improving our understanding of the Solar System's largest and most planet-like moon (Grasset et al., 2013). The scientific objectives of JUICE, and the spacecraft's specific contributions to Ganymedean research, will be discussed in Chapter 7 in context of relevant research presented in this thesis.

In the meantime, however, much of Ganymede's magnetosphere remains poorly understood. On the upstream magnetopause, Galileo findings indicate some evidence for magnetic reconnection as shown in Figure 18 (Kivelson et al., 1998; Jia, Walker, et al., 2010, Tóth et al., 2016). The spacecraft data also presents potential evidence for K-H instability although this is disputed (Kivelson et al., 1998; Volwerk et al., 2013). Without a larger observation sample, however, conclusions drawn on these interactions remain reliant on physically limited numerical simulations. Neither energy transport mechanism has been subjected to direct study. Consequently, their fundamental natures including global structures (i.e., extended or patchy), properties (i.e., steady or transient), and influence on magnetospheric plasma convection remain largely unknown.

When Galileo crossed into Ganymede's magnetosphere, the spacecraft detected plasma-magnetic features consistent with expectations for slow-rotating environment, including low-

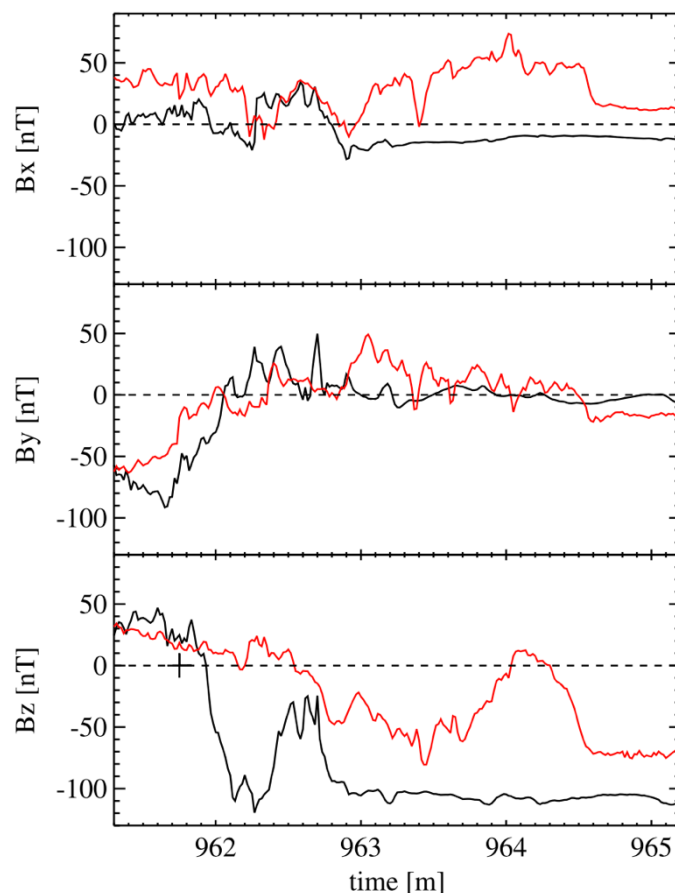


Figure 18: Comparison between magnetometer data from Galileo's G8 flyby (black) to MHD-EPIC simulation outputs (red). (B_x , B_y , B_z) are in Ganymede-centred Cartesian coordinates later explained in Chapter 3, Section 2. Qualitative similarities are observed between the two fluctuating plotted lines, which are interpreted as a likely in-situ observation of magnetic reconnection on Ganymede's upstream magnetopause. Image courtesy of Tóth et al., (2016), *Journal of Geophysical Research: Space Physics*.

frequency oscillations/resonances of magnetic field lines (Volwerk et al., 1999; Volwerk et al., 2013), and particle-trapping radiation belts (Williams, Mauk, McEntire, 1997; Williams et al., 1998; Williams, 2001, 2004). But similar to magnetopause processes, any understanding of these features beyond their existence is limited by lack of observational data available.

At Ganymede's polar caps, HST observations were able to capture brightness, asymmetry, and some motion of the auroral rings (McGrath et al., 2013; Saur et al., 2015; Musacchio et al., 2017). However, rings mapping attempts are incomplete as HST cannot simultaneously observe their entire structure from its near-Earth vantage point. Auroral morphology have been studied in relation to Ganymede's position in the Jovian plasma sheet (Saur et al., 2015; Musacchio et al., 2017). However, remote measurements are limited in their ability to resolve smaller structures, particularly transient emission spikes that associate with rapid plasma-magnetic interactions, that would be possible with an in-situ observer.

Ganymede is understood to have an Alfvén wing magnetotail structure (Neubauer, 1998). Numerical simulations suggest presence of a downstream magnetic reconnection site (e.g., Jia et al., 2008), and Galileo has observed evidence of moonward plasma flow (Collinson et al., 2018). However, the reconnection site has never been characterised, nor has the downstream equatorial plasma sheet. Although a Dungey-like plasma convection cycle for Ganymede is likely, its adaptations for sub-Alfvénic environment remain unclear. Magnetotail plasma losses are expected along open magnetic field lines at high Alfvén wing latitudes (Frank et al., 1997; Williams, Mauk, McEntire, 1997, Williams, Mauk, McEntire, Roelof, et al., 1997), but the extent of contribution from downstream reconnection in equatorial regions is again inconclusive.

Back at Ganymede, the moon's internal magnetic field structure also contains uncertainties. Limited data from Galileo means higher-degree contributions to the primary dynamo field cannot be resolved. The Ganymedean induction process (Kivelson et al., 2002), thought to produce a secondary magnetic field, is also poorly understood. Properties of the subsurface ocean, and the ocean's convective interactions with surrounding magnetic fields, remain severely unconstrained.

Many of these open questions cannot be conclusively answered before the arrival of JUICE in 2030. However, this does not mean that Ganymedean research is halted in the intervening years. Any information obtained about the moon will not only lay groundwork for future research, but also inform the ongoing planning for the JUICE mission. Helpful predictions can guide spacecraft trajectory, instrument calibration, and observation schedule. So, what can we still learn about this magnetic and watery moon? In the last three years, I have sought to provide answers for this question by assessing aspects of Ganymede's magnetic domain at global scale, utilising fundamental magnetic and plasma theories to achieve concrete understanding of the Solar System's most unique magnetosphere. My body of research is divided into four work chapters that will follow from this introduction.

The first work chapter concerns magnetic reconnection on Ganymede's upstream magnetopause and its role in cross-boundary energy transport. Reconnection onset conditions will be assessed and its associated electric field, proportional to its rate, will be calculated. The result is then analysed to determine where reconnection can occur on the magnetopause, and whether the interaction can be characterised by any pattern.

The second work chapter further explores the implications of magnetic reconnection on global scale plasma convection in Ganymede's magnetosphere. Are reconnection events generally localised and transient as predicted by numerical MHD simulations (Jia, Walker, et al., 2010), or do they occur in more stable extended structures? The effect of reconnection on the OCFB, and potentially aurora motion, will also be considered.

The third work chapter is similar to the first, but instead concerns K-H instability growth on Ganymede's upstream magnetopause. The onset condition for linear K-H waves will be assessed to determine their prevalence at the boundary. Then, the instability growth rate will be characterised, taking into account kinetic effects arising from heavy Jovian plasma ions, to determine viability of nonlinear K-H vortices near Ganymede.

The fourth and final work chapter considers the effects of magnetopause interactions on magnetic induction in Ganymede's subsurface ocean. Periodic changes in near-boundary conditions result in oscillation of the Chapman-Ferraro magnetic field, which then induce a

secondary magnetic field from the conducting saltwater. This induction work will focus on quantification of the inducing Chapman-Ferraro field, rather than full computation of the inductive response.

In the final conclusion chapter, I will reassess the impacts of my findings on both the ongoing knowledge of Ganymede's magnetosphere, and the planning of the JUICE mission. The latter aspect will focus on the data required to verify my conclusions, and how it can be provided by the spacecraft's instruments. With this section, the introduction to this thesis is now complete, and our investigation of Ganymede's magnetosphere can now begin.

Chapter 3: Magnetic Reconnection on Ganymede's Upstream Magnetopause

We begin this chapter with an overview of magnetic reconnection in the magnetopause environment, with emphasis on the interaction's three onset conditions and rate – the latter represented by an associated electric field. Existing literature on reconnection assessments at Solar System planets and Ganymede is then briefly reviewed. Benefits and downsides are compared among assessments based on in-situ observations, numerical simulations, and analytical modelling. Despite its estimative nature, analytical modelling is well-suited to efficiently unveil the fundamental global nature of Ganymede's upstream magnetopause reconnection, which has never been assessed before.

The chapter proceeds to introduce idealised mathematical description of Ganymede's magnetopause (Kivelson et al., 1996), followed by parametrisations of plasma and magnetic properties on both sides of the boundary. Subsequent evaluation shows all three reconnection onset conditions to be well satisfied across the entire current sheet surface throughout Ganymede's orbit of Jupiter. The reconnection electric field is also shown to be large everywhere, allowing widespread reconnection without constraint on its global structure. The average field magnitude, however, expresses dependence on Ganymede's position in the Jovian plasma sheet, and is therefore effectively controlled by Jupiter's rotation.

3.1. Reconnection Onset and Rate

Plasma convection around the slow-rotating Ganymede likely follows a Dungey-like pattern (Dungey, 1961), driven by energetic plasma injection across the upstream magnetopause that divides Jupiter-Ganymede plasma-magnetic regimes. One of the two mechanisms for cross-magnetopause particle/energy transport is magnetic reconnection, whose basic concept has been introduced in Chapter 1, Section 1.4. This chapter will provide a deeper look into reconnection and its existing research, followed by the first analytical assessment of the interaction on Ganymede's upstream magnetopause.

Figure 19a considers a simple plasma-magnetic configuration at the equatorial magnetopause in the body-centred Cartesian coordinates where X points along upstream plasma inflow direction, Z points northward along the body's rotation axis, and Y completes right-hand orthogonality along the body's orbital ecliptic plane. The upstream plasma possesses velocity $\mathbf{U}_0$ and mass density ρ_0 , carrying a magnetic field $\mathbf{B}_0$ whose orientation is strongly (although not perfectly) orthogonal to the flow direction. Ignoring the small flow-parallel component, we can define the upstream magnetic field's clock angle

$$\tan(\theta_{cl}) = \frac{B_{y,0}}{B_{z,0}} \quad (3.1)$$

Clock angles between $0^\circ - 90^\circ$ correspond to $B_{z,0} > 0$, so the upstream magnetic field is said to be northward. Thus, clock angles between $90^\circ - 180^\circ$ correspond to $B_{z,0} < 0$ and southward upstream field. As inflow plasma slows near the magnetopause, the upstream field is slightly compressed, leading to an external (to the magnetopause) field of new magnitude $\mathbf{B}_0 \rightarrow \mathbf{B}_{ex}$ and minimally changed direction. External plasma meets the magnetopause with local surface normal vector $\mathbf{N} = (N_x, N_y, N_z)^{11}$, resulting in a deflection defined by the flaring angle

¹¹ Pointing toward the magnetospheric side

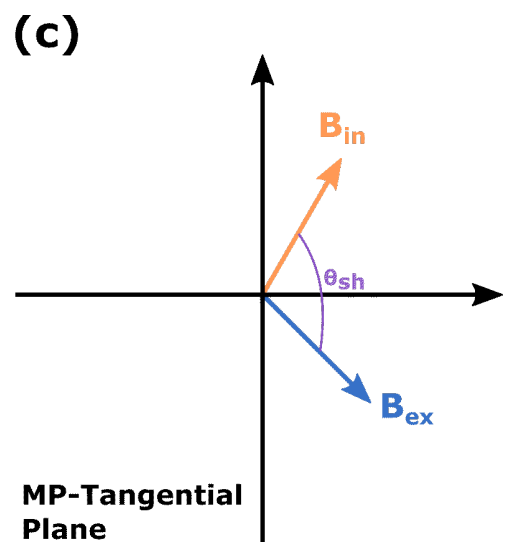
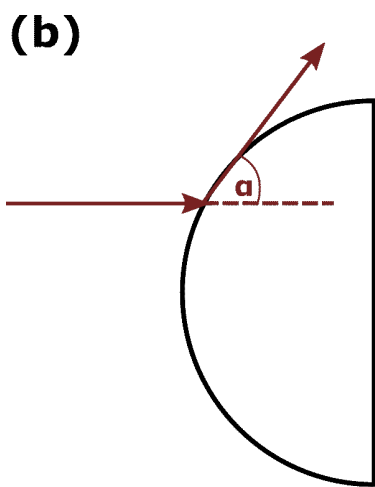
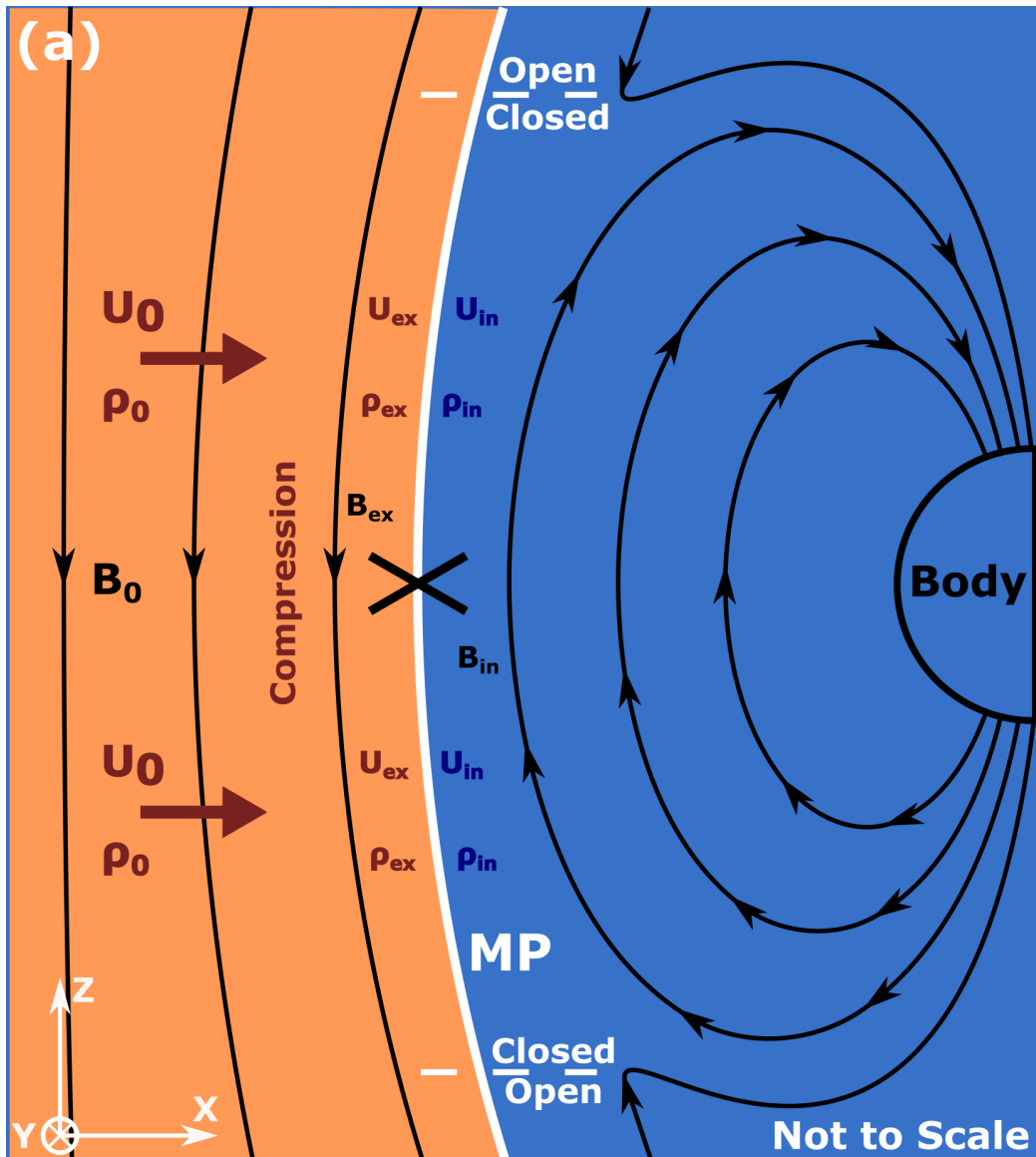


Figure 19: Schematic diagrams for (a) magnetic configuration near an upstream magnetopause, (b) flaring angle, and (c) magnetic shear angle.

$$\cos(\alpha) = \frac{\mathbf{N} \cdot \mathbf{U}_{\text{ex}}}{|\mathbf{N}| |\mathbf{U}_{\text{ex}}|} \quad (3.2)$$

where $\mathbf{U}_{\text{ex}}$ is the external plasma velocity ($\mathbf{U}_{\text{ex}} \approx \mathbf{U}_0$). The inflow plasma is generally assumed to be horizontal along the X-axis, which simplifies Eq. 3.2 and leads to a schematic diagram in Figure 19b. Just inside the magnetopause, the body's closed magnetospheric field $\mathbf{B}_{\text{in}}$ field lines push back against external pressure. Magnetic fields adjacent to the magnetopause can be projected onto a local tangential plane, shown in Figure 19c, on which we define the magnetic shear angle

$$\cos(\theta_{\text{sh}}) = \frac{\mathbf{B}_{\text{ex}} \cdot \mathbf{B}_{\text{in}}}{|\mathbf{B}_{\text{ex}}| |\mathbf{B}_{\text{in}}|} \quad (3.3)$$

where $\mathbf{B}_{\text{ex}}$ and $\mathbf{B}_{\text{in}}$ are adjacent magnetic field vectors. Magnetic fields of shear angle $\theta_{\text{sh}} = 180^\circ$ is said to be perfectly antiparallel. If the shear angle is smaller, then there exists a magnetic guide field component orthogonal to the main antiparallel direction, which exerts influence on local motions and thus reconnection physics.

The site of magnetic reconnection is denoted in Figure 19a with the symbol 'X'. As discussed in Chapter 1, Section 1.4, reconnection is often characterised by ion/electron inertial lengths (d_i , d_e), which define the scales below which particles decouple from their magnetic fields and diffuse across the magnetopause. The diffusion region is typically defined as a three-dimensional box of cross section $\sim d_{i,e} \times 10d_{i,e}$ (Fuselier & Lewis, 2011). Due to the ions' larger masses, they have larger inertial lengths than electrons (Eq. 1.14 – 1.15) and thus larger diffusion regions. Figure 20 shows an illustration of ion and electron diffusion regions in the Hall fast reconnection model. When a pair of adjacent external/internal magnetic field lines reconnect, they form two open field lines that travel away from the reconnection site, dragging away diffused plasmas in the form of bidirectional outflow jets.

At any point on the magnetopause, magnetic reconnection can occur if three specific onset conditions are satisfied. First, the field-separating current sheet must be thinner than one ion

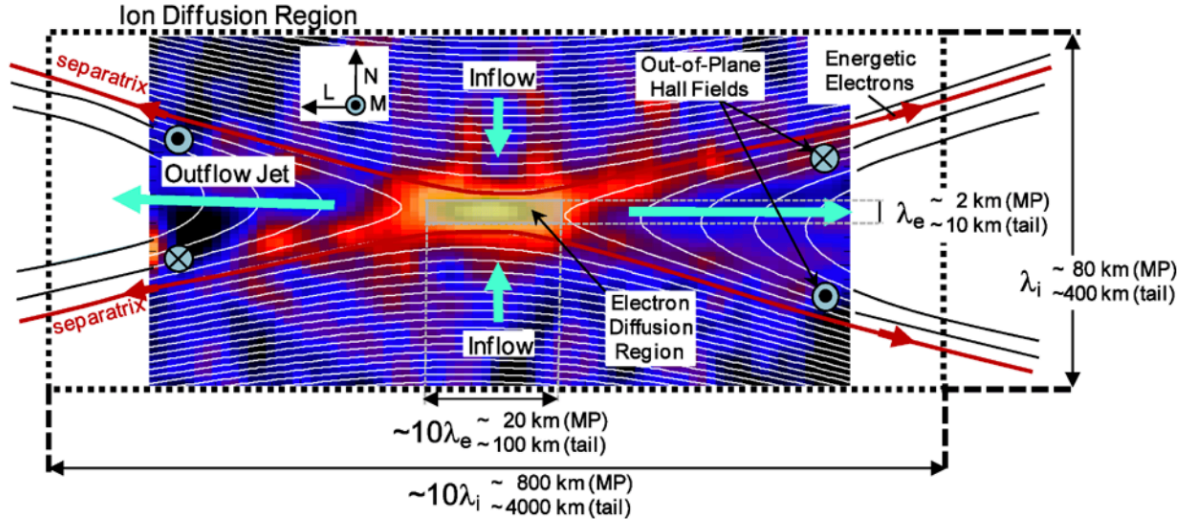


Figure 20: A magnetic reconnection site in Earth's plasma-magnetic environment. Diffusion regions are located at central conjunction of magnetic field lines. Image courtesy of Fusilier & Lewis, (2011), *Space Science Review*.

inertial length (d_i) in order to violate MHD's locality assumption, which allows plasma detachment from magnetic field and thus diffusion through the regions discussed above (Phan et al., 2011). Second, the shear angle between adjacent magnetic fields must be larger than a threshold (Swisdak et al., 2003; Swisdak et al., 2010)

$$\theta_{sh} > 2 \tan^{-1} \left(\frac{d_i \Delta\beta}{L} \right) \approx 2 \tan^{-1} (\Delta\beta) \quad (3.4)$$

where the system length L is assumed equal to the ion inertial length d_i , and $\Delta\beta = \beta_{ex} - \beta_{in}$ is the plasma beta difference across the magnetopause. Eq. 3.4 is known as the diamagnetic drift condition and measures the reconnection-suppressive effect of the magnetic guide field component parallel to diffusion region (i.e., in/out of page in Figure 20). If the guide field is sufficiently strong i.e., small θ_{sh} , local ions will drift away from the reconnection site at net velocity greater than that of the outflow jets. When this occurs, local centrifugal force overcomes magnetic tension of newly reconnected field lines, which stops plasma outflow and suppresses reconnection. (Swisdak et al., 2010). Diamagnetic drift is a kinetic phenomenon, therefore Eq. 3.4 is not reliant on MHD description of plasma.

Third, the velocity shear of magnetopause-adjacent bulk plasmas, in direction of the reconnection outflow, cannot exceed an upper limit (Doss et al., 2015). To express this condition, we define the outflow directions as unit vectors that are parallel/antiparallel to the

cross product between the vector bisecting the adjacent magnetic fields and the local magnetopause normal i.e.,

$$\mathbf{v}_{bi} = \frac{\hat{\mathbf{B}}_{ex} + \hat{\mathbf{B}}_{in}}{2} \quad (3.5)$$

$$\hat{\mathbf{v}}_{out} = \pm \frac{\mathbf{v}_{bi} \times \hat{\mathbf{N}}}{|\mathbf{v}_{bi} \times \hat{\mathbf{N}}|} \quad (3.6)$$

where the normal $\mathbf{N}$ points toward Ganymede, and the hat subscripts denote unit vectors. We define the primary outflow vector to be parallel to the external magnetic field direction. In case of Ganymede's magnetopause, this means the primary outflow points southward with the Jovian magnetic field i.e., $\hat{v}_{out,z} < 0$. Local plasma mass densities, magnetic fields, and bulk flow velocities along the primary outflow is then defined by subscripts '1' and '2' where

$$\rho_1 = \rho_{ex} \quad (3.7)$$

$$\rho_2 = \rho_{in} \quad (3.8)$$

$$\mathbf{B}_1 = \mathbf{B}_{ex} \cdot \hat{\mathbf{v}}_{out} \quad (3.9)$$

$$\mathbf{B}_2 = \mathbf{B}_{in} \cdot \hat{\mathbf{v}}_{out} \quad (3.10)$$

$$\mathbf{v}_1 = \mathbf{v}_{ex} \cdot \hat{\mathbf{v}}_{out} \quad (3.11)$$

$$\mathbf{v}_2 = \mathbf{v}_{in} \cdot \hat{\mathbf{v}}_{out} \quad (3.12)$$

The bulk flow shear condition hence follows

$$v_{1-2} = \frac{|v_1 - v_2|}{2} < v_{out} \left(\frac{\rho_1 B_1 + \rho_2 B_2}{2\sqrt{\rho_1 B_1 \rho_2 B_2}} \right) \quad (3.13)$$

$$v_{out} = \sqrt{\frac{B_1 B_2 (B_1 + B_2)}{\mu_0 (\rho_1 B_2 + \rho_2 B_1)}} \quad (3.14)$$

where $\mu_0 = 4\pi \times 10^{-7}$ H/m is the permeability of free space and v_{out} is magnitude of the outflow velocity vector. As v_{1-2} approaches its limit, reconnection outflow jets become overrun by local background flows, preventing reconnection from occurring. The outflow speed v_{out} can be replaced typically by the local Alfvén speed C_A (Eq. 1.49) (Doss et al., 2015).

Furthermore, in special case of symmetric reconnection, the bracketed fraction in Eq. 3.13 can be reduced to one, making the Alfvén speed an estimated upper limit for flow-based reconnection suppression.

Plasma inflow into a magnetic reconnection site generates an electric field that is parallel to the magnetopause current sheet. This is known as the reconnection electric field. Since this field measures plasma/magnetic fluxes entering the reconnection site, its magnitude is often defined as the unnormalised reconnection rate. The electric field can be computed via the formula (Doss et al., 2015)

$$E_{MR} = 2kv_{out} \left(\frac{B_1 B_2}{B_1 + B_2} \right) \left(1 - \left(\frac{v_1 - v_2}{v_{out}} \right)^2 \frac{\rho_1 B_2 \rho_2 B_1}{(\rho_1 B_2 + \rho_2 B_1)^2} \right) \quad (3.15)$$

where k is the reconnection efficiency (i.e., the ratio between magnetopause-normal magnetic field and the total internal magnetic field) between 0 and 1. Note that E_{MR} is only one component of the larger convective electric field, which contains additional contributions from plasma flowing parallel to the magnetopause (but not into the reconnection site). Eq. 3.15 includes the effect of velocity shear onset condition. As the shear $v_1 - v_2$ approaches the outflow velocity v_{out} , which is an approximate upper limit for reconnection suppression, the electric field rapidly drops toward zero. Eq. 3.15 is built on the MHD plasma fluid description and does not account for diamagnetic drift effect. Hence, the reconnection electric field should be computed only after the diamagnetic drift onset condition has been assessed.

Across a magnetopause surface, locations that satisfy all three onset conditions are potential magnetic reconnection sites, whose associated electric field can then be calculated. This information can provide some imagery of reconnection structure on the boundary. For instance, consider a magnetopause containing a large patch of viable reconnection sites, but whose electric fields strongly peak along a much narrower line. In this case, reconnection may form an elongated ‘X-line’ structure that stretches across the magnetopause¹². The X-line has been observed across the Earth’s magnetopause and is known to be steady up to multiple

¹² Hence, the ‘X’ symbols are used to signify magnetic reconnection in schematic diagrams.

hours (e.g., Phan et al., 2000; Dunlop et al., 2011). If the reconnection electric field is assumed to be parallel to the X-line, the reconnection voltage (i.e., electric potential produced by reconnection) can be calculated via the line integral

$$V_{\text{MR}} = - \oint \mathbf{E} \cdot d\mathbf{l} \sim -EL \quad (3.16)$$

which can be approximated as a simple product where L is length of the X-line. A pattern of reconnection sites and rates, however, does not guarantee a specific reconnection structure. In addition to X-lines, reconnection can occur as transient flux-transfer events (FTEs) within timescales of minutes and seconds. The relation between steady and transient reconnection at global scale will be more explicitly discussed in Chapter 4.

3.2. Analytical Modelling of Ganymede's Magnetopause

Assessing onset and rate of magnetic reconnection at global scale requires extensive temporal mapping of the magnetopause surface. Where in-situ observations are sufficient in number, reconnection onset and rate have been studied at the magnetopauses of Earth (Phan, Paschmann, et al., 2013), Mercury (DiBraccio et al., 2013), and Saturn (Masters et al., 2012; Sawyer et al., 2019). While these analyses offered important characterisations of individual reconnection events, they do not yield global pictures of reconnection structure, or assess how the pictures vary over time.

Meanwhile, numerical simulations are capable of capturing entire magnetopause surfaces. Global MHD codes have been used to study large-scale reconnection at the magnetopauses of Earth (e.g., Ogino et al., 1986; Park et al., 2006; Eggington et al., 2020), Mercury (Jia et al., 2019), Saturn (Fukazawa et al., 2007) and Ganymede (Jia et al., 2008; Jia et al., 2009). In these studies, reconnection events are detected via tell-tale plasma-magnetic signals, specifically evidence of rapid outflow jets, magnetic null points, and magnetic field line tracing. Reconnection onset conditions and electric field are typically not evaluated directly, especially as MHD simulations are not equipped to account for reconnection microphysics (e.g., diamagnetic drifts). Unfortunately, global kinetic codes remain computationally

expensive and thus generally unsuitable for magnetospheric systems larger than Mercury's and Ganymede's.

We need a method to capture global magnetopause conditions, then efficiently evaluate reconnection onset conditions and subsequently its electric field at every point on the boundary. This can be achieved by analytical modelling, which utilises mathematical parametrisations to create a magnetopause surface, and then approximate adjacent plasma-magnetic condition parameters (bulk flow velocities, mass densities, magnetic field vectors, etc). As an analytical model only once converts inputs to outputs, its codes run very efficiently (often within seconds), and input adjustments can be made with ease. However, simple parametrisations do not reiterate feedback effects, so the model does not capture full dynamics around the magnetopause. Its output values are only estimates and should not be taken precisely. Nevertheless, an analytical model can provide useful general conclusions and support for results from numerical simulations.

Analytical reconnection assessment has its origin at Earth's magnetopause (Alexeev et al., 1998; Trattner et al., 2007a; Trattner et al., 2007b). Trattner et al., (2007a, 2007b) in particular evaluate Polar spacecraft observations against analytically computed magnetic shear angles, showing reconnection X-lines to favour sites where the diamagnetic drift condition is optimised (i.e., maximum shear angle). However, it is at outer planets where analytical model usage is the most prominent. Reconnection onset and rate have been assessed analytically at the magnetopauses of Jupiter (Desroche et al., 2012; Masters, 2017), Saturn (Desroche et al., 2013; Masters, 2015b), Uranus (Masters, 2014), and Neptune (Masters, 2015a). Along these boundaries, locations with three satisfied onset conditions are considered to be viable for reconnection. These regional patches partially cover the magnetopause and can be equatorial or polar depending on the external interplanetary magnetic field (IMF) orientation. For instance, Jupiter's dayside magnetopause reconnection is equatorial when the IMF points strongly northward or southward from the solar ecliptic plane (i.e., almost perfectly parallel/antiparallel to the magnetospheric field). However, the events become mainly polar when the IMF lies in the ecliptic plane (Masters, 2017). These results are shown to be consistent with global MHD simulations of Jupiter (Sarkango et al., 2017).

For the rest of this chapter, we return to Ganymede. Subsequent construction of the moon's analytical magnetopause model and assessment of magnetic reconnection are produced through MATLAB. Relevant mathematical equations are implemented as functions, and then run together through streamlined scripts that can be deployed repeatedly using prescribed upstream Jovian plasma-magnetic condition inputs. In this thesis, description of modelling will focus on its scientific achievements, with coding minutiae relegated to footnotes where necessary.

As mentioned in Chapter 2, Section 2.3, Ganymede's magnetosphere is first described as an ovoid column with magnetopause location at (Kivelson et al., 1998)

$$\left(\frac{X' - X_0}{a}\right)^2 + \left(\frac{Y' - Y_0}{b}\right)^2 = 1 \quad (3.17)$$

where the offsets (X_0 , Y_0) are functions of vertical distance from Ganymede's equator

$$X_0(Z') = X_0(0) + |Z'| \tan(\theta_A) \quad (3.18)$$

$$Y_0(Z') = \frac{2}{\pi} Y_{0,\max} \sin(\lambda_{\text{III}} - 248^\circ) \tan^{-1}\left(\frac{Z'}{\delta}\right) \quad (3.19)$$

The X_0 offset describes bending of the Alfvén wing. The Y_0 offset stabilises when $Z' \gg \delta$. It is important to note that Eq. 3.17 – 3.19 are written in a Ganymede-centred GphiB coordinate system (X' , Y' , Z'), where Y' and Z' axes are dynamically oriented such that the external Jovian magnetic field lines always lie in the X' - Z' plane. This differs significantly from the fixed Cartesian GphiO coordinates (X , Y , Z), whose Y and Z axes are fixed to Ganymede-Jupiter vector and Jupiter's rotation axis respectively. GphiB coordinates can be rotated into GphiO following

$$\theta_R = \tan^{-1}\left(\frac{|B_{0,z}|}{B_{0,y}}\right) - 90^\circ \quad (3.20)$$

where $B_{0,y}$ and $B_{0,z}$ are (GphiO) Y and Z components of the ambient Jovian magnetic field. Matrix rotation then can be applied to convert (X', Y', Z') into (X, Y, Z) coordinates.

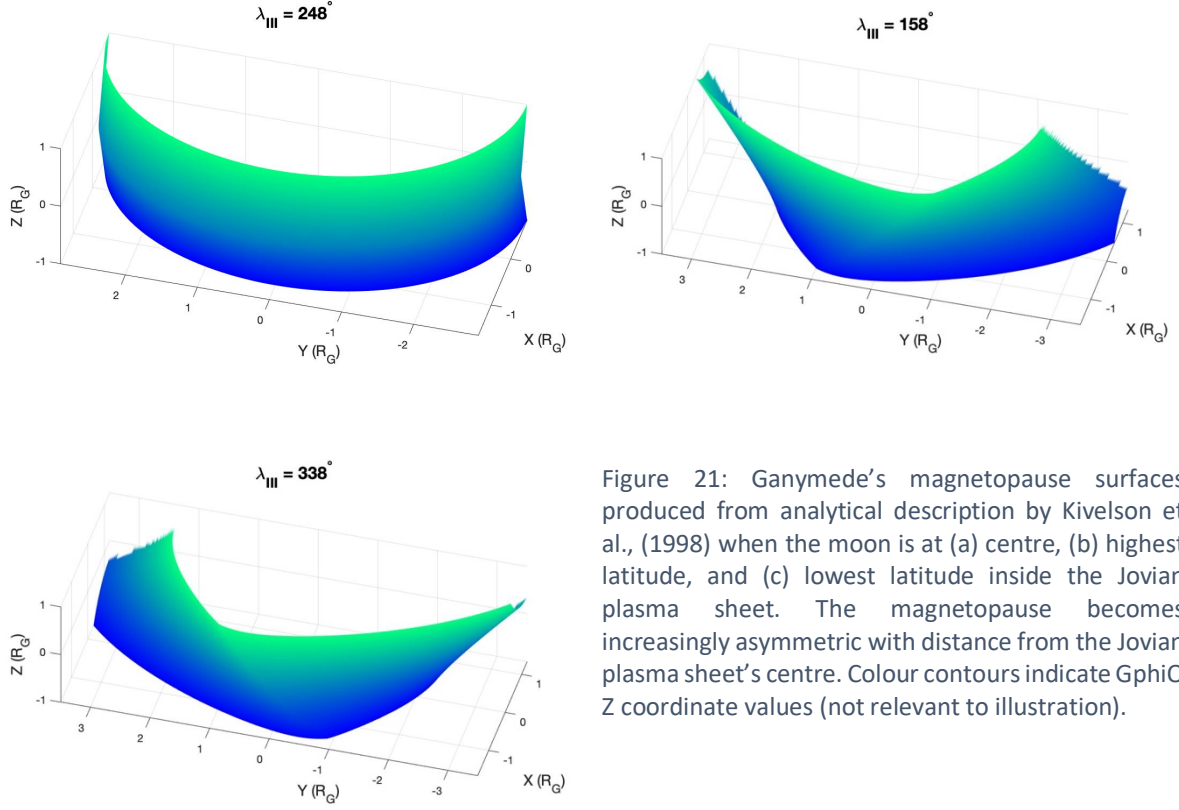


Figure 21: Ganymede's magnetopause surfaces produced from analytical description by Kivelson et al., (1998) when the moon is at (a) centre, (b) highest latitude, and (c) lowest latitude inside the Jovian plasma sheet. The magnetopause becomes increasingly asymmetric with distance from the Jovian plasma sheet's centre. Colour contours indicate GphiO Z coordinate values (not relevant to illustration).

Kivelson et al., (1998) chose $a = 2.2 R_G$ and $\delta = 0.5 R_G$, then derived the other parameters by least-square fits to Galileo observations, yielding $b = 2.90 R_G$, $X_0(0) = 0.544 R_G$, $Y_{0,max} = 0.914 R_G$, and $\theta_A = 0.298$ radians (or 17°). Imperfections of small-sample fits mean derived parameters are only estimates. This leaves only one unconstrained parameter – the System III Jovian longitude λ_{III} , which describes Ganymede's position relative to the Jovian meridian¹³. If Ganymede's orbital period is assumed to be very long compared to Jupiter's rotation period, then the longitude also approximates Ganymede's position relative to centre of the Jovian plasma sheet, and accounts for north-south asymmetry of the upstream magnetopause. In Eq. 3.19, the Y_0 offset is zero when $\lambda_{III} = 248^\circ$, corresponding to Ganymede at the plasma sheet's centre, and the magnetopause is symmetric. Y_0 reaches its two maximums when $\lambda_{III} = 158^\circ$ and $\lambda_{III} = 338^\circ$, corresponding to Ganymede at its highest

¹³ In this work, I use the eastward convention for System III longitude, where longitudes increase anticlockwise around Jupiter. Some other works, including Jia et al., (2008), utilise a westward convention.

and lowest Jovian magnetic latitudes. At these locations, the magnetopause achieves (opposite) maximum north-south asymmetry.

These three longitudes are considered ‘extreme’ cases, and their surfaces are shown in Figure 21. In each subplot, we consider a domain of $-3.5 R_G < Y < 3.5 R_G$ and $-1.0 R_G < Z < 1.0 R_G$ with $0.01 R_G$ resolution. The chosen domain covers the entire closed-field region of Ganymede’s upstream magnetopause, while including some open-field polar areas. Since Eq. 3.9 – 3.11 do not describe elongation of Ganymede’s magnetotail, only upstream coordinate solutions with $X < 0$ are considered. Figure 21 illustrates the cylindrical nature of Ganymede’s magnetosphere, and relatively small Alfvén wing bend back consistent with $\theta_A \sim 25^\circ$ derived from Eq. 2.7. All three boundaries are idealised, and their locations are artificially fixed. For ease of analysis, each magnetopause can be projected onto a two-dimensional Y–Z plane, chosen as the two variables define the domain ranges. An example projection is shown in Figure 22a when Ganymede lies at the Jovian plasma sheet’s centre ($\lambda_{III} = 248^\circ$), with X as the contour parameter. As expected, the magnetopause protrudes furthest upstream at the subflow point ($Y = 0, Z = 0$) with standoff distance $r_{MP} = 1.65 R_G$. This is somewhat smaller than numerical simulation results where $r_{MP} \sim 2 R_G$ (Jia et al., 2008). The X-coordinate values then increase monotonically away from the subflow point in all directions, as the magnetopause curves toward the downstream.

The next step is to parametrise plasma and magnetic parameters adjacent to the magnetopause. Regions immediately external and internal to the boundary are defined to be on the ‘Jovian-side’ and ‘Ganymedean-side’ respectively. We assume the magnetopause is in steady state with no ongoing magnetic reconnection. Upstream of Ganymede, Jovian plasma flows toward the magnetopause at bulk velocity $v_{J,0} \approx 140$ km/s in Ganymede’s rest frame (Williams, Mauk, McEntire, 1997; Williams, Mauk, McEntire, Roelof, et al., 1997, Jia et al., 2008), independent of the moon’s latitude in the Jovian plasma sheet. The inflow plasma is deflected around the magnetopause at speed parametrised by the expression

$$v_J = v_{J,0} \sin(\alpha) \quad (3.21)$$

where α is the flaring angle and v_j is the Jovian-side speed after deflection. The sine relation is a simplification of sinusoidal relations used for bullet-shaped planetary magnetopauses (Petrinec & Russell, 1997), adopted for sub-Alfvénic cylindrical plasma-magnetic configuration. Eq. 3.21 prescribes the Jovian-side plasma to be stagnant at the subflow point ($\alpha = 0^\circ$) due to head-on collision, but almost unperturbed near magnetopause flanks where inflows are almost parallel to the current sheet ($\alpha \rightarrow 90^\circ$). The deflected flow vector components are then computed with two constraints. One, the flow must be parallel to the magnetopause at all times. Two, the flow must remain coplanar to the inflow upstream vector $\mathbf{v}_{J,0}$, which is parallel to the GphiO X-axis. Figure 22b shows deflected plasma flows across the Jovian-side magnetopause when Ganymede lies at the Jovian plasma sheet's centre ($\lambda_{III} = 248^\circ$), with speed and direction denoted by the colour bar and black arrows respectively.

As Ganymede orbits Jupiter, it traverses up-down inside the Jovian plasma sheet, whose density decreases away from the densest centre. Following the convention by Jia et al., (2008), the upstream Jovian plasma mass density is chosen to be $\rho_{J,0} = 56 \text{ amu/cm}^3$ when Ganymede lies at centre of the Jovian plasma sheet, and $\rho_{J,0} = 28 \text{ amu/cm}^3$ when Ganymede is at highest/lowest Jovian magnetic latitudes. Near the magnetopause, the upstream plasma becomes slightly compressed. Unlike bulk flow speed, the density increase is modelled by the cosine relation

$$\rho_J = A \cos(\alpha) + \rho_{J,0} \quad (3.22)$$

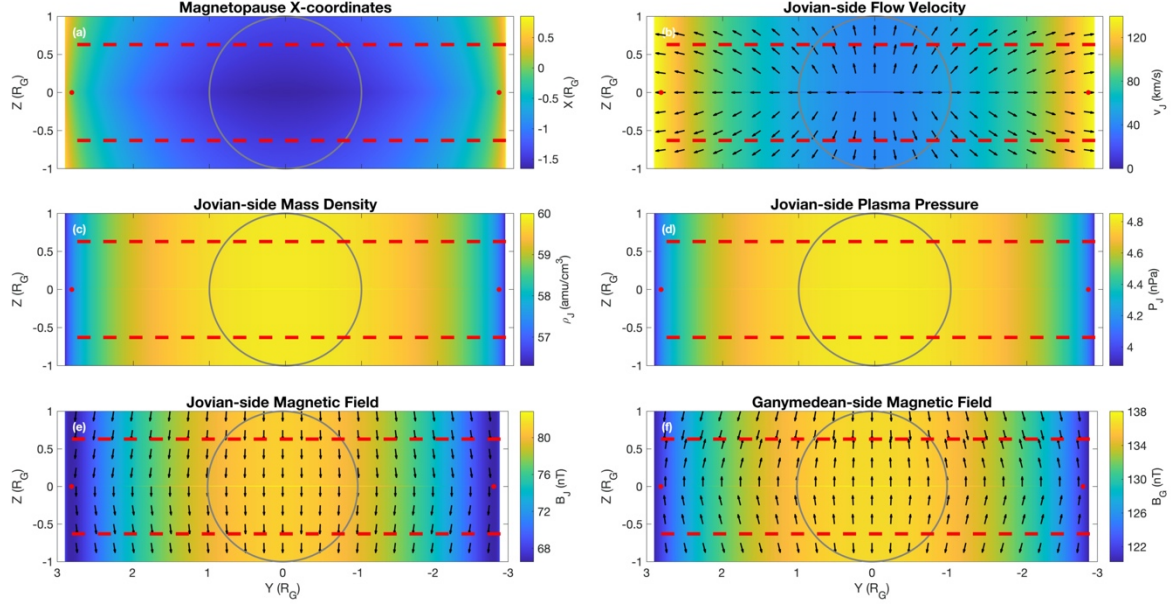


Figure 22: Parametrised plasma and magnetic conditions near Ganymede's magnetopause when the moon lies at centre of the Jovian plasma sheet ($\lambda_{III} = 248^\circ$). Parameters are (a) magnetopause X-coordinates, (b) Jovian-side bulk flow velocity, (c) Jovian-side plasma mass density, (d) Jovian-side plasma pressure, (e) Jovian-side magnetic field, and (f) Ganymede-side magnetic field. Red dashed lines indicate boundaries between Ganymede's open and closed magnetic field regions at $|Z| = 0.63 R_G$. Ganymede is outlined in grey. Image courtesy of Kaweeyanun et al., (2021), *Journal of Geophysical Research, Space Physics*.

where A is the density compression constant, and ρ_J is Jovian-side plasma mass density. Eq. 3.22 yields largest compression at the subflow point ($\alpha = 0^\circ$) where upstream plasma collides head-on with the magnetopause. The compression then decreases toward zero at the magnetopause flanks with boundary-parallel inflow ($\alpha \rightarrow 90^\circ$), and the Jovian-side density approaches its original upstream value. The density compression constant is taken to be $A = 4 \text{ amu/cm}^3$ based on Tóth et al., (2016) numerical results. An example density compression pattern is shown in Figure 22c when Ganymede lies at the Jovian plasma sheet's centre ($\lambda_{III} = 248^\circ$).

The Jovian plasma carries pressure from its thermal and energetic particle populations, whose sum is defined as the plasma pressure. The upstream plasma pressure reaches its maximum at $P_{p,J,0} = 3.8 \text{ nPa}$ when Ganymede lies at centre of the Jovian plasma sheet and decreases to $P_{p,J,0} = 1.9 \text{ nPa}$ when Ganymede is at highest/lowest Jovian magnetic latitudes. Similar to mass density, Jovian-side plasma pressure is compressed by the cosine relation

$$P_{p,J} = B \cos(\alpha) + P_{p,J,0} \quad (3.23)$$

where B is the pressure compression constant, and P_J is the Jovian-side plasma pressure. We take $B = 1.05 \text{ nPa}$ by approximating the super-Alfvénic B function as the Alfvén Mach number approaches one (Petrinec & Russell, 1997). This is an imperfect method, but it produces plasma pressures that are generally consistent with numerical simulations (Tóth et al., 2016). Figure 22d shows plasma pressure compression pattern when Ganymede lies at the Jovian plasma sheet's centre ($\lambda_{III} = 248^\circ$).

The final relevant parameters on the external magnetopause are the Jovian-side magnetic field and its pressure. The Jovian magnetospheric field has been modelled by analytically (Khurana, 1997), and its vector components computed for upstream of Ganymede (Jia et al., 2008). Using results from Jia et al., (2008), we choose upstream field strength $B_{J,0} = 70 \text{ nT}$ when Ganymede lies at the Jovian plasma sheet's centre, and $B_{J,0} = 105 \text{ nT}$ when Ganymede is at highest/lowest Jovian magnetic latitudes. The upstream Jovian magnetic field is assumed to have negligible X component i.e., $B_{J,0,x} = 0 \text{ nT}$. The remaining two components are then parametrised by

$$B_{J,0,y} = -84 \sin(\lambda_{III} - 248^\circ) \quad (3.24)$$

$$B_{J,0,z} = 3 \cos(\lambda_{III} - 248^\circ) - 79 \quad (3.25)$$

where units are in nT. As Ganymede orbits around Jupiter, the upstream Jovian magnetic field oscillates in the Y–Z plane. The Z-component of the field is always negative, and the Y-component produces opposite asymmetry as Ganymede move north and south from the Jovian plasma sheet's centre ($\lambda_{III} = 248^\circ$). At maximum magnetic latitudes $\lambda_{III} = 158^\circ$ and $\lambda_{III} = 338^\circ$, Eq. 3.24 leads to upstream clock angles $\theta_{cl} = 227^\circ$ and $\theta_{cl} = 133^\circ$ respectively. The north-south $B_{J,0}$ asymmetry is consistent with that of the magnetopause, which is crucial as the latter occurs in response to the former. As the clock angle does not deviate beyond 45° from $\theta_{cl} = 180^\circ$, the upstream Jovian magnetic field is always strongly southward.

Like its plasma counterpart, the upstream Jovian magnetic field is compressed just outside the Ganymedean magnetopause. The extent of compression can be calculated by conservation of plasma-magnetic pressure, which requires the sum of plasma, magnetic, and

dynamic ram pressures to remain unchanged before and after magnetopause compression i.e.,

$$P_{p,J,0} + P_{B,J,0} + P_{r,J,0} = P_{p,J} + P_{B,J} + P_{r,J} \quad (3.26)$$

where left- and right-hand sides are the upstream and the Jovian-side conditions respectively. Note that the Jovian-side ram pressure ($P_{r,J}$) only accounts for its magnetopause-parallel component, and the remaining magnetopause-orthogonal component is folded into the Jovian-side magnetic pressure term ($P_{B,J}$). Eq 3.26 can be rearranged for $P_{B,J}$, which is then converted into Jovian-side magnetic field strength (B_J). The Jovian-side field direction is constrained to be parallel to the magnetopause, calculated by projecting the upstream magnetic field vectors onto the boundary. Figure 22e shows the Jovian-side magnetic field vector ($\mathbf{B}_J$) when Ganymede lies at the Jovian plasma sheet's centre ($\lambda_{III} = 248^\circ$). The Jovian-side magnetic field has typical magnitude $B_J \sim 80$ nT, which is consistent with Galileo observations (Kivelson et al., 1998).

Only pressures orthogonal to the magnetopause can impact the boundary location. Thus, we define the Jovian-side acting pressure

$$P_{J,act} = P_{p,J} + P_{B,J} \quad (3.27)$$

This pressure must be balanced on the Ganymedean side to maintain a steady magnetopause. Since Ganymede's magnetospheric plasma is very cold and its pressure negligible (Jia et al., 2008), the Ganymedean-side magnetic pressure can be assumed equal to the external acting pressure i.e.,

$$P_{J,act} \approx P_{B,G} \quad (3.28)$$

and hence the Ganymedean-side magnetic field strength (B_G) can be straightforwardly computed. Just like the Jovian-side field, the Ganymedean-side field direction must be magnetopause-parallel at all locations. Each field vector is also assumed to have no azimuthal component to be consistent with a compressed internal dipole. This implies that

$$\frac{B_{G,y}}{B_{G,x}} = \frac{Y}{X} \quad (3.29)$$

where Y and X are GphiO coordinate values. Finally, the Z -component $B_{G,z}$ is fixed to be southward inside Ganymede's closed-field line region, defined by coordinate range $|Z| < 0.63 R_G$, and northward elsewhere. The Ganymedean-side magnetic field is shown in Figure 22f when the moon lies at the Jovian plasma sheet's centre ($\lambda_{III} = 248^\circ$). Open-closed field boundaries (OCFBs) have also been retroactively added as red dashed lines in all Figure 22 subplots. The typical Ganymedean magnetic field strength $B_G \sim 130$ nT is much stronger than the Jovian-side field. Magnetic reconnection on Ganymede's upstream magnetopause will therefore be asymmetric such that there is net plasma diffusion entry into the moon magnetosphere.

Plasma properties inside Ganymede's magnetosphere are not as well-constrained as the external Jovian plasma, but can be estimated from numerical simulation results. Jia et al., (2008) suggests that Ganymede-side plasma has mass density $\rho_G \approx 32$ amu/cm³ and flows slowly relative to the Jovian-side plasma. For simplicity, I assume Ganymedean-side plasma velocity to be stagnant i.e., $v_G \approx 0$ km/s. Internal mass density and velocity are both considered uniform across the magnetopause, hence they are not included in Figure 22.

Together, subplots in Figure 22 represent maps of relevant plasma and magnetic conditions across Ganymede's upstream magnetopause. The figure can be replicated for every Jovian System-III longitude λ_{III} , so the magnetopause can be modelled throughout Ganymede's orbit, albeit as a sequence of individual time snapshots. Looking forward, we can evaluate magnetic reconnection onset conditions and subsequently its electric field at every grid point in the analytical model domain. This will be the first global assessment of magnetic reconnection that utilises fundamental plasma theory which accounts for kinetic microphysics, to unveil the structure of Ganymedean reconnection which thus far remains poorly understood.

3.3. Magnetic Reconnection Assessment

The first onset condition for magnetic reconnection requires Ganymede's upstream magnetopause to be thinner than one ion inertial length (Phan et al., 2011). Using data from Figure 22, we estimate the Jovian-side ion inertial length (Eq. 1.17) and obtain $d_i = 426$ km. Meanwhile, Galileo measurements estimate a Ganymede magnetopause thickness that is below 400 km (Kivelson et al., 1996). As the result, we assume the first reconnection onset condition is satisfied throughout the upstream magnetopause. In the rest of this section, we will assess the two remaining reconnection onset conditions, compute the reconnection electric field, and derive general conclusions on the nature of Ganymede's upstream reconnection.

In Figure 23, we evaluate the diamagnetic drift onset condition when Ganymede lies at the Jovian plasma sheet's centre ($\lambda_{III} = 248^\circ$). The kinetic condition requires local magnetic shear angle to exceed a minimum threshold defined by the cross-magnetopause plasma beta. For this analysis, only Ganymede's low-latitude closed-field region with antiparallel boundary-adjacent magnetic fields is considered. Figure 23a illustrates the beta difference ($\Delta\beta$) along the magnetopause surface. Given cold plasma on the Ganymede side, we assume $\beta_G = 0$ and therefore $\Delta\beta = \beta_J$. The average beta difference over the magnetopause is $\langle\Delta\beta\rangle = 2.02$, with highest values the magnetopause flanks where the Jovian magnetic field is weakest. A small trough in $\Delta\beta$ value also appears along Ganymede's equator, likely a product of sharp idealised magnetopause surface.

Using plasma beta, the minimum shear threshold (hence defined as $\theta_{sh,min}$) for the diamagnetic drift condition is shown in Figure 23b. The average magnetic shear angle required to satisfy the condition is $\langle\theta_{sh,min}\rangle = 90.3^\circ$. The shear threshold shares the same

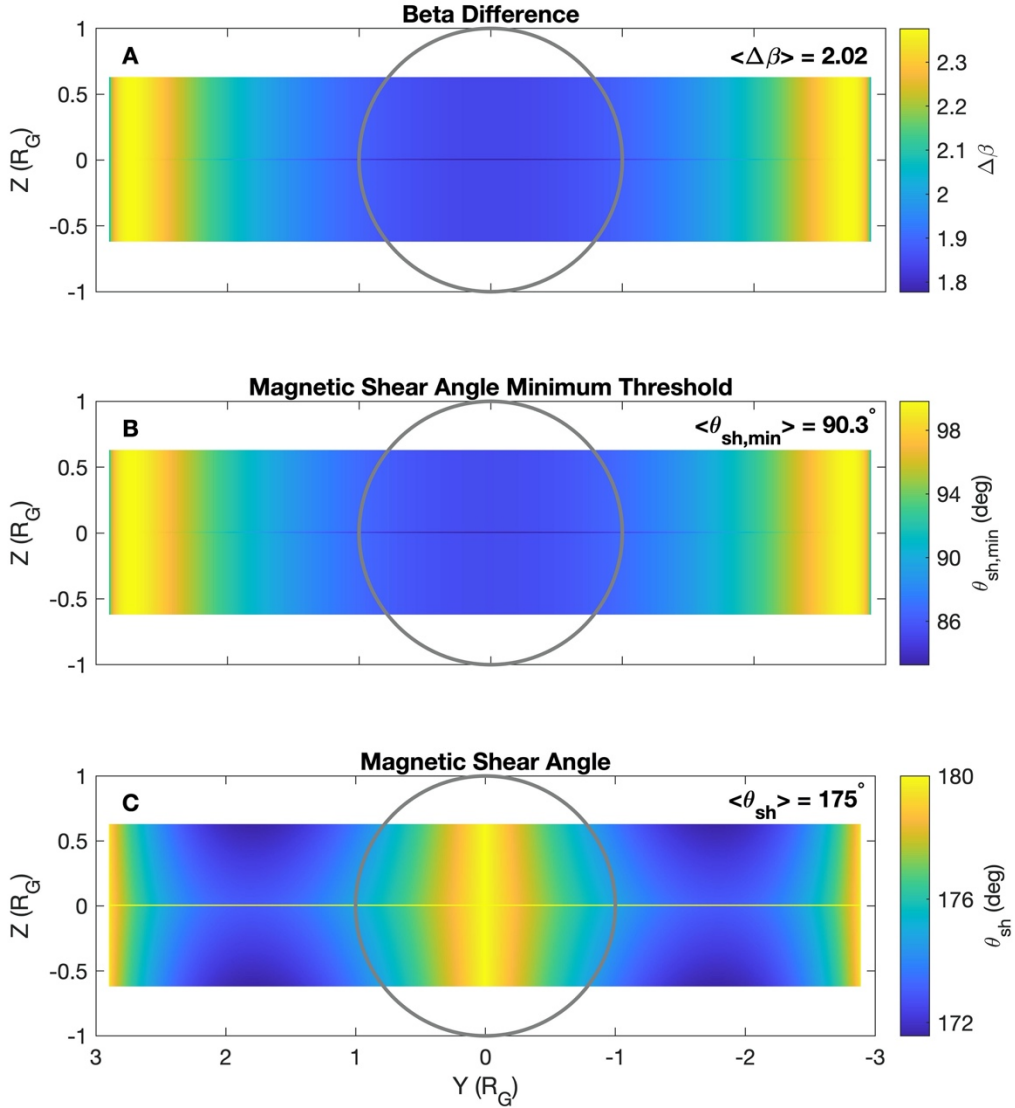


Figure 23: Evaluation of the diamagnetic drift reconnection onset condition when Ganymede lies at centre of the Jovian plasma sheet ($\lambda_{\text{III}} = 248^\circ$). Plots shown are (a) cross-magnetopause difference, (b) minimum magnetic shear angle threshold for reconnection onset, and (c) magnetic shear angle on the magnetopause. The average parameter value in each plot is denoted on the top-right. Ganymede is outlined in grey. Image courtesy of Kaweeyanun et al., (2020). *Geophysical Research Letters*.

variation pattern as the beta difference, which is unsurprisingly given their relationship in Eq. 3.4. The shear threshold range ($\sim 85^\circ - 100^\circ$) also indicates that reconnection is viable with little to no antiparallelism between Jovian and Ganymedean magnetic fields. Finally, Figure 23c shows magnetic shear angles calculated from information in Figure 22e and Figure 22f. The average shear angle is $\langle \theta_{\text{sh}} \rangle = 175^\circ$, and the adjacent magnetic fields are almost completely antiparallel near the subflow point and around the magnetopause flanks. Since the lowest magnetic shear calculated ($\sim 172^\circ$) is much higher than the highest minimum shear threshold, the diamagnetic drift condition is satisfied throughout the magnetopause.

As Ganymede moves toward higher/lower Jovian magnetic latitudes, the average magnetic shear angle decreases with larger magnetopause asymmetry ($\langle\theta_{sh}\rangle = 149^\circ$). However, the much stronger Jovian-side magnetic field significantly reduces the average beta difference ($\langle\Delta\beta\rangle = 0.66$), therefore lowering the average shear threshold ($\langle\theta_{sh,min}\rangle = 36.4^\circ$). Consequently, the changes in the Jovian longitude (λ_{III}) does not impact satisfaction of the diamagnetic drift condition.

In Figure 24, we evaluate the bulk flow shear onset condition when Ganymede lies at the Jovian plasma sheet's centre ($\lambda_{III} = 248^\circ$). The flow shear condition imposes an upper limit on the shear velocity, defined by local magnetic field strengths and plasma mass densities. Figure 24a reveals the primary outflow velocity on the upstream magnetopause. Only southward outflow vectors are shown here for clarify of figure, but each southward vector has an equivalent, antiparallel northward vector. The average outflow speed is $\langle v_{out}\rangle = 312$ km/s and peaks near the subflow region, where adjacent magnetic fields are strongest and most strongly aligned with the outflow vector direction. A small peak in v_{out} also appears along Ganymede's equator, again likely a product of sharp idealised magnetopause surface.

The outflow-aligned flow shear upper limit (hence defined as $v_{1-2,max}$) is calculated in Figure 24b with average $\langle v_{1-2,max}\rangle = 367$ km/s. The maximum shear shares the same variation pattern as Figure 24a, as the parameter is proportional to v_{out} , B_1 , and B_2 . Figure 24c then shows the outflow-aligned bulk flow shear. Given stagnant Ganymedean-side plasma, flow shear velocities are given by Jovian-side values in Figure 22b. The average outflow-aligned flow shear is $\langle v_{1-2}\rangle = 13.7$ km/s. Largest v_{1-2} values occur near the subflow point, where bulk flow shears are well-aligned with the outflow jets. Whereas along the magnetopause flanks, despite much faster bulk flow shears, have smaller v_{1-2} values because outflow directions are orthogonal to the bulk flows. Regardless, the outflow-aligned flow shear is at least one order

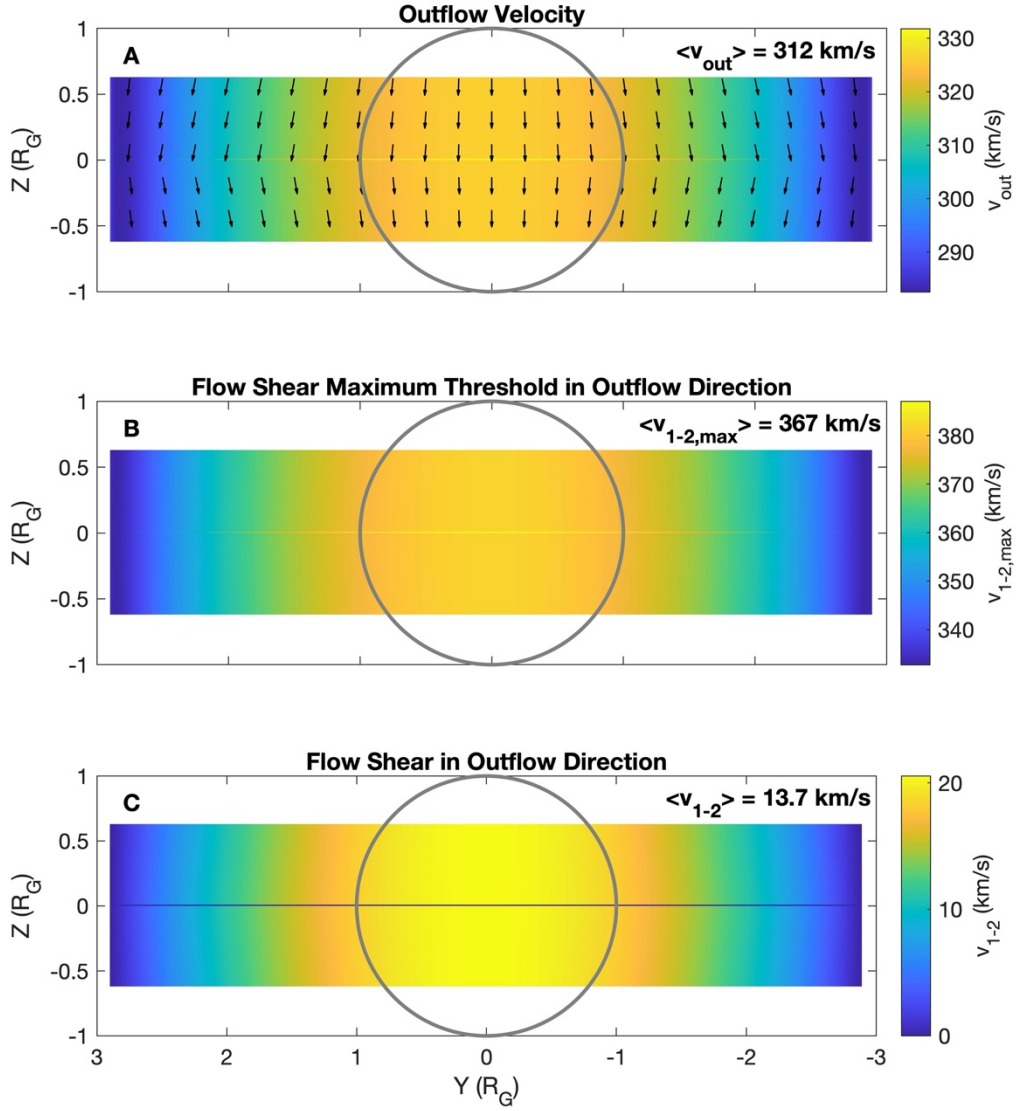


Figure 24: Evaluation of the bulk flow shear reconnection onset condition when Ganymede lies at centre of the Jovian plasma sheet ($\lambda_{\text{III}} = 248^\circ$). Plots shown are (a) outflow velocity, (b) maximum bulk flow shear threshold in outflow direction for reconnection onset, and (c) bulk flow shear in outflow direction. The average parameter value in each plot is denoted on the top-right. Ganymede is outlined in grey.

of magnitude smaller than its upper limit, therefore the bulk flow shear condition for reconnection is well-satisfied throughout the magnetopause.

Just like with the diamagnetic drift case, we evaluate the bulk flow shear condition as Ganymede transitions to higher/lower Jovian magnetic latitudes. At $\lambda_{\text{III}} = 158^\circ$ and $\lambda_{\text{III}} = 338^\circ$, the average outflow velocity increases ($\langle v_{\text{out}} \rangle = 429$ km/s), as the outflow speed is more sensitive to magnetic field strength increase than mass density decrease. The average for flow shear upper limit also increases ($\langle v_{1-2,\text{max}} \rangle = 432$ km/s), contrasting against a reduction seen in average outflow-aligned bulk flow shear ($\langle v_{1-2} \rangle = 11.7$ km/s).

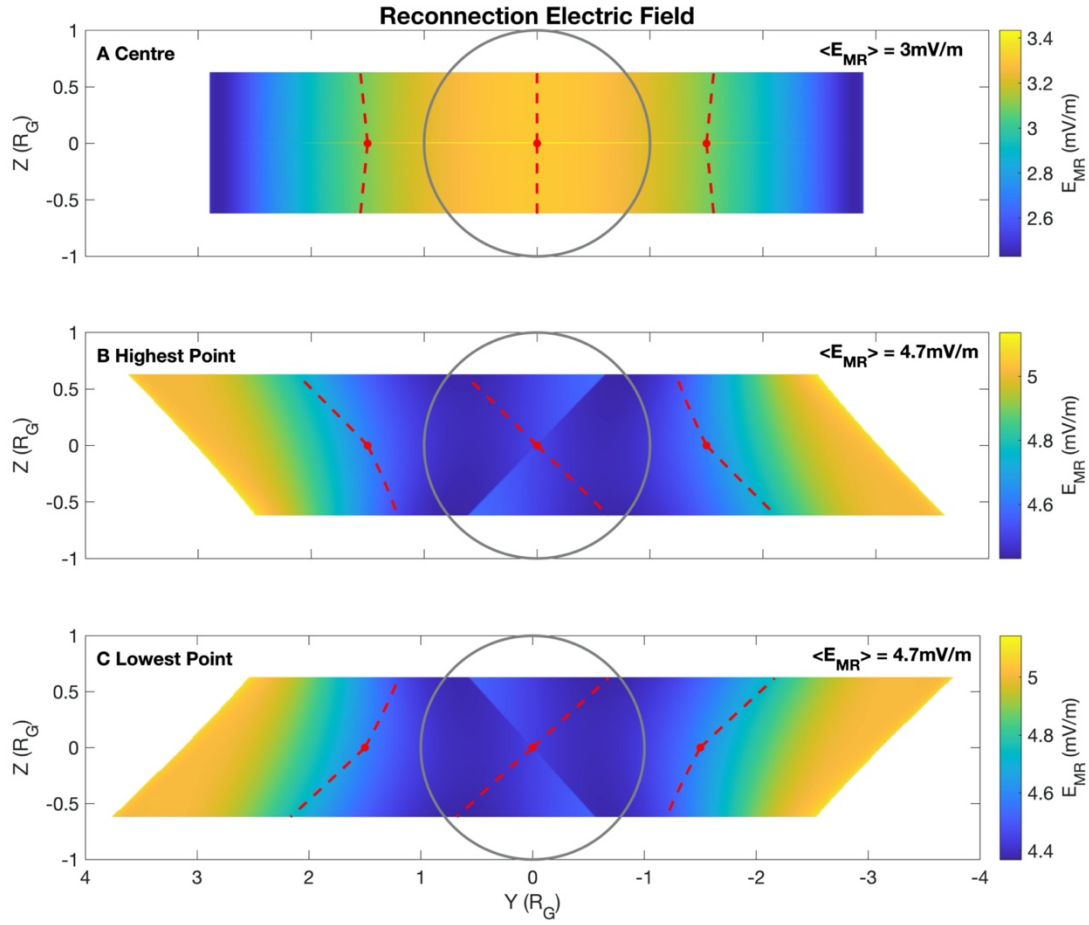


Figure 25: Magnetic reconnection electric field calculated when Ganymede is at (a) centre ($\lambda_{\text{III}} = 248^\circ$), (b) highest latitude ($\lambda_{\text{III}} = 158^\circ$), and (c) lowest latitude ($\lambda_{\text{III}} = 338^\circ$) inside the Jovian plasma sheet. The average parameter value in each plot is denoted on the top-right. Bidirectional motions of tracked outflow plasma parcels from three equatorial reconnection sites ($Y = -1.5, 0, 1.5 R_G$) are shown as red dashed lines. Ganymede is outlined in grey.

Consequently, the changes in Jovian longitude (λ_{III}) also does not impact satisfaction of the bulk flow shear condition.

With all three onset conditions satisfied, magnetic reconnection can occur everywhere on Ganymede's upstream magnetopause for three extreme Ganymede positions in the Jovian plasma sheet. In Figure 25, we calculate the reconnection electric field, which are effectively the unnormalised reconnection rate, at these locations. The choice of $k = 0.1$ is adopted from near-Earth observations as reconnection efficiency has no known plasma beta dependence (Paschmann et al., 2013; Masters, 2017). When the moon is at the plasma sheet's centre ($\lambda_{\text{III}} = 248^\circ$), Figure 25a shows the average reconnection electric field of $\langle E_{\text{MR}} \rangle = 3.0 \text{ mV/m}$ across the symmetric magnetopause. The electric field is strongest near the subflow point,

where adjacent magnetic fields are strongest and reconnection outflows are fastest. Multiplying the average electric field by the magnetopause width ($\sim 6 R_G$) yields an estimate reconnection voltage of $V_{MR} \sim 50$ kV. Red dashed lines denote direction of plasma parcels travelling outward along reconnection outflow tracked by a force balance model¹⁴ (Cowley & Owen, 1989; Cooling et al., 2001). Three equatorial reconnection sites are sampled for tracking – one at the subflow point, and two at $Y = \pm 1.5 R_G$. Plasma parcel paths show reconnection outflows to be predominantly north-south, with some shifts toward nearest magnetopause flanks due to momentum from external Jovian-side plasma flows.

When Ganymede moves to highest and lowest Jovian magnetic latitudes, the average reconnection electric field increases to $\langle E_{MR} \rangle = 4.7$ mV/m in Figure 25b ($\lambda_{III} = 158^\circ$) and Figure 25c ($\lambda_{III} = 338^\circ$) as the magnetopause gains maximum north-south asymmetries. The stronger electric field corresponds to $V_{MR} \sim 74$ kV. Plasma parcels are also tracked in these figures, illustrating similar influence of Jovian-side flows in case of asymmetric magnetopauses.

With Figure 25, we have completed fundamental assessments for magnetic reconnection on Ganymede's upstream magnetopause. However, additional interpretation is still required to understand the nature of Ganymede reconnection, which will be discussed in the next section.

3.4. The Nature of Ganymede reconnection

According to our analytical model, plasma and magnetic conditions are favourable for magnetic reconnection across the entirety of Ganymede's upstream magnetopause. This conclusion holds true when the moon is at centre, highest latitude, and lowest latitude in the Jovian plasma sheet. Away from these extreme positions, each plasma and magnetic parameter oscillates inside a range bookended by its value at the plasma sheet's centre, and at one of the maximum magnetic latitudes. Hence, other Ganymede locations can be

¹⁴ The model offers equation for plasma parcel velocities tangential to the magnetopause for northward and southward reconnection outflows. From each sampled reconnection site, I computed parcel position change vector over 0.1 s (for detailed tracks), and then projected the vector onto the magnetopause to find the parcel's next location. This process was then iterated until the parcel departed the magnetopause domain.

considered ‘intermediate’ cases. Since all three reconnection onset conditions are well satisfied in extreme cases, they must also be satisfied in intermediate cases. Consequently, reconnection can occur everywhere on Ganymede’s magnetopause at all times. This finding is consistent with numerical simulations, whose reconnection events are found to be widely spread in locations (Jia, Walker, et al., 2010; Tóth et al., 2016).

In Figure 25, Ganymede’s reconnection electric field ranges between $E_{MR} = 2.5 - 5.1$ mV/m. For comparison, reconnection electric field strengths near planetary magnetopauses are listed in Table A. The Ganymede values exceed those at other planets (except Mercury) by multiple orders of magnitude, indicating significant (unnormalised) reconnection rates

Table A: Analytically derived magnetic reconnection electric fields (E_{MR}) for Solar System dayside/upstream magnetopauses

System	E_{MR} (mV/m)
Ganymede	2.5 – 5.1
Mercury ^a	0.3 – 3
Earth ^b	<0.5
Jupiter ^c	0.01 – 0.03
Saturn ^d	0.01
Uranus ^e	0.03
Neptune ^f	0.01

Sources: ^aGershman et al., 2016; ^bEggington et al., 2020; ^cMasters, 2017; ^dMasters 2015a; ^eMasters, 2014; ^fMasters, 2015b

throughout the moon’s magnetopause. However, this also means that a dominant reconnection path or structure cannot be identified from the electric field maps. At a given time, Ganymede reconnection can be one long X-line, multiple X-lines, or chaotic spreads of transient flux-transfer events.

Despite its unconstrained nature, magnetic reconnection at Ganymede can be characterised by variation of the average electric field, which is higher when the moon is away from the centre of the Jovian plasma sheet (Figure 25). Eq. 3.15 for the electric field is most sensitive to changes in adjacent outflow-aligned magnetic fields B_1 and B_2 . The former depends on the

Jovian-side field B_J , which increases monotonically with Ganymede's distance from the plasma sheet's centre. The latter depends on the Ganymede-side field B_G , which increases alongside B_J to maintain fixed magnetopause surface in our analytical model. Magnitude of the electric field is therefore governed by Ganymede's orbital position and the Jovian System-III longitude (λ_{III}). Then, given Jupiter's rapid rotation compared to Ganymede's orbital period, each λ_{III} approximately corresponds to a distinct time of day on Jupiter. Thus, the average reconnection rate at Ganymede exhibits a Jovian-diurnal variation period and is effectively driven by the planet's rotation. This result has been independently supported by remote observations of radio emissions from Ganymede's interaction footprints in the Jovian polar regions (Zarka et al., 2018).

Magnetic reconnection is an exothermic process that heats up local ions and electrons around the diffusion regions. Particle temperature changes induced by reconnection can be empirically related to the outflow speed by (Phan, Shay, et al., 2013; Phan et al., 2014)

$$\Delta T_i = 0.13 m_i (v_{out})^2 \quad (3.30)$$

$$\Delta T_e = 0.017 m_i (v_{out})^2 \quad (3.31)$$

where $m_i = 14$ amu is the local ion mass. Eq. 3.30 and 3.31 generates $\Delta T_i = 2000 - 4200$ eV for ions and $\Delta T_e = 250 - 560$ eV for electrons, with minimum (maximum) values corresponding to when Ganymede is at centre (highest/lowest latitude) of the Jovian plasma sheet. These temperature changes exceed upstream Jovian ion and electron temperatures of 60 eV and 300 eV respectively (Kivelson et al., 2004). Heating from Ganymede reconnection should therefore be detectable to future in-situ observations.

The analytical model used in this reconnection assessment has used some simplifications to expedite modelling efficiency. Unlike in numerical simulations, the magnetopause position is fixed. The model does not re-calculate boundary position after a change in pressure balance, instead adjusting the Ganymede-side magnetic field strength to maintain the same magnetopause. This simplification does not change boundary-adjacent magnetic field strengths (as they are positionally independent), and therefore should not impact our results. However, fixing the magnetopause means B_G is dependent on B_J , and both magnetic

parameters are dependent on ρ_J and P_J . Hence, we cannot isolate model sensitivity to individual plasma-magnetic parameters.

In this analysis, Jovian-side plasma and magnetic properties are parametrised by simple sinusoidal functions of the flaring angle, which adequately captures compression and deflection arising from idealised horizontal inflow. However, real upstream Jovian plasma is not perfectly laminar or uniformly directed along the GphiO X-axis. This can result in more vertical (Z-dependent) or asymmetric magnetopause compressions that are not accounted for single sine/cosine relations. Although more complex parametrisations can be introduced, they are unlikely to significantly impact plasma-magnetic parameter magnitudes, and therefore our results given that both diamagnetic and bulk flow shear onset conditions are extremely well satisfied.

The model in this chapter takes a set of defined constant upstream inputs and produces snapshot pictures of steady-state Ganymedean magnetopauses. Changes in plasma-magnetic conditions as Ganymede travels up-down inside the Jovian plasma sheet are not considered. Accounting for such changes will be an important step toward understanding temporal dynamics around the moon, including the effects of magnetic reconnection on plasma convection. This model development will be discussed more fully in Chapter 4.

In conclusion, magnetic reconnection has been determined as fundamentally viable throughout Ganymede's closed-field magnetopause, with its average rate controlled by Jupiter's rotation. These findings significantly advance our understanding of Ganymedean reconnection, but many open questions still remain. For instance, can reconnection structure be further constrained? Reconnection structures can be steady or transient, and onset assessment does not predict which of the two is dominant on the Ganymedean magnetopause. Furthermore, reconnection has a driving role on Ganymede's Dungey-like plasma convection cycle, and knowledge of periodic upstream reconnection rate may be vital for improved understanding of not only convection dynamics, but also downstream reconnection in Ganymede's magnetotail. These open areas of research will be the primary foci of our next chapter.

Chapter 4: Effects of Magnetic Reconnection on Ganymede's Magnetosphere

This chapter provides further exploration of Ganymede's upstream magnetic reconnection and its driving influence on plasma convection around the moon. My two-part analysis is conducted through the lens of magnetic flux opening, transport, and closure. The first half surveys flux-transfer events (FTEs) from MHD-EPIC simulations of Ganymede, and through computation of their flux contents measures localisation of the moon's upstream reconnection. Each FTE is shown to contain $<1 - 15\%$ of total open flux in Ganymede's polar caps, and FTEs can collectively contribute $\sim 40\%$ of upstream flux opening at any particular time. Localised reconnection bundles are therefore significant, perhaps even majority, contributors of open upstream flux, and their detection should be a focus for future in-situ observations of the Ganymedean magnetopause.

The second half of this chapter analyses magnetic flux transport to characterise Ganymede's magnetospheric plasma convection. Most importantly, Ganymede's Dungey cycle must be rapid given upstream magnetic reconnection's ability to open up closed magnetic field lines within minutes. Then, to maintain relatively steady locations of OCFBs and aurora ovals, downstream reconnection must consistently generate significant voltage to provide balancing magnetic flux closure. Finally, Ganymede's OCFBs and/or aurora ovals will likely exhibit periodic variation if upstream and downstream reconnection voltages are approximately equivalent. Although this result may require a predominantly steady global reconnection structure, it has been supported by resistive MHD simulation results which show ample evidence of transient FTEs.

4.1. Magnetic Flux and its Transport

Assessment of idealised steady-state plasma conditions near Ganymede's upstream magnetopause shows magnetic reconnection to be highly viable, with widespread large characteristic occurrence rates that leave its global structure unconstrained. Furthermore, the average reconnection rate periodically varies with Ganymede's magnetic latitude inside the Jovian plasma sheet, and is therefore governed by Jupiter's rotation. Given that

Ganymede likely possesses a Dungey-like plasma convection pattern, upstream reconnection is the primary driver of plasma-magnetic dynamics in the moon magnetosphere. In this chapter, I will explore the consequences of these findings, and determine whether I can use upstream reconnection to generate more information on Ganymede's still mysterious magnetosphere.

My subsequent analysis relies on the concept of magnetic flux, which describe the amount of magnetic field that 'flows' through a surface area

$$\Phi = \iint \mathbf{B}(\mathbf{r}) \cdot d\mathbf{A} \quad (4.1)$$

where $\mathbf{B}(\mathbf{r})$ is a spatially variable magnetic field and $d\mathbf{A}$ describes an infinitesimal element of a surface defined by orthogonal area vector $\mathbf{A}$. One flux unit Weber (Wb) is defined by a magnetic field of 1 T strength crossing through a surface of area 1 m². The magnetic flux is non-zero if its defining surface does not enclose a finite volume¹⁵.

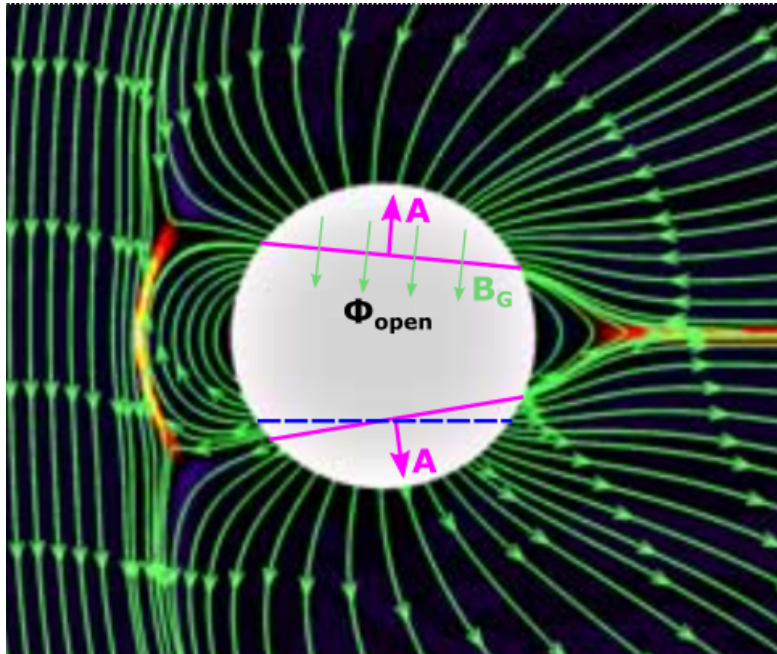


Figure 26: A simple annotation of Ganymede's polar caps and the concept of open magnetic flux. The background image is taken from global resistive MHD simulation by Jia et al., (2008), *Journal of Geophysical Research*.

¹⁵ In this case, the contour is said to be 'open', but the term will not be used in this document to prevent confusion with open magnetic field lines.

Ganymede's cylindrical magnetosphere possesses two open polar caps, each confined by the open-closed field boundary (OCFB) that divides open and closed magnetic field lines produced by the moon's internal dynamo. Figure 26 illustrates open magnetic field lines flowing through the polar caps, producing open flux Φ_{open} for the magnetospheric system. Since the same field lines flow through both surfaces at all times, Φ_{open} is calculated using only one polar cap to avoid double counting. For the rest of this document, I will use the northern polar cap as the primary flux surface.

The open magnetic flux represents the 'amount' of open magnetic field lines in Ganymede's magnetosphere. Φ_{open} increases when closed magnetic field lines are opened and decreases when open magnetic field lines are closed. In a Dungey-like cycle, flux opening and closure are achieved by magnetic reconnection at the magnetopause and in the magnetotail respectively (Dungey, 1961). As the result, the change in open flux $\Delta\Phi_{\text{open}}$ within a time interval Δt is given by

$$\Delta\Phi_{\text{open}} = \Phi_{\text{up}} - \Phi_{\text{dn}} \quad (4.2)$$

where Φ_{up} and Φ_{dn} are fluxes open and closed by upstream and downstream reconnection respectively. Eq. 4.2 represents the conservation of magnetic flux and documents changes to the magnetosphere's flux budget. The equation can be rearranged such that if two of its terms are known, then the third term can be calculated.

I will first focus on magnetic flux opening by upstream reconnection. The flux opening exhibits a simple relationship with the reconnection voltage such that

$$\Phi_{\text{up}} = \int_t^{t+\Delta t} V_{\text{up}}(t)dt \quad (4.3)$$

where $V_{\text{up}} = V_{\text{MR}}$ from Chapter 3 and Δt is the again the time interval considered. The voltage is computable from the reconnection electric field ($E_{\text{up}} = E_{\text{MR}}$) through Eq. 3.16 if the global reconnection structure is predominantly constrained along one or more X-lines. A variant of Eq. 4.3. has been applied to Earth's and Saturn's magnetopause, in which the right-hand side is rewritten as a function of external plasma flow speed, magnetic field strength, and magnetic field orientation – all of which are obtained empirically from observational data (e.g., Jackman et al., 2004; Milan et al., 2007).

The Kronian application does not account for reconnection in the polar lobes, and vitally assumes that reconnection occurs along an X-line, which is not always the case for other magnetopauses, like Ganymede's, where reconnection is expected to be widespread and localised (Jia, Walker, et al., 2010). Zhou et al., (2020) estimates reconnection voltage in Eq. 4.3. by defining an upstream half-open flux surface (shown in Figure 27), then evaluating magnetic flux changes through this surface, arriving at $V_{\text{up}} = 82.9 \pm 18.3$ kV for Hall-MHD model and $V_{\text{up}} = 83.6 \pm 8.2$ kV for MHD-EPIC model. Both value derivations are reliant on global numerical simulations, which give imperfect accounts of reconnection microphysics.

Localised magnetic reconnection is known as flux-transfer events (FTEs)¹⁶, distinguishable via bimodal oscillations in the magnetic field component normal to the current sheets from which they originate (e.g., Paschmann et al., 2013). The formation of FTEs at a magnetopause remains under much debate.

Three popular mechanisms are illustrated in Figure 28. Model A visualises an FTE as two

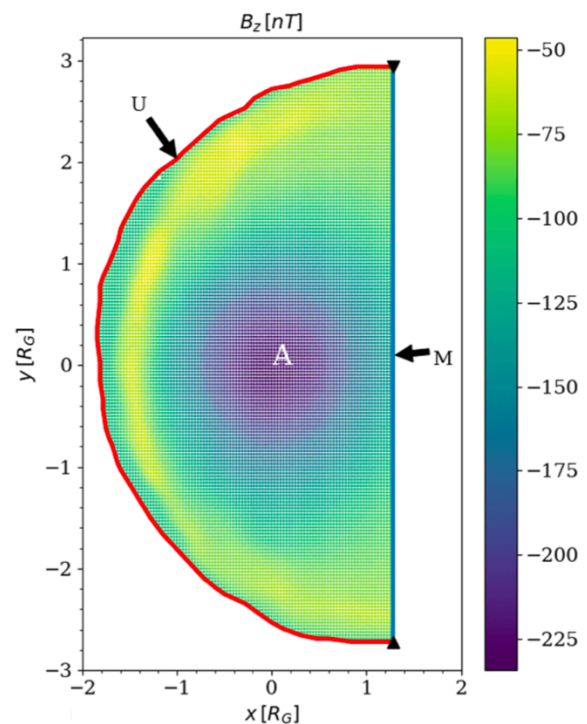


Figure 27: Half-open magnetic flux surface used to calculate upstream reconnection voltage. Image courtesy of Zhou et al., (2020), *Journal of Geophysical Research: Space Physics*.

¹⁶ Not all FTEs arise from reconnection, which requires a local magnetic null point ($|B| = 0$). For magnetopause-like boundaries, however, reconnection and FTEs can be assumed correlated.

narrow elbow-shaped tubes that connect magnetic field lines across the magnetopause. These tubes are thought to arise from intermittent reconnection events independent of larger X-lines (Russell & Elphic, 1979). Meanwhile, Model B proposes a ‘flux ropes’ structure, whose twisted field lines are produced by ‘bursty’ reconnection events above and below an X-line (Scholer, 1988; Southwood et al., 1988a). Lastly, Model C shows FTE formation when two or more X-lines in close proximity merge into a long and wide tube structure (Lee & Fu, 1985; Raeder, 2006).

FTEs have been most extensively studied at Earth with numerous in-situ spacecraft (e.g., Russell & Elphic, 1979; Wild et al., 2005; Hasegawa et al., 2010), alongside indirectly evidence of magnetic field opening from the planet’s ionospheres (e.g., Lockwood et al., 1990b; Milan et al., 2000; Oksavic et al., 2004). These events are shown to have a typical lifetime of ~ 8 minutes (Rijnbeek et al., 1984), with evidence for all three models of formation (e.g., Fear et al., 2007; Øieroset et al., 2011; Varsani et al., 2014). In fact, FTEs of Models B and C can occur concurrently around the same X-line (Fear & Milan, 2012; Trenchi et al., 2016). Away from Earth, minutes-long FTEs have been observed at Jupiter (Walker & Russell, 1985; Huddleston et al., 1997) and Saturn (Jasinski et al., 2016), albeit much less frequently due to lesser prevalence of magnetopause reconnection. In smaller magnetospheres, FTEs have been

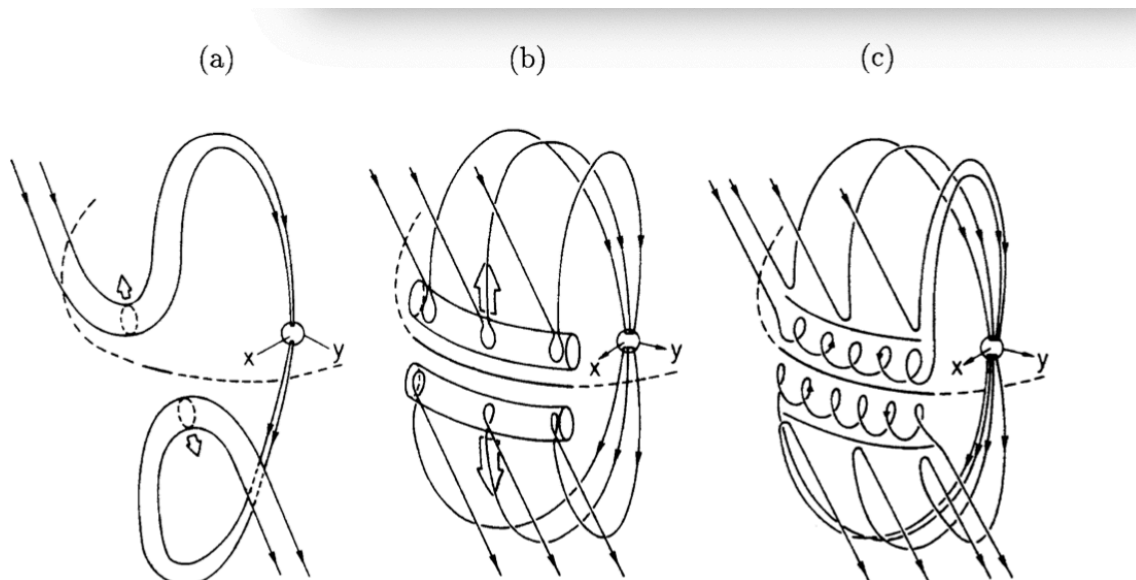


Figure 28: Diagrams illustrating three popular models of FTE formation. (a) Elbow-shaped flux tubes from transient magnetic reconnection extending across the magnetopause (dashed line), (b) Flux ropes (shown as tubes here) above and below one cross-magnetopause X-line (solid line) arising from ‘bursty’ reconnection, and (c) Large flux rope/tube formed via merging of two or more X-lines (solid line) in close proximity. Image courtesy of Paschmann et al., (2013), *Space Science Review*.

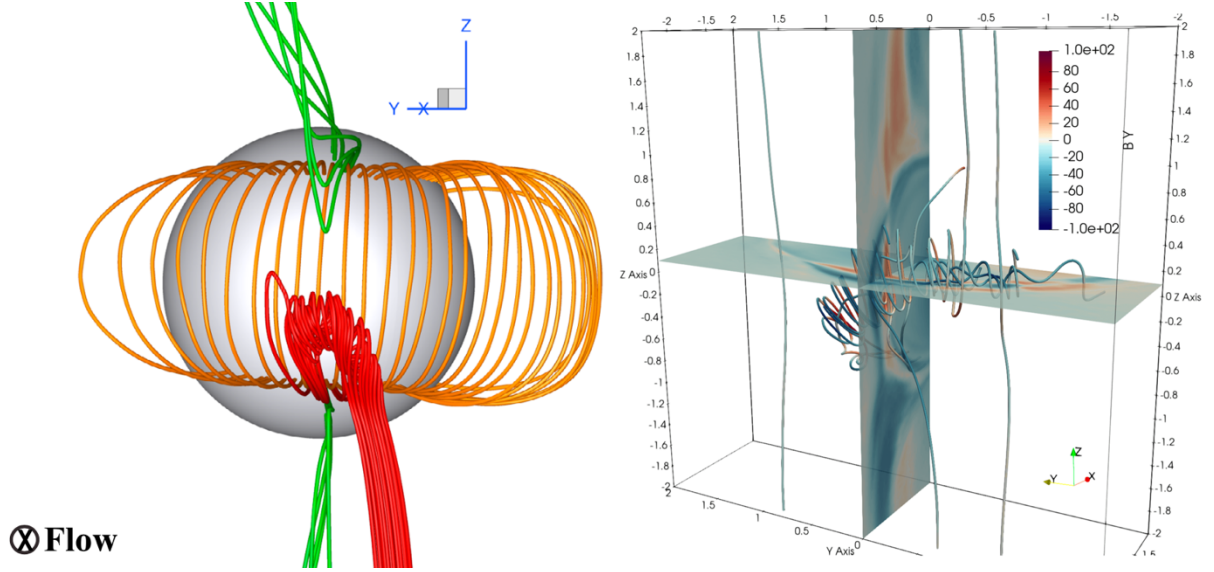


Figure 29: Examples of twisted magnetic flux rope/tube from (a) resistive MHD and (b) MHD-EPIC numerical simulations. The FTEs are shown in red and blue-orange lines respectively. Images courtesy of Jia et al., (2010b), *Journal of Geophysical Research* and Zhou et al., (2020), *Journal of Geophysical Research: Space Physics*.

abundantly observed at Mercury (Russell & Walker, 1985; Slavin et al., 2009; Imber et al., 2014), but with much shorter lifetime of ~ 10 seconds (Imber et al., 2014).

Short-living FTEs have not been observed directly at Ganymede, but they are seen in both resistive MHD and hybrid MHD-EPIC simulations (Jia, Walker, et al., 2010; Zhou et al., 2019; Zhou et al., 2020). The resistive MHD model Ganymedean FTEs are predicted to be widespread with lifetimes in the tens of seconds, forming localised flux events (Model 1) of cross-section radius $r \sim 0.2 R_G$ (Jia, Walker, et al., 2010). One resistive MHD flux tube is traced by red magnetic field lines in Figure 29a. Meanwhile, the MHD-EPIC model sees FTEs altering conditions near long X-lines (Models 2/3) with production rate every 20 – 50 s (Zhou et al., 2020). A flux rope example from MHD-EPIC modelling is shown in Figure 29b. Despite varying prediction on FTE structures, both models consider transient FTEs to be prominent over steady X-lines even though upstream Jovian plasma conditions are steady.

The amount of magnetic flux opened by each FTE depends on its structure. The flux inside a smaller Model A flux tube can be estimated by (Fear et al., 2017)

$$\Phi_{\text{tube}} = B_{\text{axial}} \pi r^2 \quad (4.4)$$

where B_{axial} is magnetic field strength along the elbow-shaped tube and r is the tube's radius. For Model 1, Eq. 4.4 represents all flux opened by transient reconnection. For Models B and C, however, Φ_{tube} is only a minority of all flux opened by reconnection. A revised expression is therefore introduced for an FTE based on X-line(s) (Fear et al., 2017)

$$\Phi_{\text{FTE}} = rL_X B_N \quad (4.5)$$

where r is the FTE radius and L_X is the length of FTE-generating X-line(s). Product $r \times L_X$ represents the area on the magnetopause threaded by open magnetic field lines of boundary-normal strength B_N , providing a comprehensive capture of all open flux. This flux content can be compared to the total open flux in the magnetosphere by computing ratio

$$F = \frac{\Phi_{\text{FTE}}}{\Phi_{\text{open}}} \quad (4.6)$$

which yields an estimate for the FTEs' importance in driving the Dungey cycle. Generally, FTEs are considered significant drivers in a magnetospheric system if $F > 10\%$. Numerous FTE studies at Earth's magnetopause suggests varying open flux contribution between $F \sim 1 - 10\%$, leaving FTE's role in plasma convection inconclusive (Fear et al., 2017). Meanwhile, the FTE open flux contribution is much greater at Mercury's magnetopause where $F \sim 15\% - 25\%$ (Fear et al., 2019), indicating plasma convection significantly driven by transient magnetic reconnection.

Given intimate entanglement between FTEs and X-lines (Models B and C), separating X-line flux contribution to Φ_{FTE} is infeasible unless the two structures are perfectly orthogonal. This is difficult to determine from observations or modelling, especially for magnetopause systems where reconnection occurs frequently. Hence, Φ_{FTE} cannot be partitioned into transient and steady components, and only states how much flux is contained in localised reconnection bundles.

Returning to Eq. 4.2, the remaining unknowns are change in open magnetic flux ($\Delta\Phi_{\text{open}}$) and flux closure by downstream reconnection (Φ_{dn}). The former is derivative from the total open magnetic flux in the magnetospheric polar caps where

$$\Delta\Phi_{\text{open}} \approx \Phi_{\text{open}}(t + \Delta t) - \Phi_{\text{open}}(t) \quad (4.7)$$

For magnetic systems of Earth, Mercury, and Saturn, Φ_{open} has been tracked via in-situ and remote observations, which alongside empirical models produce open flux estimates in Table B.

Table B: Estimates of open magnetic flux in planetary polar caps.

System	Φ_{open} (MWb)
Earth ^{a,b}	~600
Mercury ^c	4 – 6
Saturn ^d	~40,000

^aMilan et al., 2007; ^bHuang et al., 2009; ^cAlexeev et al., 2010; ^dBadman et al., 2005

The subsequent value of open flux change depends on the time interval considered. Once $\Delta\Phi_{\text{open}}$ is calculated, Eq. 4.2 will yield the downstream flux closure (Φ_{dn}), which is convertible to downstream reconnection voltage (V_{dn}) via Eq. 4.3 (or similar variants). Analysis of flux closure has been carried out for magnetospheres of Earth, Mercury, and Saturn (Badman et al., 2005; Milan et al., 2007; Imber et al., 2017). Derivation of V_{dn} represents an important characterisation of magnetotail reconnection, and opens the gateway toward understanding of large-scale downstream dynamics.

Unfortunately, neither open polar magnetic flux or downstream flux closure has been computed for Ganymede – whether by observations or simulations of its magnetosphere. With only one known parameter, Eq. 4.2. cannot be solved in its current state. One of $\Delta\Phi_{\text{open}}$ or Φ_{dn} must be calculated, or otherwise assumed, to provide constraint on Ganymede’s downstream reconnection. Hence, the moon’s plasma convection picture remains profoundly unclear without additional research.

Earlier in this section, the open magnetic flux surface is defined encircled by the OCFB. As Φ_{open} increases (decreases) in a magnetosphere, size of the body's polar caps increases (decreases) accordingly – shifting the OCFB toward the body's equator (magnetic poles). This process can be observed visually through aurora ovals, which were briefly discussed in Chapter 2 for Dungey-like magnetospheres. The complete physics of auroral emission is complex and unique for each celestial body, but the general principle involves excitation of atmospheric molecules by bombardment of magnetospheric plasma returning to the body after undergoing magnetic reconnection in the magnetotail. With their visibility, aurora ovals serve as important features for remote documentation of plasma convection dynamics.

Given complexity of auroral physics, I will focus my review on Ganymede's magnetosphere. In-situ and remote observations of Ganymedeaurora show the ovals to correlate closely in location with modelled OCFBs, albeit with minor differences attributed to methodical inaccuracies and unaccounted interactions with Jupiter's magnetosphere (Khurana et al., 2007; McGrath et al., 2013). The proxy relationship can be assumed if the auroral emission occurs near the moon's surface (Saur et al., 2015; Musacchio et al., 2017), which will be taken as true in this document.

Previously in Figure 17, I show snapshots of Ganymedeauroras taken by the Hubble Space Telescope (HST). Due to limited exposure intervals, the HST does not capture entire aurora ovals with each measurement pass (McGrath et al., 2013). Hence, images across long time periods (months/years) are often summed together to study the aurora's global-scale morphology. Data from the HST reveals a highly tilted oval structure, where the aurora is significantly pushed poleward on the upstream side and equatorward on the downstream side (McGrath et al., 2013; Saur et al., 2015).

I quantify the tilt by defining spherical coordinates centred at Ganymede where r is radial distance, θ is the co-latitude measured from the rotation axis (assumed parallel to GphIO Z-coordinates), and ϕ is the longitude on the moon's surface. The aurora ovals have upstream co-latitude $\theta \sim 40^\circ - 50^\circ$ and downstream co-latitude $\theta \sim 70^\circ - 80^\circ$ (McGrath et al., 2013). The significant asymmetry is thought to arise from sub-Alfvénic plasma-magnetic interactions

between Ganymede's and Jupiter's magnetospheres (Neubauer, 1998; Jia et al., 2009). Note that smaller but still significant auroral tilts are also observed for Mercury despite less notable sub-Alfvénic effects (Gershman et al., 2015). Since asymmetries of aurora ovals do not impact density of magnetic field lines passing through the polar caps, the amount of open magnetic flux will not change if tilted aurora ovals are replaced by 'representative' rings at constant co-latitude (θ_0). This oval is denoted by the blue dashed line in Figure 26. Recalling the expression for magnetic dipole field in Eq. 1.7, I set $g_1^0 = -719$ nT and integrate the equation through the oval surface with $r = 1 R_G$, $\theta = \theta_0$, and $0 \leq \phi \leq 2\pi$ to obtain¹⁷

$$\Phi_{\text{open}} = -719 \frac{\mu_0 \sin(\theta_0)}{2R_G} \quad (4.9)$$

where $R_G = 2634$ km and $\mu_0 = 4\pi \times 10^{-7}$ H/m. The open flux expectedly increases as the oval co-latitude increases toward Ganymede's equator ($\theta_0 = 90^\circ$). Considering the asymmetry observed by HST, I take as example $\theta_0 = 55^\circ$ for the northern aurora oval. This co-latitude yields a polar cap open flux of $\Phi_{\text{open}} = -12.83$ MWb, with the negative sign accounting for the fact that the magnetic field lines are southward and antiparallel to the northern polar cap area vector.

A study of aurora morphology through summed images discards transient variations at hours timescale. Although multiple analyses have investigated aurora oval motions with respect to Ganymede's position inside the Jovian plasma sheet (Saur et al., 2015; Musacchio et al., 2017), these studies are limited by HST's inability to capture Ganymede at regular Jovian magnetic latitude intervals, and therefore unable to visualise periodic vertical OCFB oscillations likely arising from orderly magnetic flux transport.

Equipped with relevant understanding of magnetic flux and its transport, the rest of this chapter will cover my novel research in two parts. In Section 4.2., I complete a survey of localised FTEs on Ganymede's magnetosphere based on MHD-EPIC simulation results and

¹⁷ Alternatively, one can use Stokes' law to rewrite $\oint \mathbf{B} \cdot d\mathbf{S} = \oint \mathbf{A} \cdot d\mathbf{l}$, where vector potential $\mathbf{A}$ is derived via one-dimensional Biot-Savart law assuming $r' \gg r$ and $d\mathbf{l}$ is in ϕ direction. This is the method I implemented in MATLAB.

calculate their magnetic flux contents. Then in Section 4.3., I evaluate Ganymede's magnetic flux transport through Eq. 4.2. and determine potential motions of the OCFBs for multiple scenarios of yet-unknown magnetic flux closure in the moon magnetotail. Both studies will further our understanding of upstream magnetic reconnection and help us obtain a clearer picture of plasma-magnetic convection in the moon's magnetic domain.

4.2. Survey of Flux-Transfer Events from Ganymedean Numerical Simulations

Given magnetic reconnection favourability of Ganymede's upstream magnetopause, extended X-lines can form at any point and time on the current surface. While there is no indication that Ganymedean FTEs prefer X-line based formation models over elbow-shaped flux tubes, I shall consider Φ_{FTE} from Models B/C to represent the upper limits of magnetic flux contained by localised FTEs. Three parameters in Eq. 4.5 are required to compute open flux of an FTE – the typical FTE radius (r), the length of X-line that produces the FTE (L_x), and the magnetopause-normal magnetic field component inside the FTE (B_N).

For this study, I analyse outputs from one MHD-EPIC Ganymede simulation ran by Zhou et al., (2020). The simulation takes plasma-magnetic inputs corresponding to when Ganymede lies at the centre of the Jovian plasma sheet (and Galileo G8 flyby) and operates for 20 minutes, with 1 s iteration interval after the upstream magnetopause reaches a steady state. For each model iteration, the magnetopause can be extracted by determination of the isosurface where $B_z = 0$ in GphiO coordinates (see Chapter 2, Section 2.2 for explanation). At this point, a LMN local coordinate system can be defined such that N is along the local magnetopause normal vector, L runs parallel to the surface toward Ganymede's northern magnetic pole, and M completes right-hand orthogonality. Magnetic field can then be projected onto LMN components, and a magnitude map of B_N on the magnetopause isosurface is produced as exemplified in Figure 30. Such snapshots obtained for 71 timesteps in which bimodal B_N oscillation is observed, which are then compared to Figure 6b in Zhou et al., (2020) to confirm detection of FTEs. The process in this paragraph is conducted and their outputs kindly given by Hongyang Zhou.

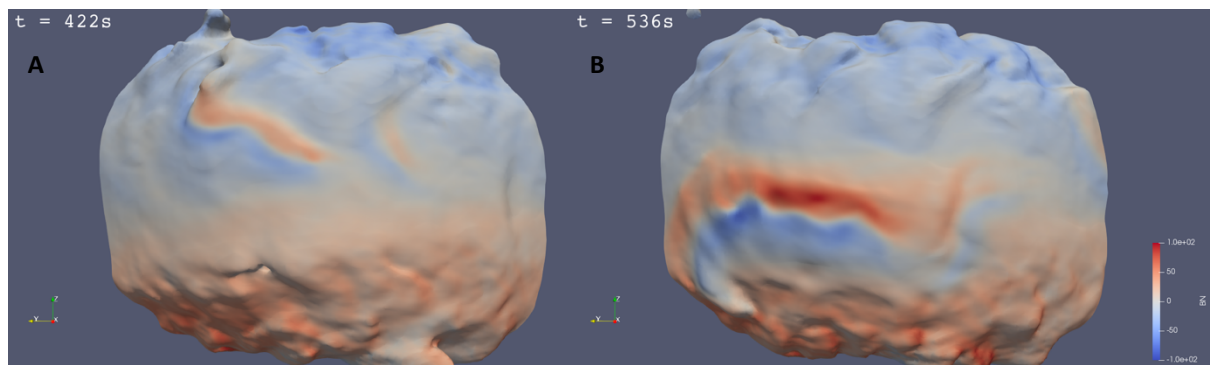


Figure 30: Two FTEs from MHD-EPIC B_N data. Bimodal magnetic oscillations are seen in forms of isolated red patches on the magnetopause. A comparatively small and faint FTE is shown in the left (a) figure, contrasting from its large and prominent counterpart in the right (b) figure. Image courtesy of Hongyang Zhou, personal communication, November 27, 2020.

For each timestep, the most prominent bimodal magnetic pattern is spotted and its oscillation is confirmed via corresponding tabulated magnitude data files. Since I investigate an upper limit of FTE flux content, I choose peak B_N within the structure to represent the reconnection bundle, except for certain occasions where the magnetic peak appears numerically artificial. In such cases, a ‘representative’ B_N value is chosen within 1 – 2 nT of the peak magnitude. A histogram illustration of the B_N distribution amongst the 71 studied FTEs are shown in Figure 31.

The distribution has the mean of 47 nT, the median of 40 nT, and the mode of 35 nT. Discrepancies among the three central values indicate the distribution’s skewness toward smaller B_N values. This finding is consistent with visual inspection that reveals more frequent occasions of smaller, fainter FTEs (Figure 30a) compared to their larger, stronger counterparts (Figure 30b). The relative rarity of large FTEs is likely due to the general frequency of reconnection, resulting in new events tearing up existing FTEs before they can grow in size (Jia, personal communication, July 16, 2021). The largest of Ganymedean FTEs can extend up to half the magnetopause width, but it must be noted that the boundary in Zhou et al., (2020) is significantly narrower than the idealised analytical surface from Kivelson et al., (1996) used in Chapter 3 by a factor of ~ 2 . Thus, the lengths of modelled FTEs may not represent true sizes of the structures.

Nevertheless, Figure 31 provides an important magnitude order estimate for magnetic field strength inside an FTE. Even if the other two parameters in Eq. 4.5 cannot be well constrained by numerical simulations, they can be studied as independent variables. Jia et al., (2010b) approximately evaluate a typical FTE radius to be $r \sim 0.2 R_G$. Taking this as a central value, I establish the variable range $0 < r < 0.4 R_G$ for an FTE radius. At its widest, the FTE will span about half the northern/southern hemisphere of Ganymede's magnetopause, which is consistent with structures seen on model snapshots. For the X-line length, the set variable range is $0 < L_x < 8 R_G$. The upper limit is derived by computing the equatorial arc length for a magnetopause of $\sim 6 R_G$ width. These two independent variable ranges are chosen to cover all possible sizes of FTEs, even if many of them are physically unrealistic.

Figure 32 illustrates amounts of magnetic flux inside an FTE of varying width and length, assuming the FTE to be a cylindrical tube of uniform magnetopause-normal magnetic field strength. I consider three values of $B_N = 35, 65, 95$ nT to cover the magnitude spread represented in the Figure 31 histogram. Φ_{FTE} is linearly proportional to L_x and r in each subplot, and proportional to B_N between the three subplots. Consider a 'typical' FTE shown in Figure 30a. If this structure has $B_N = 35$ nT with length $L_x \sim 2 R_G$ and width $r \sim 0.2 R_G$, then Figure 32a reveals open flux content $\Phi_{FTE} \sim 0.1$ MWb. On the other hand, a larger FTE

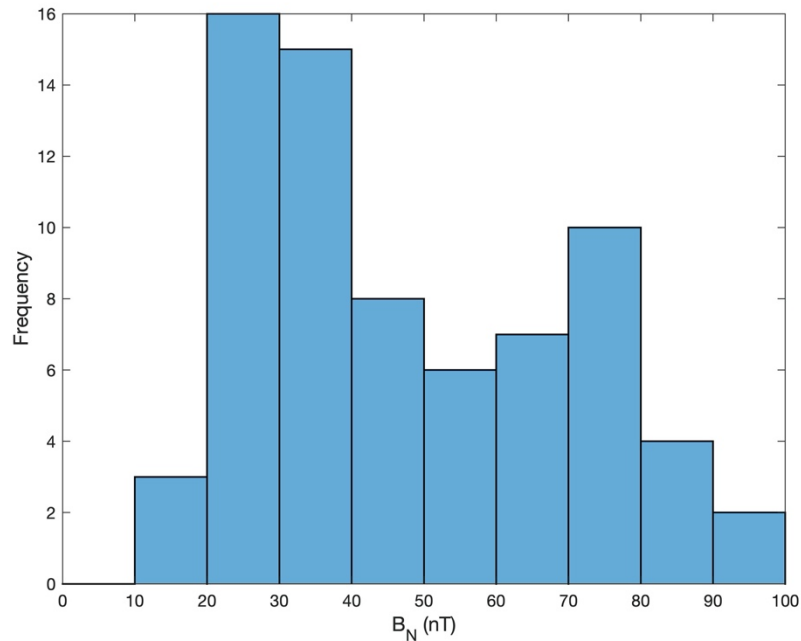


Figure 31: Distribution of magnetopause-normal magnetic field inside 71 FTEs detected from MHD-EPIC model outputs. The data shows two peaks around $B_N \sim 30$ nT and $B_N \sim 75$ nT.

in Figure 30b ($B_N = 95$ nT, $L_x \sim 5 R_G$, $r \sim 0.2 R_G$) possesses up to $\Phi_{FTE} \sim 0.7$ MWb (Figure 32c). The flux upper limit in Figure 32 is $\Phi_{FTE} \sim 2$ MWb, however this is unlikely to be achieved under real conditions.

If the polar cap flux content is taken to be $\Phi_{open} = 12.83$ MWb, I derive ratios $F < 1\%$ and $F \sim 5\%$ for FTEs in Figures 29a and 29b respectively, although the largest possible FTE may produce up to $\sim 10\% - 15\%$ of the total open flux. Hence, large FTEs can be significant drivers of Ganymede's magnetic flux convection at the moment of their formation. Although smaller individual FTEs may not contribute substantial flux, their impacts can rapidly accumulate given the FTE production rate of ~ 3 events/min (Zhou et al., 2020). The effects of largest $\Phi_{FTE} \sim 2$ MWb event, for instance, can be replicated by three successive $\Phi_{FTE} \sim 0.7$ MWb events. The Ganymedeian F values are fairly similar to those in Mercury's magnetosphere, where large FTEs are thought to open $\sim 5\%$ of the planet's total open flux with maximum limits of $\sim 15\% - 25\%$ (Fear et al., 2019).

Using the upstream magnetic reconnection voltage $V_{up} \approx 83$ kV from Zhou et al., (2020), I estimate a flux opening of $\Phi_{up} \approx 5$ MWb per one minute of magnetopause activity. Assuming large FTEs like Figure 30b, these localised structures may contribute up to $\sim 40\%$ of upstream magnetic flux opening, but with smaller typical values due to abundance of small FTEs like Figure 30a. It may be tempting to state that quasi-steady X-lines generate the other 60% of opened magnetic flux. However, such conclusion overlooks the facts that (1) Φ_{FTE} contains an inseparable contribution from X-lines and (2) the MHD-EPIC simulation is not the perfect accountant of all FTEs on Ganymede's magnetopause.

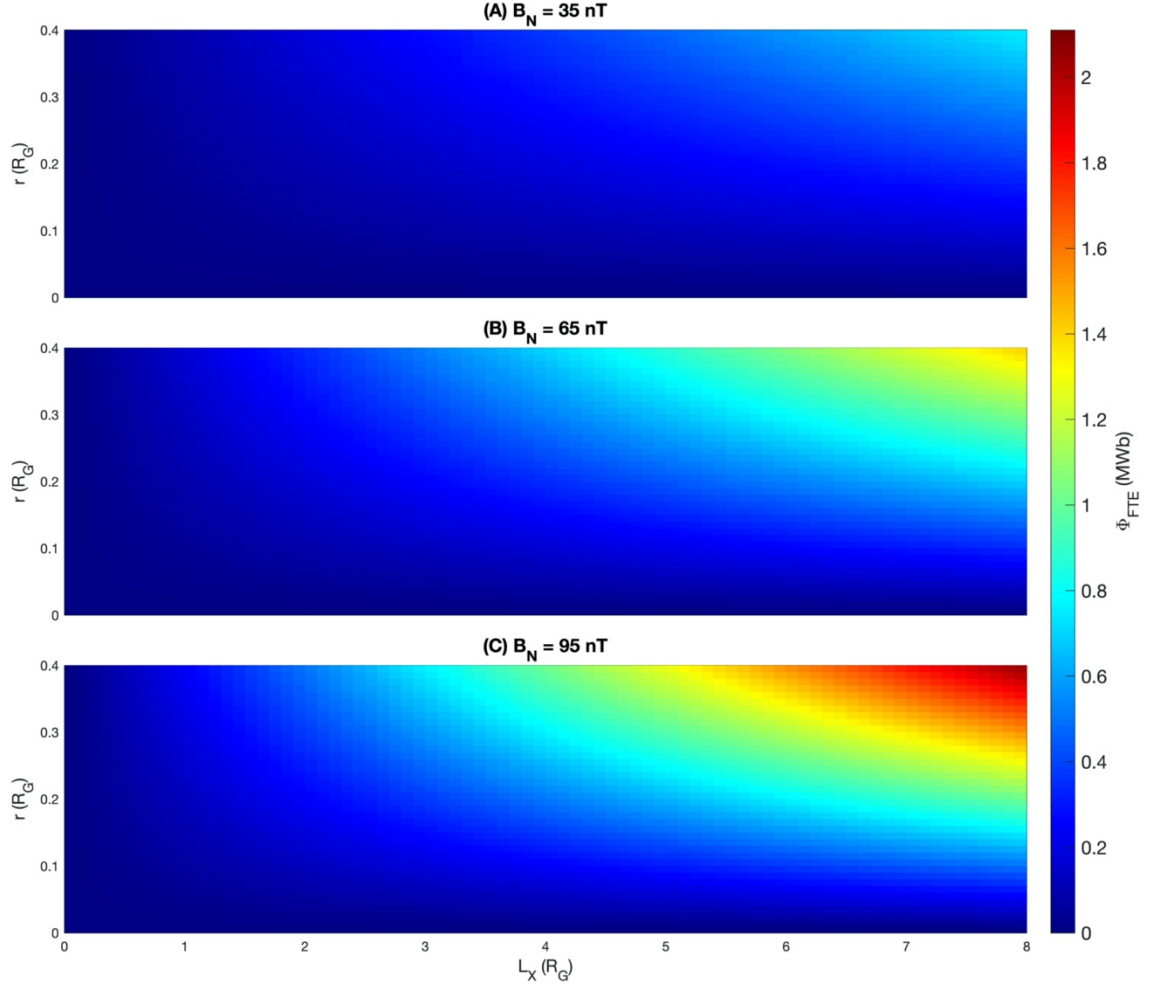


Figure 32: Examples of magnetic flux opened by FTEs calculated from Eq. 4.5 for (a) $B_N = 35$ nT, (b) $B_N = 65$ nT, and (c) $B_N = 95$ nT across ranges of independent FTE radius (r) and reconnection X-line length (L_x).

The above analysis assumes that Ganymede lies at centre of the Jovian plasma sheet ($\lambda_{III} = 248^\circ$). As Ganymede transitions to higher/lower magnetic latitudes, the guide field becomes more influential on magnetopause plasma-magnetic interactions. The presence of B_Y facilitates larger FTE growths, often by joining multiple small FTE magnetic ‘islands’ into one larger flux rope/tube structure (Jia, personal communication, July 16, 2021). The guide field also enhances the magnetic field component axial to the flux rope, resulting in larger Φ_{tube} (Eq. 4.4) which likely leads to larger total FTE flux content Φ_{open} . Consequently, the FTE flux contributions derived from $\lambda_{III} = 248^\circ$ plasma-magnetic conditions may not represent true upper limits, and localised reconnection may become even more influential to magnetic flux opening when Ganymede is away from the Jovian plasma sheet’s centre.

My work should be considered as a general gauge, given a number of simplifications used in this analysis. First, the magnetopause-normal magnetic field component is not uniform inside

an FTE, and the usage of peak B_N leads to a likely overestimation of FTE flux content. Similarly, Eq. 4.5 assumes long FTEs generated by X-lines when FTEs can also arise as small elbow-shaped flux tubes, which also leads to an overestimation. Fortunately, the conclusions reached on the FTEs' influence only require the upper limit results, hence are not impacted by overestimation.

Any isolated bimodal oscillation in B_N found in MHD-EPIC magnetopause snapshots are assumed to be FTEs. However, numerous oscillations are also observed in close proximities, and it can be uncertain whether they represent distinct FTEs, or they are other fluctuations not caused by localised magnetic reconnection. Although the 71 FTE timesteps are verified with Zhou et al., (2020), the final FTE count is not guaranteed to be truly accurate. Nevertheless, the sample size analysed is sufficiently large to determine B_N distributions. Moreover, unaccounted FTEs are much more likely to be small and faint, and therefore unimpactful to the main conclusions drawn.

Ultimately, much remains to be answered on the magnetic reconnection structure on Ganymede's upstream magnetopause. Magnetic flux analysis of the MHD-EPIC simulation suggests important contributions from localised FTEs, but there is still significant amount of unaccounted flux opened on the surface that is strongly reconnection-favourable. I compare flux contents of FTEs to the open flux in the polar caps, which thus far are assumed as stationary flat surfaces bounded by OCFBs at constant co-latitudes. In the next section, I will consider the upstream flux opening on the OCFBs, and how their oscillations can constrain Ganymede's Dungey-like plasma convection cycle.

4.3. Oscillation of Ganymede's OCFBs and the Effects of Magnetic Flux Closure

Irrespective of its global structure, magnetic reconnection on Ganymede's upstream magnetopause is driven by Jupiter's rotation. Hence, as the moon traverses inside the Jovian plasma sheet, reconnection-associated electric field and voltage on average exhibits a Jovian-diurnal variation. Although real time-plots of E_{up}/V_{up} will possess significant fluctuations from FTEs, for this analysis I shall assume smooth E_{up}/V_{up} plots associated with one reconnection X-line across the equatorial magnetopause. Such global structure is permitted at all times

given results from the reconnection onset assessment (Figure 24), and may contribute to the fraction of Φ_{open} not accounted by Φ_{FTE} .

If the reconnection electric field is assumed parallel to the X-line, the upstream voltage is calculable by the analytical model (Chapter 3, Section 3.2). I first obtain an array of distances between each grid point on the $Z = 0$ magnetopause, then multiply each ‘step length’ by E_{up} value at the start of each step to form a voltage array, which is then summed up to find V_{up} value. This process is repeated across all Jovian System-III east longitude (λ_{III}), resulting in a periodic variation shown through the blue line in Figure 33 where Ganymede starts at $\lambda_{\text{III}} = 248^\circ$ when $t = 0$. The X-line voltage has minimums $V_{\text{up}} \sim 60$ kV when Ganymede is at the Jovian plasma sheet’s centre ($\lambda_{\text{III}} = 68^\circ, 248^\circ$) and maximums $V_{\text{up}} \sim 105$ kV when Ganymede is at its highest and lowest points inside the plasma sheet ($\lambda_{\text{III}} = 158^\circ, 338^\circ$). The minimum voltage is $\sim 70\%$ of V_{up} computed under comparable magnetopause conditions by Zhou et al., (2020), which is a sensible estimate for non-FTE magnetic flux contribution¹⁸.

Adapting the analytical model for intermediate Ganymedeian orbital position requires more granular parametrisations of upstream Jovian plasma-magnetic conditions than previously described. Instead of defining individual inputs for each of the three extreme λ_{III} values, upstream parameters are now prescribed as functions

$$\rho_{J,0} = (\rho_{J,0})_{\lambda_{\text{III}}=248^\circ} \exp\left(\frac{-d}{H}\right)^2 \quad (4.10)$$

$$P_{J,0} = (P_{J,0})_{\lambda_{\text{III}}=248^\circ} \exp\left(\frac{-d}{H}\right)^2 \quad (4.11)$$

$$B_{J,0} = |35 \sin(\lambda_{\text{III}} - 248^\circ)| + (B_{J,0})_{\lambda_{\text{III}}=248^\circ} \quad (4.12)$$

¹⁸ Again, if the X-line is to generate FTEs, then part of its flux would be included in Φ_{FTE} .

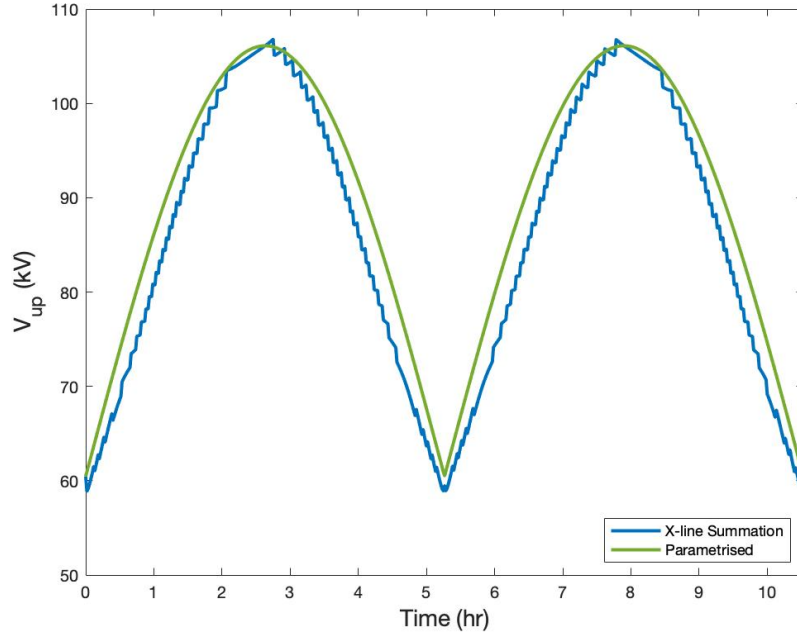


Figure 33: Temporal variation of upstream magnetic reconnection voltage produced by a magnetopause-wide equatorial X-line as Ganymede traverses the Jovian plasma sheet, with the reconnection electric field assumed parallel to the X-line. Blue line indicates voltage if the magnetopause analytical model (Chapter 3, Section 2) is iterated over one synodic period at $\Delta\lambda = 1^\circ$ (105.3 s) interval. Green line indicates voltage curve parametrised by Eq. 4.14.

where d is Ganymede's distance from the centre of the Jovian plasma sheet generalised from Eq. 2.6

$$d = (15 R_J) \sin(6.3^\circ) \sin(\lambda_{III} - 248^\circ) \quad (4.13)$$

and $H = \frac{1.62}{\sqrt{\ln 2}}$ is chosen as the plasma sheet's scale height as it yields extreme plasma mass density and pressure consistent with values used during magnetic reconnection assessment. Eq. 4.10 – 4.13 replace constant upstream values in the model's MATLAB code, reducing its inputs to upstream Jovian plasma flow speed and east longitude. The magnetopause domain used in the model remains unchanged.

Given that the reconnection electric field is most sensitive to upstream Jovian magnetic field strength, whose relationship with the Jovian east longitude (and thus time) follows a sine relationship (Jia et al., 2008), the reconnection voltage can be fairly well parametrised by the equation

$$V_{\text{up}}(t) = (V_{\text{max}} - V_{\text{min}}) \left| \sin \left(\frac{2\pi t}{10.53 \times 3600} \right) \right| + V_{\text{min}} \quad (4.14)$$

where $V_{\text{max}} = 105$ kV, $V_{\text{min}} = 60$, and the unit of time t is in seconds. Parametrised upstream reconnection voltage is shown as the green line in Figure 33. Although both voltage lines have the same period, parametrised voltage oscillations have greater widths than those produced by numerical X-line integration.

To produce Figure 33, I kept Ganymede's northern and southern OCFBs fixed at $\theta_0 = 55^\circ$ and $\theta_0 = 125^\circ$ respectively throughout the moon's trip through the Jovian plasma sheet. Although this choice does not capture OCFB motions arising from magnetic flux opening, it does not impact reconnection voltage calculated along Ganymede's equatorial magnetopause. In reality, opening magnetic flux will push the OCFBs toward Ganymede's equator. How influential is upstream reconnection on the moon's magnetospheric flux content? One way to answer this question is to consider an unrealistic scenario where none of the flux opened is closed i.e., if downstream reconnection does not exist. In this case, how long does it take for every closed magnetic field line to be opened?

Using Eq. 4.9, I compute absolute amounts of open magnetic flux associated with different northern OCFB co-latitudes in Figure 34. As the boundary moves from the pole ($\theta_0 = 0^\circ$) to the equator ($\theta_0 = 90^\circ$), $|\Phi_{\text{open}}|$ increases from zero to a maximum of 15.7 MWb. When open flux is added or removed from Ganymede's polar cap, Eq. 4.9 allows the changes to be translated into OCFB movements.

My next step is to evaluate the OCFBs' temporal evolution under upstream reconnection voltage from Figure 33. Without flux closure, the two OCFBs will eventually meet at

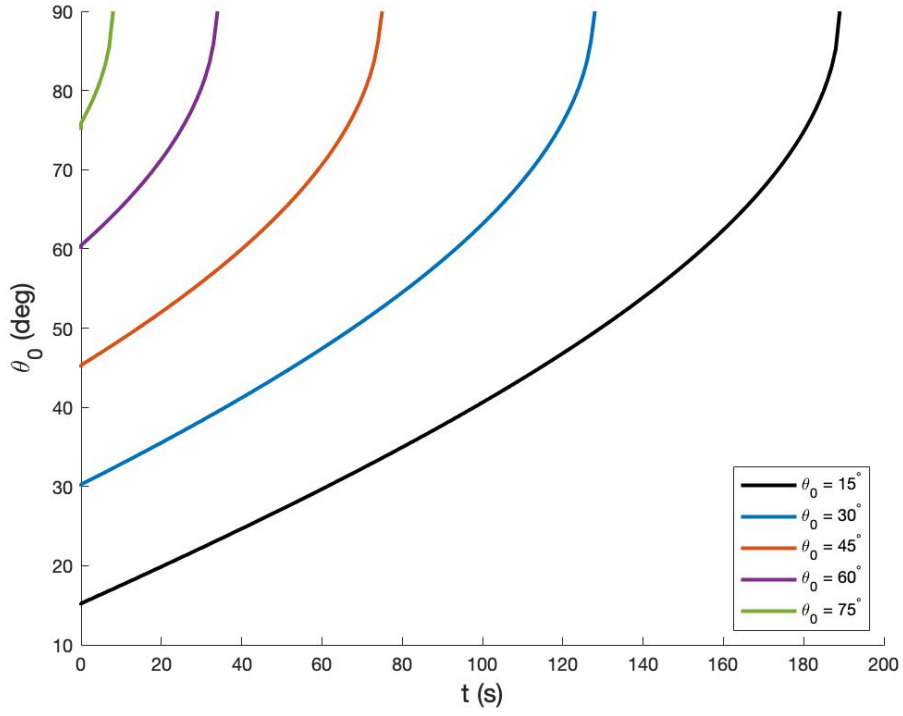


Figure 34: Time required for an idealised Ganymede northern OCFB to reach the moon's equator if there is no magnetic flux closure in the magnetotail. Different lines correspond to different starting OCFB co-latitudes.

Ganymede's equator after a certain time interval. In Figure 34, I consider the northern OCFB at different starting co-latitudes ($\theta_0 = 0^\circ, 15^\circ, 30^\circ, 45^\circ, 60^\circ, 75^\circ$) and track its journey toward the equator at $\Delta t = 1$ second resolution. Irrespective of its starting location, the OCFB is expected to reach Ganymede's equator in no more than $\sim 1 - 4$ minutes. The flux opening occurs so rapidly that the upstream reconnection voltage does not deviate from its initial minimum value.

Given that the HST does not detect auroras anywhere near Ganymede's equator, there must be rapid magnetic flux closure in the downstream magnetotail. Consider a simple scenario where magnetic flux is opened and closed at one upstream and one downstream reconnection structures with no other flux sources or sinks. Assuming also that the flux transport occurs in one direction i.e., all opened flux moves uniformly downstream, then the downstream reconnection voltage can be expressed by

$$V_{dn}(t) = V_{up}(t + \Delta t_{oc}) \quad (4.15)$$

where Δt_{oc} is time delay between flux opening and closure. This flux transfer time can be roughly estimated by assuming that upstream and downstream reconnection sites are equidistance from Ganymede, and that magnetic field lines travel at bulk speed of the upstream Jovian plasma ($v_{J,0}$). Hence,

$$\Delta t_{oc} \sim \frac{2r_{MP}}{v_{J,0}} \quad (4.15)$$

where r_{MP} is the magnetopause standoff distance. For simplicity, I use $r_{MP} = 1.65 R_G$ for the analytical model and $v_{J,0} = 140$ km/s, resulting in $\Delta t_{oc} \sim 62$ s. Note that using the observational estimate $r_{MP} \sim 2 R_G$ will result in a slightly longer $\Delta t_{oc} \sim 75$ s. Compared to the times in Figure 34, magnetic flux transfer time is sufficiently short so Ganymede's closed-field region will not be depleted. Note that Δt_{oc} can be translated into a 'phase delay' in System-III Jovian east longitude where

$$\frac{\Delta \lambda_{III}}{360^\circ} = \frac{\Delta t_{oc}}{10.53 \text{ hr}} \quad (4.16)$$

where 10.53 hours is the Jovian synodic period in Ganymede's rest frame. This substitution is for convenience since the analytical model defines upstream plasma-magnetic conditions by the east longitude (Eq. 4.10 – 4.12). Thus, the interval $\Delta t_{oc} \sim 62$ s translates to $\Delta \lambda_{III} \sim 0.5^\circ$, which also illustrates the rapid nature of Ganymede's magnetic flux transport.

For any time interval Δt , a difference between V_{up} and V_{dn} will produce a change in opened magnetic flux $\Delta \Phi_{open}$. This is shown *schematically* in Figure 35, where differences between curves are exaggerated since the real V_{up} and V_{dn} time-plots are virtually identical. For the northern OCFB, $\Delta \Phi_{open}$ is negative

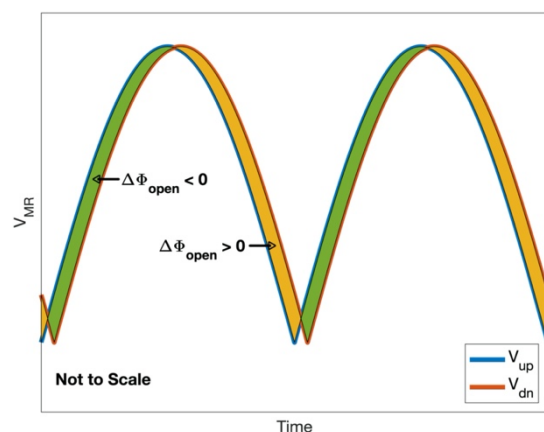


Figure 35: An illustration of net magnetic flux opening (green) and closure (yellow) arising from difference between upstream and downstream magnetic reconnection voltages. Diagram is not to scale.

(green) when $V_{\text{up}} > V_{\text{dn}}$ and positive (yellow) when $V_{\text{up}} < V_{\text{dn}}$. Once the two voltages are obtained, I compute the potential difference $V_{\text{diff}} = V_{\text{up}} - V_{\text{dn}}$ over one Jovian synodic period and integrate the area under the V_{diff} curve $\Delta t = 1$ min resolution (lengthened from previous 1-second intervals for computational efficiency) using the trapezium estimation. After each model iteration, $\Delta\Phi_{\text{open}}$ is added to the total open flux, and subsequently the change to the OCFB oval co-latitude is calculated via Eq. 4.9.

Starting from $\theta_0 = 55^\circ$, Ganymede's northern OCFB motion tracking is illustrated in Figure 36a for both numerically summed (blue) and parameterised (green) upstream reconnection voltages. In both cases, the OCFB oscillates with the period of 5.27 hours, or half the Jovian synodic period. This means the boundary does not produce a net shift over large timescales, which is consistent with long-term auroral stability observed by the HST. Numerically summed and parametrised reconnection voltages produce similar OCFB motions when the boundary is far from Ganymede's equator. However, the discrepancy spiked near co-latitude maximums $\theta_0 \approx 87^\circ$ and equatorial OCFB is predicted for the X-line summed voltage. This is a result of integration error, in which step-like natures of summed V_{up} (and thus V_{dn}) produces numerous irregular spikes in V_{diff} . Poor function behaviour means that trapezium integration is no longer an accurate approximation. In contrast, parametrised V_{up} and V_{dn} produce smooth and well-bounded V_{diff} which generates less extreme OCFB behaviour.

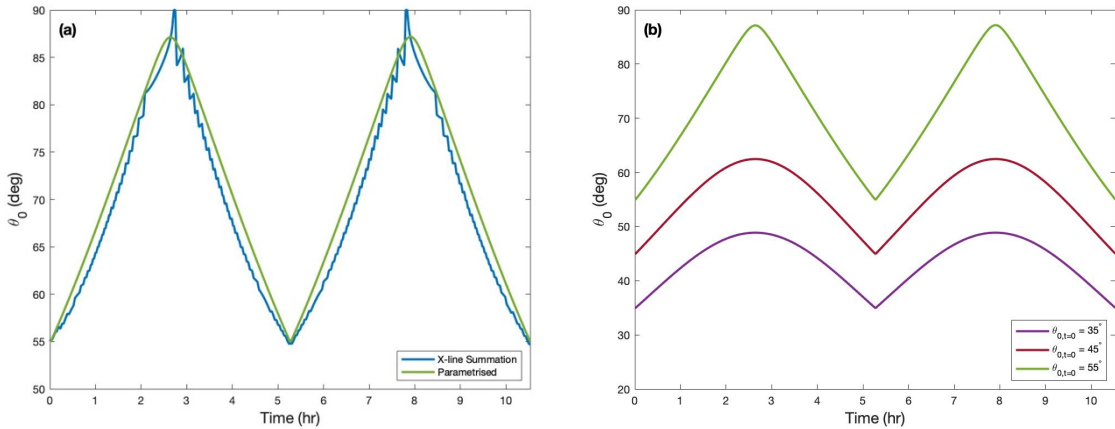


Figure 36: (a) Ganymede's northern OCFB oscillation for time-delayed equivalent upstream/downstream reconnection voltages, which are generated by manual X-line summation and sine modulus parametrisation. The boundary starts at $\theta_0 = 55^\circ$ when $t = 0$. (b) OCFB oscillations for different initial boundary oval co-latitudes under parameterised reconnection voltages.

Figure 36b exhibits oscillations of Ganymede's northern OCFB for three initial co-latitudes $\theta_0 = 35^\circ, 45^\circ, 55^\circ$. The oscillation amplitudes vary between $10^\circ - 30^\circ$ and increasing with larger initial co-latitude. This is due to the sine relationship between θ_0 and Φ_{open} described in Eq. 4.9, where oval movements require less flux change as θ_0 approaches 90° . Oscillations of OCFBs have been noticed in resistive MHD simulation (Jia et al., 2009), which suggests that the idea of equivalent downstream flux closure likely has some merit.

However, the significant oscillation amplitude deserves some attention. Even well-behaved parametrised reconnection voltages are not immune to integration error, which may potentially overestimate magnetic flux change. More fundamentally, each OCFB oscillation is accumulated through four long intervals of small but consistent voltage difference (Figure 35), which is only possible for smooth X-line voltage curves. Under real conditions, magnetic flux transport will be complicated by non-idealised factors including transient FTEs, non-equivalent upstream and downstream voltages, and presences of multiple flux sources and sinks operating at different timescales. These factors may create more intersection points between V_{up} and V_{dn} , resulting in smaller disjointed intervals of up-and-down OCFB movements that generate little net motion. At the same time, significant voltage fluctuation can generate irregular spikes in V_{diff} , which can increase the oscillation amplitude as seen in Figure 36a. Ultimately, the translation of my Figure 36 findings to physical results remains unclear.

It is also uncertain how OCFB oscillation will manifest once upstream-downstream asymmetry is taken into account. Tilted OCFB may simply undergo vertical shifts with unchanged asymmetry, or complications may arise given that the asymmetry is produced by upstream Jupiter-Ganymede plasma interactions. Finally, while the OCFBs may be strongly linked with aurora ovals, it is not certain if their correlation will be impacted by the OCFBs' motions. If the boundaries' movement do not translate fully to the auroras', then my predictions cannot be verified until the next spacecraft arrives at Ganymede.

Ultimately, the detailed picture of Ganymede's plasma convection remains murky despite my findings. The analytical model used in this section is limited by its simplicity, even though it has been expanded to capture temporal upstream plasma-magnetic condition changes. The

inroads needed toward understanding of downstream reconnection, and general magnetic flux transport, are unlikely to be achieved without significant additional in-situ observations. With this section and chapter, my research on Ganymede's magnetic reconnection has been fully discussed. However, information unveiled by the analytical model has not been exhausted. Next, I return to Ganymede's upstream magnetopause and consider the other notable energy-transfer interaction – Kelvin-Helmholtz instability.

Chapter 5: Kelvin-Helmholtz Instability on Ganymede's Upstream Magnetopause

In this chapter, I cover the path of Ganymede's K-H instability from onset to nonlinear vortex formation. The interaction's picture remains largely murky, but the process has revealed multiple useful insights. Linear K-H instability onset inequality is expected to be satisfied on Ganymede's magnetopause flanks, forming waves that propagate downstream alongside the external Jovian plasma flow. Then, the instability is assessed with respect to the finite Larmor radius (FLR) effect, in which kinetic effects from heavy Jovian ions produce asymmetric growth favouring the sub-Jovian magnetopause flank. Finally, formation of nonlinear K-H vortices is considered first in isolation, and later with magnetic reconnection effects taken into account – culminating in the finding that K-H instability is unlikely to drive Ganymede's Duney cycle.

5.1. From Perturbation to Vortices

Kelvin-Helmholtz (K-H) instability is a phenomenon of two phases – linear waves and nonlinear vortices. As discussed in Chapter 1, the nonlinear phase can facilitate plasma and energy transport across a boundary. In this section, I review K-H instability on a magnetopause from initial perturbation to nonlinear vortices.

Recall the magnetopause configuration in Figure 19 of Chapter 3, Section 3.1, in which plasmas are considered collisionless fluids under the MHD framework. A perturbation to the current sheet boundary generates a wavevector ($\mathbf{k}$) which can initiate K-H instability growth provided if (Farrugia et al., 1998)

$$[\mathbf{k} \cdot (\mathbf{U}_{\text{ex}} - \mathbf{U}_{\text{in}})]^2 > \frac{1}{\mu_0} \left(\frac{1}{\rho_{\text{ex}}} + \frac{1}{\rho_{\text{in}}} \right) [(\mathbf{k} \cdot \mathbf{B}_{\text{ex}})^2 + (\mathbf{k} \cdot \mathbf{B}_{\text{in}})^2] \quad (5.1)$$

where $\mathbf{U}_{\text{ex,in}}$ are bulk plasma flow velocities, $\rho_{\text{ex,in}}$ are plasma mass densities, and $\mathbf{B}_{\text{ex,in}}$ are magnetic field vectors external and internal to the magnetopause respectively. The

permeability of free space is $\mu_0 = 4\pi \times 10^{-7}$ H/m. Eq. 5.1 is known as the thin model¹⁹ and is commonly used to predict K-H instability (e.g., Masters, 2017).

The onset inequality reflects the limit beyond which local magnetic tension forces cannot stabilise a perturbed magnetopause. K-H instability becomes more favoured under faster bulk flow shear, smaller plasma mass density, and weaker boundary-adjacent magnetic fields. Orientation of the K-H wavevector also has material impact, destabilising the magnetopause when it is more parallel to $\mathbf{U}_{sh} = \mathbf{U}_{ex} - \mathbf{U}_{in}$ and/or more orthogonal to $\mathbf{B}_{ex,in}$. Note that the thin model is an approximation whose imperfections have been discussed (Gratton et al., 2004). Nevertheless, the model is generally accurate for cases where K-H wavevectors are strongly field-orthogonal ($\mathbf{k} \cdot \mathbf{B} \approx 0$, also known as flute mode) and the magnetic shear angle (θ_{sh}) is small.

Once formed, linear K-H instability waves propagate along the magnetopause. The wave transmission can be characterised by the phase velocity, which is equal to the centre-of-mass (zero-momentum frame) velocity of plasma near the unstable boundary

$$\mathbf{v}_{KH} = \frac{\rho_{ex}\mathbf{U}_{ex} + \rho_{in}\mathbf{U}_{in}}{\rho_{ex} + \rho_{in}} \quad (5.2)$$

where mass parameters have been replaced by densities given that I consider a common volume element straddling the magnetopause. The direction of propagation is therefore along the vector sum of external and internal bulk flows.

As a K-H wave propagates, its amplitude grows exponentially following the expression

$$A(t) = A_0 e^{\gamma t} \quad (5.3)$$

where A_0 is the initial perturbation amplitude and γ is the linear K-H instability growth rate, which can be obtained from the dispersion relation of the MHD momentum equation. After

¹⁹ In this model, the magnetopause is much thinner than the K-H wavelength $\lambda_{KH} = 2\pi/k$.

some time, the amplitude will become large such that the magnetopause enters the nonlinear turbulent vortex phase. The threshold magnitude likely depends on the magnetospheric system, but for simplicity I can define constant $A(t_f)$, where t_f is the time interval between K-H wave and K-H vortex formations. If the K-H wave propagates at constant phase velocity, then propagation distance follows

$$d_{KH} = |\mathbf{v}_{KH}|t_f \quad (5.4)$$

In reality, there is no analytical transition between linear and nonlinear K-H instability. Waves and vortices essentially exist as separate frameworks. Given the mathematical complexity of the nonlinear theory, studies of K-H vortices are generally conducted through numerical simulations, which directly compute equations of motions, and in-situ observations, which search for trademark sawtooth magnetic fluctuation signature. Early simulation results suggest that K-H instability strongly favour magnetopause conditions where adjacent magnetic fields are exactly parallel or antiparallel due to stabilising effects of guide field components (Thomas & Winske, 1993; Miura, 1995; Eastwood et al., 2015). Hence, linear and nonlinear K-H events are generally expected along low-latitude magnetopause flanks. This prediction is also consistent with the thin model, given the field-orthogonality of plasma flows in the flank regions.

Before furthering the discussion on magnetopause K-H instability, I must clarify the naming convention of its two flanks. For planetary magnetopauses, the solar wind plasma is emitted radially from the Sun, so the planet orbits toward the ‘dawn’ flank and away from the ‘dusk’ flank at all times. This is not the case for Ganymede, however. Since Jupiter’s plasma rotates with the planet, Ganymede orbits toward the magnetotail rather than one of the flanks. Hence, the Ganymedean magnetopause flanks are defined as ‘sub-Jovian’ and ‘anti-Jovian’, where the sub-Jovian flank lies between Ganymede and Jupiter at all times.

Linear and nonlinear K-H instability events have been observed on Earth’s magnetopause flanks, including evidence of cross-magnetopause particle mixing in K-H vortices as shown in Figure 37 (Hasegawa, Fujimoto, Phan, et al., 2004; Owen et al., 2004). Most instability

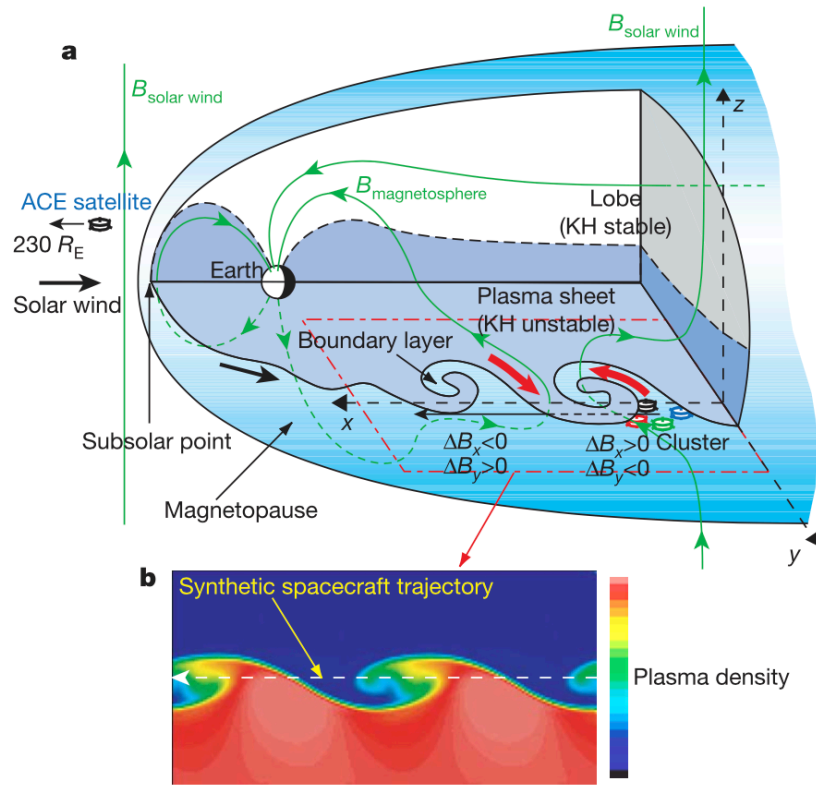


Figure 37: Illustrations of (a) K-H instability along Earth's magnetopause flank under northward external IMF, (b) trajectory of Cluster spacecraft through simulated nonlinear K-H vortex structure. Image courtesy of Hasegawa, Fujimoto, Phan, et al., (2004), *Nature*.

observations occur under parallel adjacent magnetic field configuration i.e., northward external IMF, with only intermittent evidence under antiparallel configuration (Hwang et al., 2011). Since K-H instability theoretically do not differentiate between the two alignments, the asymmetry is generated from external factors. Most commonly, magnetic reconnection has been shown to inhibit growth of K-H vortices under southward IMF (Nakamura et al., 2020).

Further away, K-H instability has also been observed on Saturn's magnetopause flanks in both linear and nonlinear phases (Masters et al., 2009; Masters et al., 2010; Wilson et al., 2012; Delamere et al., 2013). No preference between parallel and antiparallel magnetic configurations has been identified, but the instability appears to favour the dusk flank due to rapid planetary rotation dynamics (Masters et al., 2010; Delamere et al., 2013). As of this writing, no K-H instability events have been directly observed on Jupiter's magnetopause flanks, although their presence are expected from simulations with Saturn-like preference for the dusk flank (Desroche et al., 2012; Masters, 2017; Zhang et al., 2018).

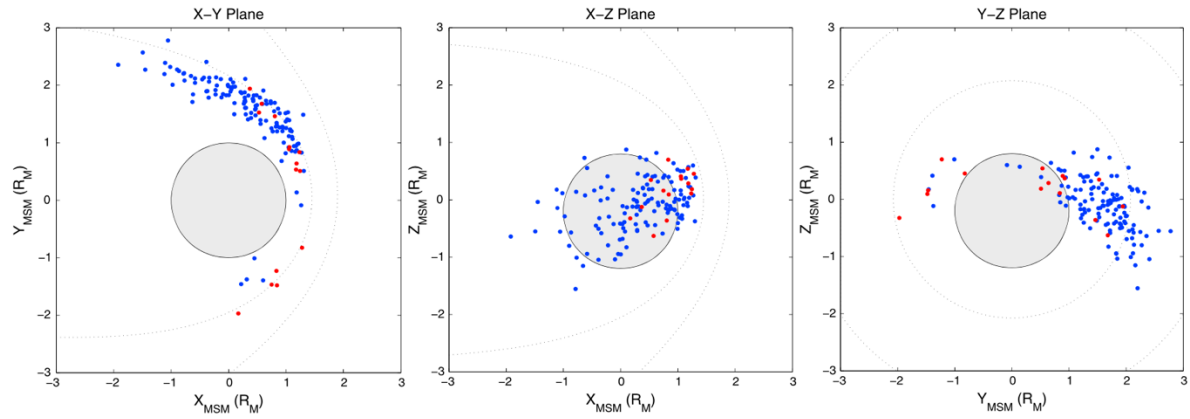


Figure 38: Location of linear K-H waves (blue) and nonlinear K-H vortices (red) observed by MESSENGER at Mercury. Inner and outer dashed lines show the magnetopause and the bow shock respectively. The incidents' distribution show K-H instability's clear preference for dusk magnetopause flanks ($Y_{MSM} > 0$). Image courtesy of Liljeblad et al., (2014), *Journal of Geophysical Research: Space Physics*.

Closer to the Sun, the MESSENGER spacecraft has captured K-H instability events on Mercury's magnetopause flanks (Slavin et al., 2008; Boardsen et al., 2010; Sundberg et al., 2012; Liljeblad et al., 2014). Among 146 encounters, 93% occur in the dusk magnetopause flank (Figure 38) despite Mercury not possessing fast rotations like Jupiter and Saturn (Liljeblad et al., 2014). The inter-flank asymmetry is indicative of kinetic influence on K-H instability, specifically the finite Larmor radius (FLR) effect, which occurs when local plasma ion gyroradii are non-negligible compared to the magnetopause thickness (Nagano, 1978, 1979; Huba, 1996). Given the size similarity between Mercury and Ganymede ($R_M = 2,440$ km and $R_G = 2,634$ km), it is reasonable to expect an FLR effect on K-H instability growth along Ganymede's magnetopause flanks.

In this thesis, I describe the FLR effect as a superposition of vorticities arising from bulk and ion plasma flow shears. This idea is illustrated schematically in Figure 39a for Ganymede's magnetopause, and Figure 39b for Mercury's magnetopause for comparison. For both systems, I consider the case where an infinitesimally thin magnetopause is north-south symmetric with adjacent magnetic fields that are parallel/antiparallel to each other and to the body's rotation axis. This simplification is more representative for Ganymede, as the Jovian magnetic field never deviates beyond 45° from the moon's rotation axis (Eq. 3.24 – 3.25) while the solar IMF rotates fully across the clock angle spectrum near Mercury.

In each sub-figure, the external plasma forms two flows of equal velocity ($\mathbf{U}_J$, $\mathbf{U}_{MSH}$) along both equatorial magnetopause flanks at respective upstream-downstream boundary points. These

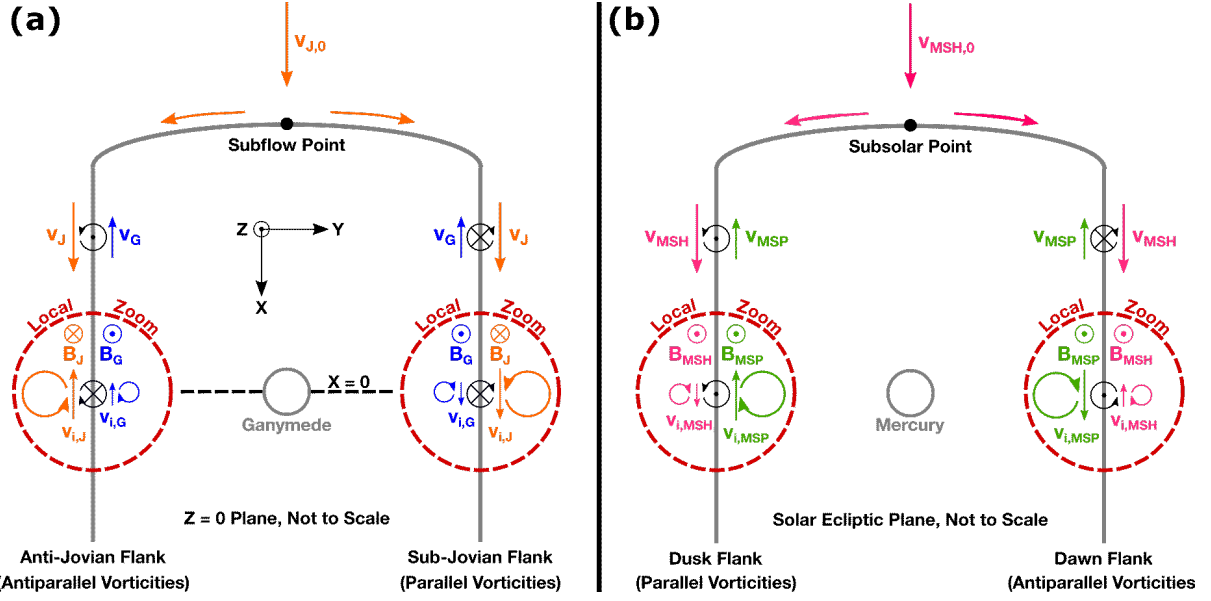


Figure 39: Schematic diagram of the FLR effect along equatorial magnetopause flanks of (a) Ganymede and (b) Mercury. Zoom windows indicate local ion motions at the *same* locations as bulk flow shears. K-H instability growth is enhanced (reduced) when bulk and ion flow shear vorticities are parallel (antiparallel). The figure is not to scale. Image courtesy of Kaweeyanun et al., (2021), *Journal of Geophysical Research: Space Physics*.

flows generate bulk shear against internal flows ($\mathbf{U}_G$, $\mathbf{U}_{MSP}$) which are assumed to be flank symmetric. Then, I compute bulk flow shear vorticity

$$\boldsymbol{\omega}_b = \nabla \times \mathbf{U}_{b,sh} \quad (5.5)$$

where $\mathbf{U}_{b,sh} = \mathbf{U}_{ex} - \mathbf{U}_{in}$ is the bulk flow shear whose magnitude is flank symmetric. The bulk vorticity points southward (into page) on the sub-Jovian/dawn flanks, and northward (out of page) on the anti-Jovian/dusk flanks. Next, I zoom in locally at equatorial magnetopause flank points to consider ion gyromotion under boundary-adjacent magnetic fields. Gyration (clockwise or counterclockwise) occurs under the left-hand thumb rule²⁰ with gyroradius

$$R = \frac{M_i v_i}{q_i B} = \frac{M_i}{q_i B} \sqrt{\frac{2k_B T_i}{M_i}} \quad (5.6)$$

where $\frac{M_i}{q_i}$ is the plasma's average ion mass to charge ratio, T_i is the plasma's average ion temperature, $k_B = 1.38 \times 10^{-23} \text{ JK}^{-1}$ is the Boltzmann's constant, and B is the magnetic field

²⁰ Make a thumb-up figure with the left hand. If the magnetic field is parallel to the thumb, positive ions gyrate in direction that the remaining fingers curl into the palm.

strength. The square root in Eq. 5.6 can be reduced to v_i – the local ion flow speed tangential to gyromotion (i.e., gyration speed). Subsequently, I define the gyrofrequency

$$\Omega = \frac{q_i B}{M_i} \quad (5.7)$$

and the gyro-viscous coefficient

$$\eta = \frac{R^2 \Omega}{4} \quad (5.8)$$

which represents a measure of flux momentum arising from gradient distribution of ion guiding centres (Schnack, 2003). Gyromotion of ions generate local external and internal flow vectors $\mathbf{v}_{i,ex}$ and $\mathbf{v}_{i,in}$ which together produce non-negligible ion flow shear $\mathbf{v}_{i,sh}$ with vorticity

$$\boldsymbol{\omega}_i = \nabla \times \mathbf{v}_{i,sh} \quad (5.9)$$

The ion vorticity will be parallel or antiparallel to the bulk vorticity (Eq. 5.5). Hence, the magnitude of total plasma vorticity can be expressed as the sum or difference of bulk and ion vorticities i.e.,

$$|\boldsymbol{\omega}| = |\boldsymbol{\omega}_b| \pm |\boldsymbol{\omega}_i| \quad (5.10)$$

where the flank with vorticity addition is expected to possess greater K-H instability growth assuming a proportional relation between growth rate and plasma shear. At Ganymede, bulk flow shears on sub-Jovian (anti-Jovian) flank produces southward (northward) vorticities. The ion flow shears, driven by large external gyration, produce vorticities that are southward on both magnetopause flanks. This leads to vorticity summation on the sub-Jovian flank and subtraction on the anti-Jovian flank, so the former should receive enhanced K-H instability growth. At Mercury, bulk flow shears on dawn (dusk) flank produces southward (northward) vorticities – a configuration equivalent to Ganymede. But unlike the moon, Hermean ion flow shears produce northward vorticities on both flanks as they are driven by large internal

gyration. Consequently, bulk and ion vorticities are summed on the Hermean dusk flank, consistent with enhanced K-H instability growth observed by MESSENGER.

Figure 39 does not offer a robust way to quantify the FLR effect. One simpler method is to relate K-H instability inter-flank asymmetry proportionally to the magnitude of the ion flow shear. If $v_{i,sh}$ increases, so does the magnitude of growth asymmetry. Seeking a more robust solution, I may consider methods used in existing literature. The FLR effect has been quantified by adding a gyro-viscous tensor (proportional to the gyro-viscous coefficient in Eq. 5.8) to the MHD pressure term (Nagano, 1978, 1979; Glassmeier & Espley, 2006; Sundberg et al., 2010) or alternatively through kinetic numerical simulations that bypass the fluid plasma description (Nakamura et al., 2010; Paral & Rankin, 2013). Unfortunately, the two methods produce a significant discrepancy when they were applied to Mercury. The MHD-FLR formulation predicts K-H instability growth enhancement on the Hermean dawn flank, contradicting the kinetic simulations that favour the dusk flank. The mathematical explanation behind the paradox is likely rooted in the energy equation, which takes a simpler form in the MHD-FLR theory (Umeda et al., 2016).

Furthermore, the two methods explain the FLR effect differently. Under the modified MHD-FLR formulation, K-H instability growth asymmetry occurs in its linear phase, arising from relative direction of linear K-H wave propagation and local diamagnetic drift, the latter of which orients oppositely between magnetopause flanks (Huba, 1996). The asymmetry manifests in the linear growth rate (γ) and amplifies as the instability enters its nonlinear vortex phase. In contrast, the kinetic simulations show that linear K-H wave growth should be flank symmetric because magnetopause current sheets expand rapidly beyond the kinetic length scale (Nakamura et al., 2010). Instead, difference in ion centrifugal drift in response

the FLR effect introduces growth asymmetry during the linear-to-nonlinear phase transition, caused by differing ion centrifugal drift response to the global convective electric field.

One might think the MHD-FLR analytical theory could be applicable to my analysis despite predicting the wrong K-H enhanced magnetopause flank at Mercury. However, this is not the case for two reasons. The FLR effect is encoded as a first order perturbation to the MHD

pressure tensor. Higher order contributions are assumed to be small, but these terms in fact diverge for an infinitesimally thin magnetopause (Nagano, 1978, 1979). The MHD-FLR theory was also formulated for northward external IMF i.e., parallel cross-boundary magnetic field configuration (Nagano, 1978, 1979), which is not the case for Ganymede given the strong antiparallel configuration produced by always-southward Jovian magnetic field and 176° tilt of the Ganymedean dipole axis.

Nevertheless, useful insights can be taken from the MHD-FLR framework. The theory asserts that the magnitude of K-H instability growth asymmetry is proportional to (1) the bulk flow shear and (2) the sum of gyro-viscous coefficients external and internal to the magnetopause (same on both flanks). These relationships likely hold true even though the theory is inapplicable as a whole. Consequently, I can calculate the two parameters, alongside the ion flow shear magnitude, as ‘quantifiers’ of the FLR effect. These values can then be compared to those at well-observed Mercury’s magnetopause, and return some quantitative insights into asymmetric K-H instability growth at Ganymede.

The rest of this chapter divides the assessment of Ganymede’s K-H instability growth into two halves. In Section 5.2, the linear MHD K-H instability onset will be assessed across Ganymede’s magnetopause by the analytical model previously used for magnetic reconnection (Chapter 3). Then, in Section 5.3., the kinetic FLR effect will be evaluated semi-quantitatively on Ganymede’s magnetopause flanks. Their results will be reconvened in Section 5.4, where I shall evaluate potential formation of nonlinear energy-transferring K-H vortices on Ganymede’s upstream magnetopause.

5.2. Linear K-H Instability Onset on Ganymede’s Magnetopause

Given no previous knowledge of K-H instability on Ganymede’s upstream magnetopause, the K-H wavevector ($\mathbf{k}$) in the instability onset condition (Eq. 5.1) is an unconstrained parameter. Consequently, I first evaluate the inequality for the case when the K-H wavevector is parallel to the bulk flow shear ($\mathbf{k} \parallel \mathbf{U}_{b,sh}$), and then reassess the condition for every 1° rotation of the wavevector until a full circle is completed. Recall that the idealised magnetopause grid stretches between $-4.0 < Y < 4.0 R_G$ and $-1.0 < Z < 1.0 R_G$ with $0.01 R_G$ resolution in the

Cartesian GphiO coordinates. Near-magnetopause parameters are parametrised using the same specifications discussed in Chapter 3, Section 3.2²¹, and then substituted into Eq. 5.1 with $\mathbf{U}_{\text{ex}} = \mathbf{U}_J$, $\mathbf{U}_{\text{in}} = \mathbf{U}_G$, $\rho_{\text{ex}} = \rho_J$, $\rho_{\text{in}} = \rho_G$, $\mathbf{B}_{\text{ex}} = \mathbf{B}_J$, and $\mathbf{B}_{\text{in}} = \mathbf{B}_G$.

A grid point is considered to be ‘K-H unstable’ if (1) the onset inequality is satisfied for *one* of the K-H wavevector orientation and (2) the unstable point is surrounded by at least four other points that satisfy the inequality. The second requirement is imposed to eliminate ‘isolated unstable points’, where the inequality is satisfied by the smallest margin through only one specific wavevector orientation. While these points may be truly unstable, they cannot facilitate K-H instability growth as global structures hence their omission.

Where the magnetopause is K-H unstable, I calculate the corresponding K-H wave phase velocity following Eq. 5.2. Recall that the Ganymede-side bulk flow is assumed to be stagnant i.e., $U_G = 0$ km/s. Hence, the velocity expression is reduced to

$$\mathbf{v}_{\text{KH}} = \left(\frac{\rho_J}{\rho_J + \rho_G} \right) \mathbf{U}_J \quad (5.11)$$

which indicate that K-H wave phase velocity is parallel to the external Jovian-side bulk plasma flow.

Similar to reconnection, the linear K-H instability onset condition is assessed when Ganymede lies at the centre of the Jovian plasma sheet ($\lambda_{\text{III}} = 248^\circ$), and when the moon is at its highest and lowest in-sheet magnetic latitudes ($\lambda_{\text{III}} = 158^\circ, 338^\circ$). The results are shown in Figure 40 where coloured regions signify K-H unstable surfaces, and the colour scale indicates phase velocity of propagating K-H waves. In Figure 40a ($\lambda_{\text{III}} = 248^\circ$), the north-south symmetric magnetopause is almost entirely K-H unstable except the area immediately surrounding the subflow point. The phase velocity is highly variable between $|\mathbf{v}_{\text{KH}}| = 0 - 89$ km/s where faster speeds occur closer to the two magnetopause flanks, and propagation directions point toward the downstream as governed by the Jovian-side plasma flow. K-H instability onset is

²¹ Chronologically, the research in this chapter is conducted before those in Chapter 4. Hence, upstream Jovian plasma-magnetic conditions are inputted as constants.

significantly more constrained in Figure 40b ($\lambda_{\text{III}} = 158^\circ$) and Figure 40c ($\lambda_{\text{III}} = 338^\circ$). The maximally asymmetric magnetopause is only K-H unstable along narrow strips on the magnetopause flanks ($|Y| > \sim 2 R_G$). The K-H wave phase velocity correspondingly possesses narrower ranges of $|\mathbf{v}_{\text{KH}}| = 50 - 66 \text{ km/s}$ with increasing speeds toward the flanks.

I approximate $|\mathbf{v}_{\text{KH}}| \sim 0.65|\mathbf{U}_{\text{b,sh}}|$ inside K-H unstable regions when Ganymede lies at the centre of the Jovian plasma sheet, and $|\mathbf{v}_{\text{KH}}| \sim 0.48|\mathbf{U}_{\text{b,sh}}|$ when Ganymede transitions to highest and lowest magnetic latitudes. These results indicate that K-H wave phase velocity is

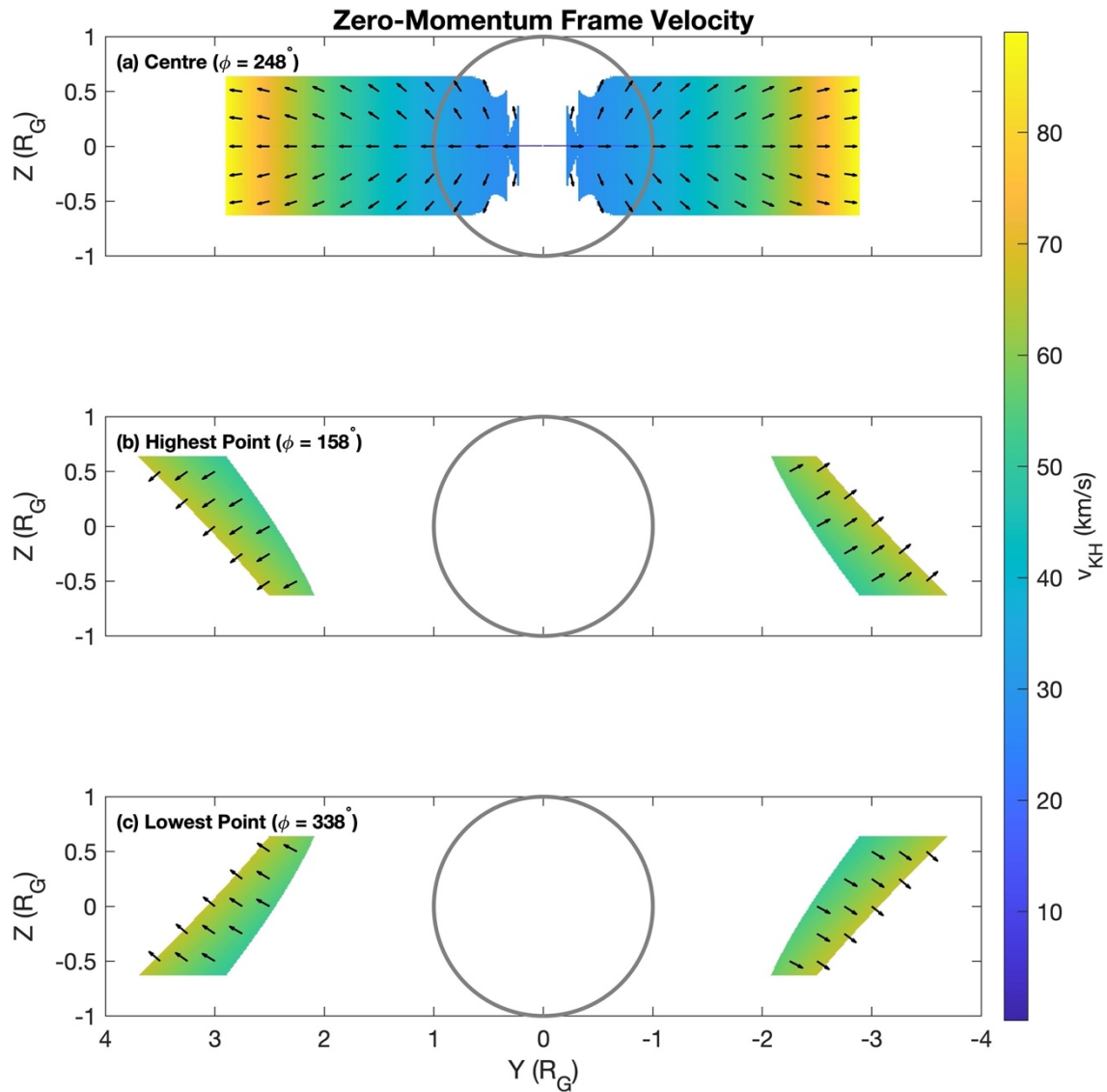


Figure 40: K-H instability onset assessment and computation of linear K-H wave phase velocity ($\mathbf{v}_{\text{KH}}$) when Ganymede is at (a) the centre of the Jovian plasma sheet, (b) highest in-sheet magnetic latitude, and (c) lowest in-sheet magnetic latitude. Ganymede is outlined in grey. Image courtesy of Kaweeyanun et al., (2021), *Journal of Geophysical Research: Space Physics*.

roughly half that of external Jovian plasma flow, which is compatible with in-situ observations (e.g., Kivelson et al., 1998).

The sizable reduction in K-H unstable areas as Ganymede moves away from the Jovian plasma sheet's centre is due to two reasons. Firstly, the Jovian plasma mass density decreases by 50% from $\lambda_{III} = 248^\circ$ to $\lambda_{III} = 158^\circ, 338^\circ$ while its magnetic field increases by 50% during the same transition. These upstream condition changes lower the threshold for K-H instability and thus the chance of interaction onset. Note that the threshold is proportional to $|\mathbf{B}_J^2|$ and $\frac{1}{\rho_J}$ respectively, so the effect of changing magnetic field strength is significantly more pronounced. Secondly, the magnetopause becomes more north-south asymmetric with Ganymede's increasing distance from the Jovian plasma sheet's centre. As the result, K-H wavevectors becomes more orthogonal to adjacent magnetic fields near the subflow region ($|Y| < \sim 2 R_G$). This can increase the Eq. 5.1 threshold and suppress K-H instability where the wavevectors are also insufficiently parallel to the bulk flow shears. Relative influences of the two factors are not precisely quantified, and beyond the reach of the estimative analytical model used in this work.

K-H instability onset is also impacted not only by Ganymede's spatial position, but also temporal changes in the Jovian plasma sheet. Figure 41 shows K-H instability assessment under $\pm 50\%$ changes in upstream plasma flow speed (a-b) and mass density (c-d) with Ganymede being fixed at centre of the Jovian plasma sheet ($\lambda_{III} = 248^\circ$). Note that temporal Jovian plasma changes are likely less than $\pm 50\%$, but this magnitude is chosen for consistency with changes induced by Ganymede's spatial variation. The temporal test reveals K-H instability to be more sensitive to flow speed changes than mass density, which is explained through Eq. 5.1 where the instability is linearly proportional to flow speed but inversely proportional to mass density. The effect of Jovian plasma flow speed is opposite to that of the magnetic field it carries, and less in magnitude given the inequality's dependence on the field strength squared. Interestingly, the impact of -50% flow speed (Figure 41a) is significantly greater than that of +50% flow speed (Figure 41b). The asymmetry occurs because the bulk flow shear is almost parallel to adjacent magnetic fields immediately

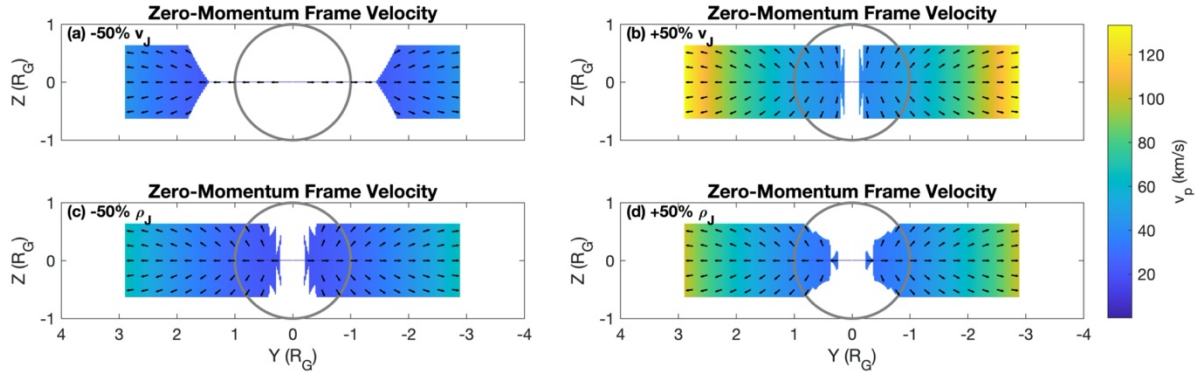


Figure 41: K-H instability onset assessment and K-H wave phase velocity under (a-b) $\pm 50\%$ changes in Jovian plasma bulk speed and (c-d) $\pm 50\%$ changes in Jovian plasma mass density. Ganymede is fixed at centre of the Jovian plasma sheet ($\lambda_{\text{III}} = 248^\circ$). Image courtesy of Kaweeyanun et al., (2021), *Journal of Geophysical Research: Space Physics*.

above/below the subflow point, so the magnetopause is highly K-H stable irrespective of K-H wavevector orientation.

As stated in Chapter 5, Section 5.1, the thin model onset inequality is an imperfect method that is most reliable when the K-H wavevector is strongly (if not perfectly) orthogonal to the adjacent magnetic fields, and the magnetic shear angle is small. Since the K-H wavevector orientation is not constrained in my analysis, the orthogonality condition is automatically satisfied. In fact, for each K-H unstable point on the magnetopause, the satisfying wavevector must be field-orthogonal as the orientation maximises the difference between two sides in Eq. 5.1. Then, as shown earlier in Figure 23, magnetic shear angle (θ_{sh}) across Ganymede's magnetopause is always within 10° of the maximum 180° , so the guide field component is always small and the magnetic configuration can be considered a variant of the small-shear regime. Hence, the thin model approximation is a robust predictor for Ganymedean K-H instability onset assessment.

I assume stagnant bulk plasma flow inside Ganymede's magnetosphere. It must be noted that numerical simulations of Ganymede predict flows of $|\mathbf{U}_G| = 20 - 50$ km/s in direction antiparallel to the external Jovian-side flow (e.g., Jia et al., 2009). The inclusion of finite-speed Ganymedean-side flow, fortunately, does not significantly impact the size of K-H unstable area.

Hence, I maintain the conclusion that linear K-H instability can occur on Ganymede's magnetopause flanks at all-latitudes, irrespective of the moon's position in the Jovian plasma sheet. The magnetopause becomes more K-H unstable when Ganymede is close to the centre of the Jovian plasma sheet. Once formed, K-H waves propagate downstream in direction of the external Jovian-side plasma flow. Onset of K-H instability is sensitive to changes in upstream plasma-magnetic conditions, most notably the magnetic field strength of Jupiter. Assessment of linear K-H instability onset, however, only marks the beginning of instability growth, whose research will be the subject of the next section.

5.3. The Finite Larmor Radius Effect

From this point onward, I focus exclusively on two equatorial points on Ganymede's magnetopause flanks with coordinates $(0, \pm Y_{eq}, 0)$ in GphiO coordinates. Ganymede is assumed to locate at the centre of the Jovian plasma sheet, where plasma is densest but its magnetic field is weakest. The magnetopause flanks are parallel to the incoming Jovian plasma flow (Figure 39), therefore external plasma and magnetic field are not compressed by collision with the boundary. Two leftmost columns in Table C summarise Jovian and Ganymedean plasma-magnetic parameters adjacent to the two equatorial magnetopause flank points, which are consistent with values used for onset assessments of magnetic reconnection and K-H instability. The Jovian and Ganymedean plasmas are both considered as singly-charged, single-species fluids with listed averaged ion masses and effective temperatures (Eviatar, Strobel, et al., 2001; Kivelson et al., 2004; Jia et al., 2008; Jia et al., 2009).

Four rightmost columns in Table C provide comparative plasma-magnetic parameters along Mercury's magnetopause flanks, each separating external magnetosheath from internal magnetosphere (Sundberg et al., 2010). Given Mercury's eccentric orbit around the Sun, conditions in the magnetosheath vary within ranges defined at the planet's peri- and aphelion positions. The external IMF is stronger than Mercury's internal field, while the solar wind temperature is colder than its magnetospheric counterpart. The situation is opposite on Ganymede's magnetopause, where the external Jovian magnetic field is weaker but its plasma possesses warmer temperature. The Hermean magnetospheric plasma is assumed to

be protons-only following Sundberg et al., (2010) – another contrast from much heavier Jovian ions. Lastly, the solar wind plasma is warmer, and the IMF is stronger, as Mercury moves from the aphelion to the perihelion.

Table C: Comparing plasma-magnetic conditions at equatorial magnetopause flanks between Ganymede (at Jovian plasma sheet's centre) and Mercury (at perihelion and aphelion points)

	Ganymede		Mercury (perihelion)		Mercury (aphelion)	
	J	G	MSH	MSP	MSH	MSP
B (nT)	67	122	46	15	21	15
T_i (10^6 K)	0.70	0.10	8	23	6	23
M_i (amu)	14	16	1	1	1	1

Note: Magnetic field strengths for Ganymede's system are taken from Figure 22e and Figure 22f (red dots). Jovian ion mass/temperature are taken from Kivelson et al., (2004), Jia et al., (2008), and Jia et al., (2009). Ganymede ion mass/temperature are taken from Eviatar et al., (2001a). Mercury system's parameters are taken from Sundberg et al., (2010). J = Jupiter, G = Ganymede, MSH = (solar) magnetosheath, MSP = (Mercury) magnetosphere.

Using parameters in Table C, I can calculate the gyroradius (Eq. 5.6), the gyrofrequency (Eq. 5.7), the gyro-viscous coefficient (Eq. 5.8), and gyration speeds for ions near Ganymede's and Mercury's equatorial magnetopause flanks in Table D. The Jovian-side ion gyroradii of 62.5 km are significant compared to Ganymede's magnetopause thickness of <400 km (Kivelson et al., 1998), confirming that the FLR effect is unlikely to be negligible. The Jovian-side gyroradii are also much greater than their Ganymede-side counterpart, so K-H instability growth asymmetry will be driven by external gyromotion. In contrast, the Hermean FLR effect is driven by internal gyromotion, whose gyroradii of 420 km are comparable to its magnetopause thickness.

The gyrofrequency of Jovian-side and Ganymede-side ions are similar in magnitude. Given the difference in gyroradii, the Jovian-side gyro-viscous coefficients are two magnitude orders larger than the Ganymede-side coefficients. The Hermean gyro-viscous coefficients are even larger, exceeding the Ganymede counterparts by 1 – 2 orders of magnitude. The same pattern also emerges when the ion flow speeds are calculated, which is expected given their similar dependence on the ion gyroradii.

Table D: Properties of ion gyromotion at equatorial magnetopause flanks of Ganymede (at Jovian plasma sheet's centre) and Mercury (at perihelion and aphelion points).

	Ganymede		Mercury (perihelion)		Mercury (aphelion)	
	J	G	MSH	MSP	MSH	MSP
R (km)	62.5	4.74	80.8	420	153	420
Ω (s ⁻¹)	0.46	0.73	4.60	1.50	2.10	1.50
η (10 ⁸ m ² /s)	4.48	0.04	75.0	661	123	661
v_i (km/s)	28.7	3.46	371	630	322	630

In Chapter 5, Section 5.1, I take three parameters to be the quantifiers of the FLR effect – the ion flow shear, the bulk flow shear, and the sum of external and internal gyro-viscous coefficients. Table E shows the values of these quantifiers for Ganymede's and Mercury's magnetopause flanks. The ion flow shears on Ganymede's magnetopause flanks are 10 times smaller than those on Mercury's. The Ganymede bulk flow shears are ~ 3 times smaller, and the sum of gyro-viscous coefficients are over 100 times smaller. All three parameters suggest that the FLR effect at Mercury should be much more pronounced than at Ganymede.

Without using the full MHD-FLR theory, the bulk flow shear and the gyro-viscous coefficients cannot offer more precise quantification of the FLR effect. The ion flow shears, however, may offer some insight given their described role in Figure 39. Assuming linear proportionality between the FLR effect and ion flow shear magnitude, the inter-flank asymmetry in K-H instability growth should be 10 times stronger at Mercury than Ganymede. If MESSENGER observations indicate that 90% (40% enhancement) of Mercury's K-H events favour the Hermean dusk flank, then 54% (4% enhancement) of Ganymede's K-H events should occur on the sub-Jovian flank. The asymmetry is small but notable and likely verifiable by future in-situ observations.

Table E: Quantifiers of the FLR effect at Ganymede's and Mercury's equatorial magnetopause flanks

	Ganymede	Mercury (perihelion)	Mercury (aphelion)
$v_{i,sh}$ (km/s)	+25.2	-258	-308
U_{sh} (km/s)	~140	~400	~400
$\sum \eta$ (10^8 m ² /s)	4.52	736	784

How sensitive are the FLR effect to changes in magnetopause conditions? I recalculate the ion flow speed (v_i) and the gyro-viscous coefficient (η) external and internal to Ganymede's magnetopause under $\pm 50\%$ changes in average ion mass (M_i), ion temperature (T_i), and local magnetic field strength (B). Sensitivity is dictated by relations $v_i \propto \sqrt{\frac{T_i}{M_i}}$ and $\eta \propto \frac{T_i}{B}$. Hence, the ion temperature affects both parameters while ion mass and magnetic field strength impact only ion flow speed and gyro-viscous coefficient respectively. The single largest change occurs when the magnetic field strength (B_J , B_G) is halved, leading to doubling of the gyro-viscous coefficient.

The ion flow speed is less sensitive due to the square root. A +50% change in ion temperature ($T_{i,j}$, $T_{i,G}$) will increase $v_i \rightarrow 1.22v_i$ for Ganymede, resulting in an FLR effect ~8 times weaker than at Mercury – which does not significantly impact our original result. Hence, the comparative strengths of the FLR effect at Ganymede and Mercury are not very sensitive to spatial and temporal condition changes. This finding is notable as re-analysis of Galileo's plasma measurements show Ganymede's internal plasma may be warmer than $T_{i,G} = 1$ eV used (Collinson et al., 2018). It also indicates that Ganymede's equatorial magnetopause flank points should maintain similar inter-flank K-H instability growth asymmetry at all times. Additionally, Table E assumes stagnant Ganymedean bulk plasma flow for calculation of U_{sh} . As mentioned in Chapter 5, Section 5.2, non-stagnant internal flow is a feature of the moon's numerical simulations. If included, the bulk flow shear will be slightly reduced and so does the FLR effect at Ganymede. The magnitude of this impact is unclear, but highly unlikely to affect the main conclusions drawn given that internal flow will be quite slow.

Quantification of the FLR effect discussed here relies heavily on comparison between Ganymede's and Mercury's similar-sized magnetopauses, which can be hindered by a couple of factors. Firstly, Figure 39b assumes the external IMF is strongly northward. In reality, the solar magnetic field continuously rotates between northward and southward orientations. Figure 39a and Figure 39b are therefore not truly equivalent. The Hermean ion flow shear may hence overestimate Mercury's FLR effect, and artificially suppress Ganymede's FLR effect when the two systems' ion flow shears are compared. Secondly, linear K-H instability growth rate depends on the K-H wavelength ($\lambda_{KH} = \frac{2\pi}{k}$) under both the MHD-FLR theory and kinetic simulations (Nakamura et al., 2010; Sundberg et al., 2010), but the parameter is not accounted for in this analysis. Since K-H wavelengths are poorly constrained at both Ganymede (1,050 – 1,400 km) and Mercury (500 – 5,000 km), highly divergent wavelengths can exaggerate or minimise comparative strength of the FLR effect.

My evaluation of the FLR effect has focused on equatorial magnetopause flank points ($X = 0$, $Z = 0$). However, recalling from Chapter 5, Section 5.2, the conditions of (1) near-180° magnetic shear angle between Jupiter's and Ganymede's magnetic field and (2) near field-orthogonal bulk plasma flow are satisfied across the entire magnetopause flanks. This means that the FLR model in Figure 39 is applicable for all flank latitudes, and so is my prediction of enhanced sub-Jovian K-H instability growth. Study of the FLR effect offers insight into the instability as it transitions from linear to nonlinear phase, but does not say much regarding K-H vortex formation, which will be the subject of this chapter's last section.

5.4. Viability of Nonlinear K-H Vortices on Ganymede's Magnetopause

Given the lack of mathematical bridge between linear and nonlinear phases of K-H instability, the formation process for K-H vortex remains unclear. Assuming a picture that large K-H wave oscillation will eventually trigger vortex formation, I can set an amplitude threshold beyond which the K-H vortices are incredibly likely, which in this study is chosen to be $A(t_f) = 100$ nT where t_f is the (unknown) time interval between K-H wave and vortex formation.

Consider a simple scenario where Ganymede lies at the centre of the Jovian plasma sheet ($\lambda_{III} = 248^\circ$). How far can a K-H wave propagate until its amplitude reaches the vortex formation threshold? To answer this question, I set as independent variables the initial K-H wave amplitude $A_i = [1, 5, 10, 15, 20]$ nT and the linear K-H instability growth rate $\gamma = [0.01, 0.1, 1, 10]$ s⁻¹. These are the two governing parameters of exponential K-H instability growth in Eq. 5.3, and their chosen ranges reflect realistic potential values for Ganymede magnetopause conditions.

Table F shows the distance (Eq. 5.4) that a K-H wave, formed near Ganymede's equatorial magnetopause flanks ($X = 0, Z = 0$) under varied initial amplitude and linear growth rate. A K-H wave is assumed to propagate at constant phase speed $|v_{KH}| = 85$ km/s parallel to the moon's magnetic equator. The Ganymede magnetosphere is about $7 - 9 R_G$ in length from the subflow point to the end of the magnetotail (Jia, Kivelson, et al., 2010). Bolded distances hence correspond to lengths greater than $7 R_G$, beyond which K-H vortices will assuredly not form within Ganymede's magnetosphere. The table suggests that a K-H wave will become a K-H vortex within reasonable distance/timeframe unless the initial wave amplitude and the linear growth rate are both small ($A_i < 10$ nT and $\gamma = 0.01$ s⁻¹). Galileo observations suggest magnetic fluctuations of amplitude $A_i = 10 - 20$ nT near Ganymede's magnetopause. Therefore, barring very low linear growth rate, nonlinear K-H instability should be viable in the moon's magnetospheric environment.

Does this mean that K-H vortices (and subsequently viscous-based energy transport) are common on Ganymede's upstream magnetopause? This chapter's discussion thus far assumes K-H instability grow in isolation from other magnetopause interactions, specifically magnetic reconnection. Analytical assessment suggests reconnection to be highly common given always-antiparallel magnetic field configuration on Ganymede's magnetopause (Chapter 3, Section 3.3). Reconnection events are shown to disrupt formation of K-H vortices by causing abrupt, drastic changes to magnetopause conditions (Nakamura et al., 2020). The question, then, is whether reconnection occurs at sufficiently frequent rate to majorly suppress K-H instability growth.

Table F: Distance (in R_G) travelled by linear K-H wave before K-H vortex formation ($A(t_f) = 100$ nT) for different initial amplitudes (A_i) and linear growth rate (γ).

γ (s^{-1}) A_i (nT)	0.01	0.1	1	10
1	14.89	1.49	0.15	0.01
5	9.68	0.97	0.10	<0.01
10	7.44	0.74	0.07	<0.01
15	6.13	0.61	0.06	<0.01
20	5.20	0.52	0.05	<0.01

Bolded numbers are greater than $7 R_G$ – the likely maximum distance between equatorial magnetopause flank and the end of Ganymede’s magnetotail.

To answer this question, I return to Mercury as an analogue. In-situ data from MESSENGER shows that only 10% of observed K-H vortices occur under southward IMF when reconnection is expected to be frequent (Liljeblad et al., 2014), indicating a strong suppressive effect of reconnection on K-H instability. The unnormalised Hermean reconnection rate is given by its electric field $E_{MR} = 0.3 - 3$ mV/m (Gershman et al., 2016), which is one order of magnitude smaller than Ganymede’s typical electric field $E_{MR} = 2 - 20$ mV/m (Chapter 3, Section 3.3; Zhou et al., 2020). As the two magnetospheres are similar in size, reconnection rates can be compared directly through electric field strengths. Reconnection at Ganymede is expected to occur more frequently than at Mercury, especially as the Jovian magnetic field is always southward while the solar wind has variable orientation. Suppression of K-H instability should therefore be even stronger at Ganymede, and few K-H vortices are expected once reconnection effects are taken into account.

Consequently, magnetic reconnection is likely the dominant mechanism for particle and energy transport into Ganymede’s magnetosphere. Thus, Chapter 4’s assumption of reconnection-governed magnetic flux opening and closure was not displaced. Despite its likely limited influence, it would be premature to dismiss K-H instability given lack of direct observational evidence. Prolonged in-situ observations of Ganymede’s magnetopause flanks

will be necessary to ascertain the impact of K-H instability, but this will not be available for many years to come.

Since Chapter 3, Ganymede's upstream magnetopause has been studied as a driver of plasma convection. However, research on the boundary's impacts has not been exhausted. An extended current sheet, the upstream magnetopause is integral to the evolving structure of Ganymede's internal magnetic field, which in turn impact induction in the moon's subsurface ocean. The next and final work chapter will capture induction as consequence of interaction between Ganymede's and Jupiter's magnetosphere, which has never been studied under current literature on Ganymede.

Chapter 6: Magnetic Induction by Ganymede's Upstream Magnetopause

In this chapter, the Chapman-Ferraro (C-F) magnetic field, produced by magnetopause currents upstream of Ganymede, is established as a potential candidate to excite secondary magnetic induction in the moon's subsurface ocean. The C-F signal varies at Jovian half-synodic (5.27 hours) due to oscillation in near-magnetopause conditions, and its diffusion is shown to penetrate through the moon's conducting ionosphere and later into the ocean. The response amplitude from the underground fluid shell is expected to lie between $|B_{\text{ind}}| \sim 1 - 10 \text{ nT}$ with the same Jovian half-synodic period, and is likely detectable under future prolonged observations.

The study asserts that Ganymede's internal magnetism may play a significant role in secondary magnetic induction, given that the C-F magnetic field depends on existence of the magnetopause boundary. However, my research is not conclusive on the C-F field's viability as a probing tool into structure of the subsurface ocean. The analytical description of the exciting field remains largely estimative, and direct numerical simulations will be likely necessary to firmly establish magnetopause-induced magnetic response.

6.1. From Magnetopause to Subsurface Ocean

When I previously discussed modelling of near-magnetopause conditions (and subsequently magnetic reconnection), a pattern of periodicity emerges. Plasma-magnetic parameters vary at Jovian half-synodic period (5.27 hr), corresponding to Ganymede's two trips through the Jovian plasma sheet per the planet's rotation. Under Ampère's law, Jupiter-Ganymede magnetic gradients across the magnetopause generate currents along the boundary known as the Chapman-Ferraro (C-F) currents (Chapman & Ferraro, 1940), whose presence in turn generates a C-F magnetic field that confines Ganymede's magnetosphere (Figure 12). Given the half-synodic periodicity of near-magnetopause conditions, the C-F magnetic field will alter Ganymede's internal magnetic field at the same period.

Since Ganymede is expected to possess a subsurface ocean, this structural oscillation becomes an important point of research due to magnetic induction. When a layer of electrically conducting fluid experiences change in its surrounding magnetic field, it creates a secondary magnetic field which negates magnetism inside the volume confined by the layer. Note that this process differs from dynamo generation discussed in Chapter 1, Section 1.2, where the primary internal magnetic field is ‘inducted’ from *convection* of the conducting fluid. For secondary induction, the fluid is not required to be in motion, with magnetism arising solely from conductive interaction with the primary field. Through induction, Ganymede’s subsurface ocean can be characterised in its structure, which represents a crucial advancement in our understanding of the potentially life-hosting liquid layer.

The C-F magnetic field can be analytically derived by assuming the magnetopause currents to form a circular loop around the moon’s equator as shown in Figure 42. Following the right-hand thumb rule, this produces a field running parallel/antiparallel to Ganymede’s rotation axis. This method has been previously used at Mercury (Glassmeier, Grosser, et al., 2007, 2007b), where secondary magnetic induction also occurs (albeit due to near-surface outer iron core rather than a subsurface ocean). The equatorial current model is in fact more

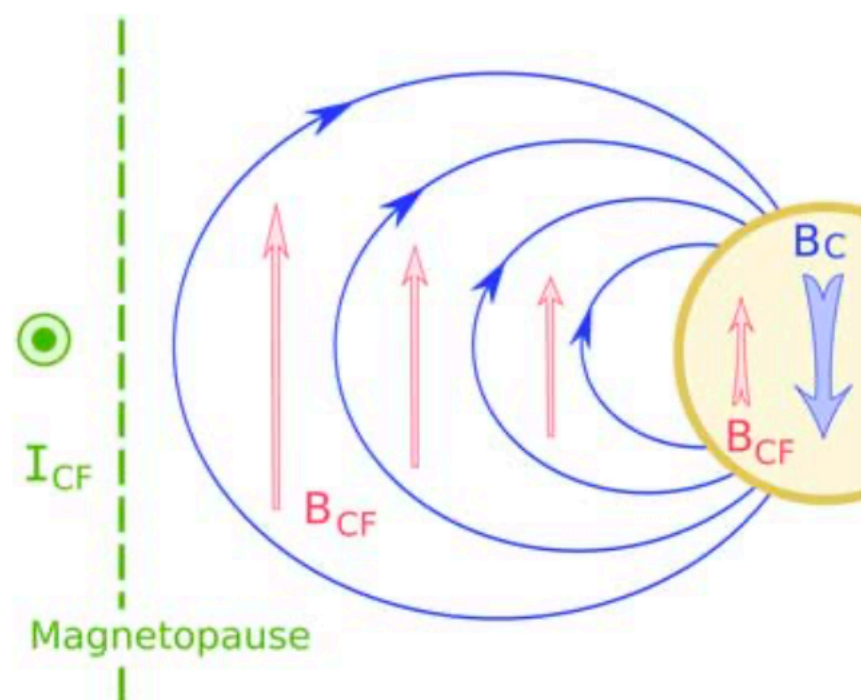


Figure 42: Equatorial current loop model that produces the Chapman-Ferraro (C-F) magnetic field (B_{CF}) parallel/antiparallel to a body’s dipole axis. Image courtesy of Glassmeier, Auster, et al., 2007, *Geophysical Research Letters*.

suitable for Ganymede as its cylindrical magnetosphere and north-south aligned magnetopause.

The model begins with computing the C-F equatorial current density (Glassmeier, Grosser, et al., 2007)

$$\mathbf{j}_{CF} = -\frac{1}{\mu_0}(\delta\mathbf{B} \times \mathbf{N}) \quad (6.1)$$

where $\delta\mathbf{B} = \mathbf{B}_I - \mathbf{B}_G$ is the magnetic gradient across Ganymede's magnetopause, and $\mathbf{N}$ is the upstream-ward magnetopause unit normal vector. For simplicity, Eq. 6.1 can be calculated only at the subflow point, obtaining a singular value of $\mathbf{j}_{CF}$. The current magnitude along the equatorial loop is then estimated through

$$I_{CF} = \oint \mathbf{j}_{CF} \cdot d\mathbf{l} = j_{CF}\pi r_{MP} \quad (6.2)$$

where the 'loop' is a semi-circular arc defined by the magnetopause standoff distance r_{MP} . Consequently, the C-F magnetic field vector has components (0, 0, B_{CF}) where

$$B_{CF} = \frac{\mu_0 I_{CF}}{2\pi} C(r_{MP}, r) \quad (6.3)$$

which is effectively Ampère's law with inverse distance function

$$C(r_{MP}, r) = \frac{1}{r_{MP} + r} \left[K\left(\frac{r}{r_{MP}}\right) + \left(\frac{r_{MP} + r}{r_{MP} - r}\right) E\left(\frac{r}{r_{MP}}\right) \right] \quad (6.4)$$

where r is the radial distance from Ganymede's centre at which the C-F magnetic field is measured (Glassmeier, Grosser, et al., 2007). Functions $K\left(\frac{r}{r_{MP}}\right)$ and $E\left(\frac{r}{r_{MP}}\right)$ are complete

elliptic integrals of first and second kind respectively²². The C-F magnetic field strength in Eq. 6.3 expectedly decreases away from the magnetopause current sheet i.e., decreasing r . More problematically, however, the field magnitude approaches infinity as $r \rightarrow r_{MP}$, which is not the case in reality. The asymptotic fault arises from the ‘thin-wire’ assumption used to approximate the magnetopause current sheet. Since $C(r_{MP}, r)$ only spikes when r is quite close to r_{MP} ($r \sim 1.5 R_G$), Eq. 6.3 is applicable at Ganymede’s surface ($r = R_G$), where the C-F magnetic field is commonly measured.

As the C-F magnetic field varies at Jovian half-synodic period, the magnetopause boundary may also oscillate in correspondence. However, while the change in upstream Jovian acting pressure is half-synodic, its correlation with magnetopause movement is complicated by multiple factors including magnetic field orientations and Alfvénic plasma-magnetic interactions. As the result, the magnetopause movement remains difficult to constrain over the half-synodic period.

Once produced, the C-F magnetic field diffuses radially toward Ganymede as summarised in Figure 43. The diffusion remains unimpeded until the magnetic signal encounters a conductive fluid layer. The first diffusive obstacle around Ganymede is actually not the subsurface ocean, but the ionosphere just above the moon’s surface. The ionosphere refers to the uppermost portion of Ganymede’s atmosphere, where gas particles become sufficiently ionised to form a conductive plasma layer (Eviatar, Vasyliūnas, et al., 2001).

²² The integrals are $K(k) = \int_0^{\pi/2} \frac{d\theta}{\sqrt{1-k^2 \sin^2(\theta)}}$ and $E(k) = \int_0^{\pi/2} \sqrt{1-k^2 \sin^2(\theta)} d\theta$. They exist as pre-packaged functions in MATLAB.

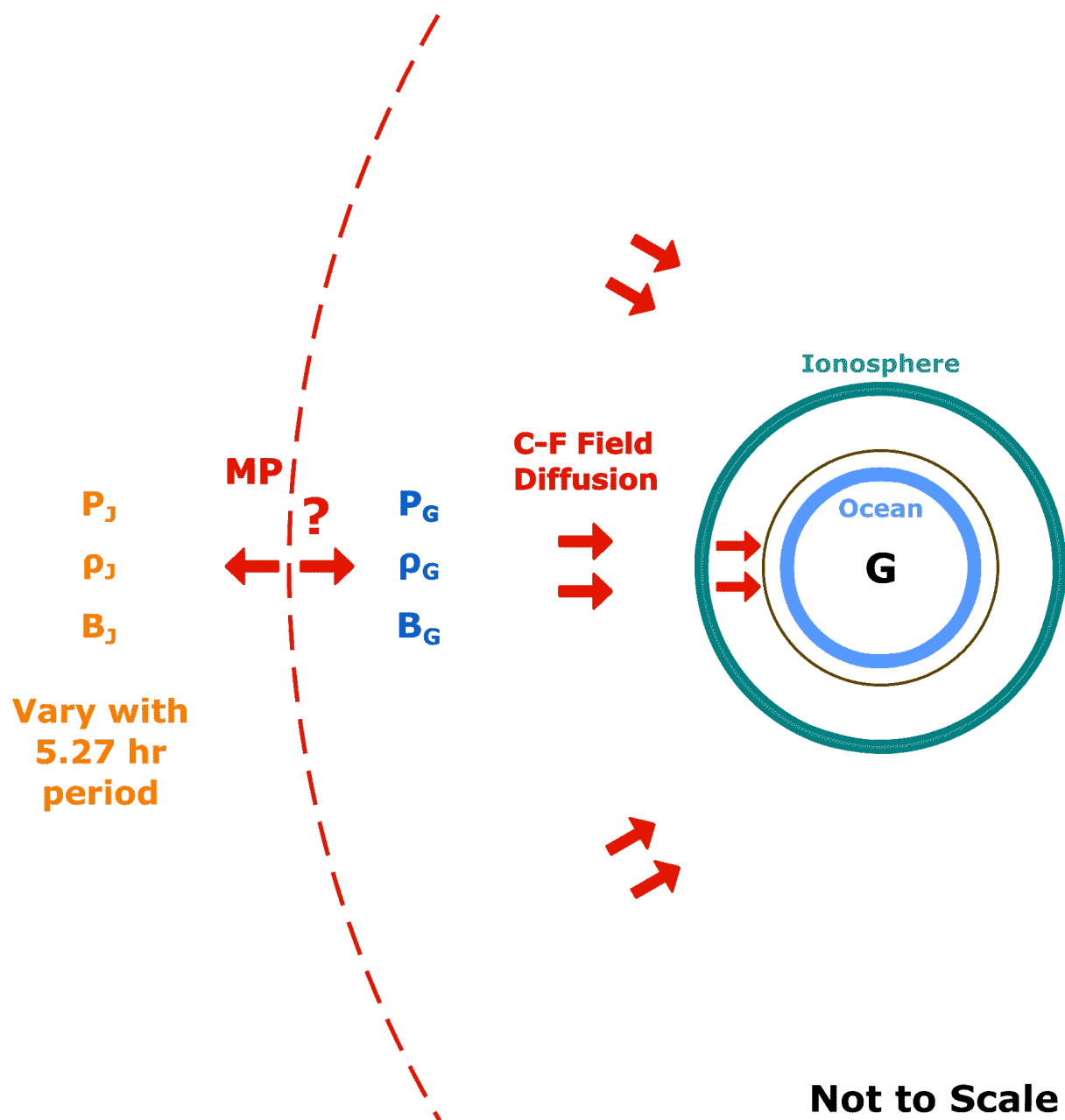


Figure 43: Summary schematic for diffusion of the Chapman-Ferraro (C-F) magnetic field produced by Ganymede's upstream magnetopause, which varies at Jovian half-synodic period. The diffusion penetrates through the moon's ionosphere but is stopped by its subsurface ocean.

When a varying magnetic field arrives at a stationary fluid shell of thickness L and constant finite conductivity σ , induction occurs under equation (Saur et al., 2010)

$$\frac{\partial \mathbf{B}}{\partial t} = -\nabla \times \left(\frac{1}{\sigma \mu_0} \nabla \times \mathbf{B} \right) \quad (6.5)$$

where $\mu_0 = 4\pi \times 10^{-7}$ H/m is the permeability of free space. In Eq. 6.5, variation in the *exciting*²³ magnetic field (left-hand side) diffuses into the fluid layer (right-hand side), leading to generation of an opposing secondary field as shown in Figure 44. Setting the fluid stationary assumes that its convection effects, which can cause additional induction, are negligible compared to those of diffusion. Oscillation of the exciting magnetic field is expressed through angular frequency

$$\omega = \frac{2\pi}{T} \quad (6.6)$$

where T is the magnetic field's oscillation period. To maintain Eq. 6.5, condition

$$\omega \ll \sigma \mu_0 c^2 \quad (6.7)$$

(where c is the speed of light) is required to neglect the Ampère displacement current term.

Magnetic diffusion is characterised by its typical depth and time scales

$$\delta_{\text{diff}} = \sqrt{\frac{2}{\sigma \mu_0 \omega}} \quad (6.8)$$

²³ Also known as inducing magnetic field.

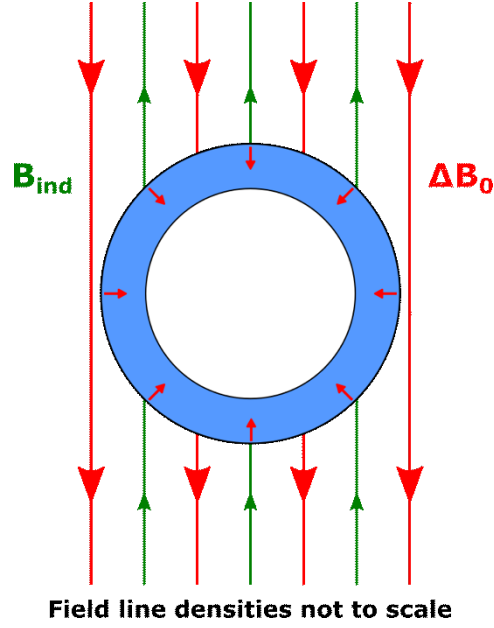


Figure 44: Illustration of magnetic diffusion (red arrows) into a conducting fluid shell (blue), leading to magnetic induction. Vertical red field lines show variational direction of the exciting field (and not the field itself). Green field lines show the induced magnetic field opposing the variation. Magnetic field line densities do not correlate to field strengths.

$$T_{\text{diff}} = \sigma \mu_0 L^2 \quad (6.9)$$

where δ_{diff} and T_{diff} are diffusion depth and time respectively, and σ is layer conductivity in Siemens per metre (S/m). Eq. 6.8 asserts that exciting fields of slower variations are in fact more penetrating. The fluid's conductivity also decides the depth of diffusion, where an infinite value implies complete non-penetration. These length and time scales can then be normalised with respect shell thickness (L) and exciting field oscillation period (T) respectively.

A magnetic signal is said to completely penetrate through a conducting fluid shell if

$$\log\left(\frac{\delta_{\text{diff}}}{L}\right) > 0 \quad (6.10)$$

and the diffusion is considered to be within reasonable time if

$$\log\left(\frac{T_{\text{diff}}}{T}\right) < 0 \quad (6.11)$$

When Eq. 6.10 and Eq. 6.11 are satisfied, induction effects from the penetrated fluid layer are said to be minimal since the opposing secondary field strength is negligible. If this is the case for Ganymede's ionosphere, then the C-F magnetic field will reach the moon's subsurface ocean because the surface ice shell is effectively non-conducting. The diffusion length and time scales can then be recalculated for the ocean. If the C-F signal does not penetrate through the liquid shell, then the ocean's induction must be significant.

The mathematics for computing induced magnetic response is highly complex and will not be discussed in detail here. In general, an induction model considers spherical bodies with one or more conducting fluid layers, whose (assumed constant) conductivity can be finite or infinite (Saur et al., 2010). For a single layer of finite conductivity, the inductive response follows (Saur et al., 2010)

$$B_{\text{ind}}(t) = \kappa B_{\text{ind}}^{\infty} \left(t + \frac{\varphi}{\omega} \right) \quad (6.12)$$

The induced field has phase lag $-\frac{\pi}{2} < \varphi < \frac{\pi}{2}$ i.e., time delay from the exciting field. The constant $\kappa < 1$ is calculated at the outer surface of the conducting layer, and indicates that the response is weaker than not only the exciting field, but also B_{ind}^{∞} – the response produced by a perfectly conducting sphere ($\frac{\delta_{\text{diff}}}{L} \ll 1$, $\frac{T_{\text{diff}}}{T} \gg 1$), which exists as a simple limit case where the secondary magnetic field is a pure dipole (Saur et al., 2010)

$$\mathbf{B}_{\text{ind}}^{\infty}(t) = \frac{\mu_0}{4\pi} \frac{[3(\mathbf{r} \cdot \mathbf{M}(t))\mathbf{r} - |\mathbf{r}|^2 \mathbf{M}(t)]}{|\mathbf{r}|^5} \quad (6.13)$$

$$\mathbf{M}(t) = -\frac{2\pi}{\mu_0} \Delta \mathbf{B}_0(t) r_0^3 \quad (6.14)$$

where $\mathbf{r}$ is a position vector, r_0 is the outer radius of the fluid shell, and $\mathbf{M}(t)$ is the magnetic moment of the induced dipole, and $\Delta \mathbf{B}_0(t)$ is variation of a spatially uniform exciting magnetic field. The negative sign in Eq. 6.14 reflects the induced field's opposition to variation in the exciting field. If we consider a point along the shell's surface equator ($\mathbf{r} \perp \mathbf{B}_0$, $|\mathbf{r}| = r_0$), Eq. 6.13 can be reduced to

$$\mathbf{B}_{\text{ind}}^{\infty}(t) = \frac{\Delta \mathbf{B}_0(t)}{2} \quad (6.15)$$

Therefore, the inductive response of an infinitely conducting shell can be linked directly to magnitude of the exciting field (at the same location) by a factor of $\frac{1}{2}$. Reduction of conductivity produces fractionally smaller inductive response, which is dependent on the magnetic system considered. For instance, the fractional reduction factor is estimated to be 0.84 for Ganymede's subsurface ocean (Kivelson et al., 2002).

Eq. 6.13 – 6.15 illustrate that understanding of the exciting magnetic field is instrumental to derivation of secondary induced magnetic response. For Ganymede's subsurface ocean, magnetic excitation occurs through periodic variations in the moon's primary magnetic field, of which there can be multiple. Exciting magnetic fields are organised by their periods of variation. Under current literature, primary excitation periods experienced by Ganymede's subsurface ocean correspond Jupiter's synodic period (10.53 hr) and Ganymede's orbital period around Jupiter (~ 171.57 hr) (Seufert et al., 2011)²⁴. Synodic oscillation occurs when the moon's dipole field realigns in response to orientation change in the Jovian magnetic field, whose period is governed by Jupiter's rotation. Meanwhile, orbital oscillation is a by-product of Ganymede's eccentric orbit, which vary the moon's distance from Jupiter and thus the Jovian magnetic field strength felt by Ganymede.

Excitation effects from these two sources have been studied through three-layer, finite-conductivity simulation (Vance et al., 2021). As shown in Figure 45, the GphiO y-component of inductive response ($B_{\text{ind},y}$) and the phase delay (φ) are computed as functions of the subsurface ocean's depth and conductivity. Contour lines in each sub-plot indicate isolines, therefore one response magnitude (or phase delay) can be caused by multiple combinations of ocean depth and conductivity. There are three sets of isolines, two for the Jovian synodic oscillation and its second harmonic (with half-synodic period of 5.27 hr), and one for the Ganymedeian orbital oscillation. Since each isoline represents an equation of two unknown

²⁴ Seufert et al., (2011) also consider induction from variation of Jupiter's dayside magnetopause (~ 640 hr), but its effect is likely negligible at Ganymede given that it is not considered in Vance et al., (2021).

In the final two sections of this chapter, I will apply the principles discussed here to characterise Ganymede's oceanic response to changes in near-magnetopause conditions. My analysis begins with a determination of the C-F magnetic field's ability to penetrate the subsurface ocean (Section 6.2), followed by an evaluation for possible inductive response magnitudes (Section 6.3). This research aims to highlight the C-F magnetic field, and thus Ganymede-Jupiter magnetopause interactions, as a potential tool for studying of induction, but will not add isoline-based constraints to depth and conductivity of the subsurface ocean.

6.2. Diffusion of the Chapman-Ferraro Magnetic Field

To determine diffusion of the C-F magnetic field into Ganymede's ionosphere, the layer's thickness and conductivity must be characterised. The thickness of the ionosphere is not currently well constrained. For this analysis, I choose a large upper limit of 400 km, even though the real ionosphere may be as thin as 50 km or less, to limit chance of a false positive diffusive penetration. Thus, the ionospheric thickness is considered as an independent variable of range 0 – 400 km with 1 km resolution. The layer's conductivity is not better known. Galileo magnetic measurements give *height-integrated* conductivity of $\sigma_h = 2 \text{ S}$ (Kivelson et al., 2004), while particle observations give a much higher value of $\sigma_h = 100 \text{ S}$ (Eviatar, Vasyliūnas, et al., 2001). Given only two data points available, I consider each conductivity value as a distinct case. Note that the height-integrated conductivity is related to conductivity by $\sigma_h = \sigma L$.

Diffusion depth (Eq. 6.8) and time (Eq. 6.9) of the C-F magnetic field in Ganymede's ionosphere are calculated as functions of the ionosphere's thickness in Figure 46. The parameters are normalised and then expressed as logarithms for ease of understanding. Diffusion length (time) monotonically decreases (increases) with increasing thickness for both $\sigma_h = 2 \text{ S}$ (red) and $\sigma_h = 100 \text{ S}$ (blue). As expected, larger conductivity lowers the penetration depth of the C-F magnetic field into the ionosphere. Nevertheless, the layer cannot stop penetration even at its upper thickness limit, as the normalised diffusion depth always remain above 1 i.e., $\log\left(\frac{\delta_{\text{diff}}}{L}\right) > 0$. The normalised diffusion time is also below 1 at all times i.e.,

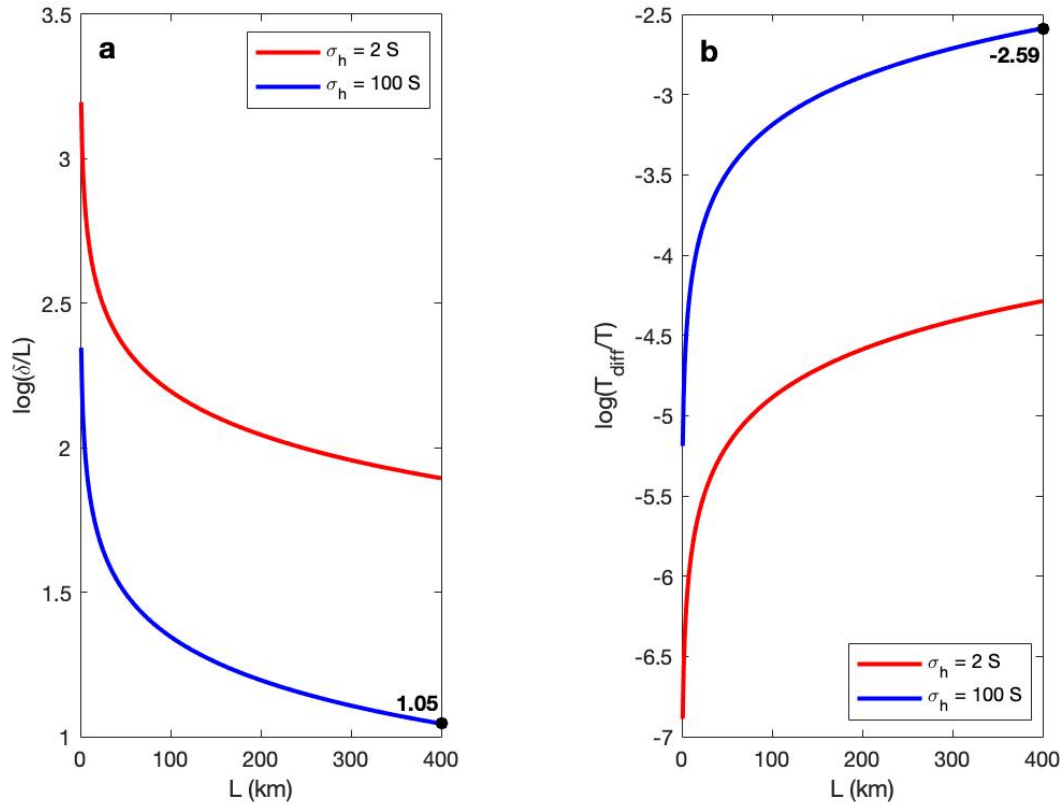


Figure 46: Normalised diffusion depth (a) and time (b) of the C-F magnetic field through Ganymede's ionosphere of varying thicknesses for low (red) and high (blue) height-integrated ionospheric conductivities.

$\log\left(\frac{T_{\text{diff}}}{T}\right) < 0$, indicating that diffusion is sufficiently rapid regardless of ionospheric thickness. Thus, Figure 46 shows the C-F magnetic field to comfortably penetrate through Ganymede's ionosphere, with little inductive response from the conductive layer. Subsequently, if the moon's atmosphere and icy surface are assumed to be non-conductive, then the diffusion should continue unimpeded until arrival at the liquid ocean's outer edge.

Similar to the ionosphere analysis, ocean thickness (depth) is set as an independent variable. There remain large uncertainties regarding the layer's thickness, whether the liquid is sandwiched between two thick ice shells or reaches completely down to the silicate mantle. Per Figure 15, the ocean may be as deep as 700 km, hence I set this value as the upper limit of the thickness range (0 – 700 km, 1 km resolution). The ocean's conductivity is completely unknown without further in-situ observations. So, following the procedure in Kivelson et al., (2002), the conductivity is assumed to be the same as Earth's i.e., $\sigma = 3 - 6$ S/m (Zheng et al., 2018). These are not height-integrated values, which mean Eq. 6.8 and Eq. 6.9 do not need

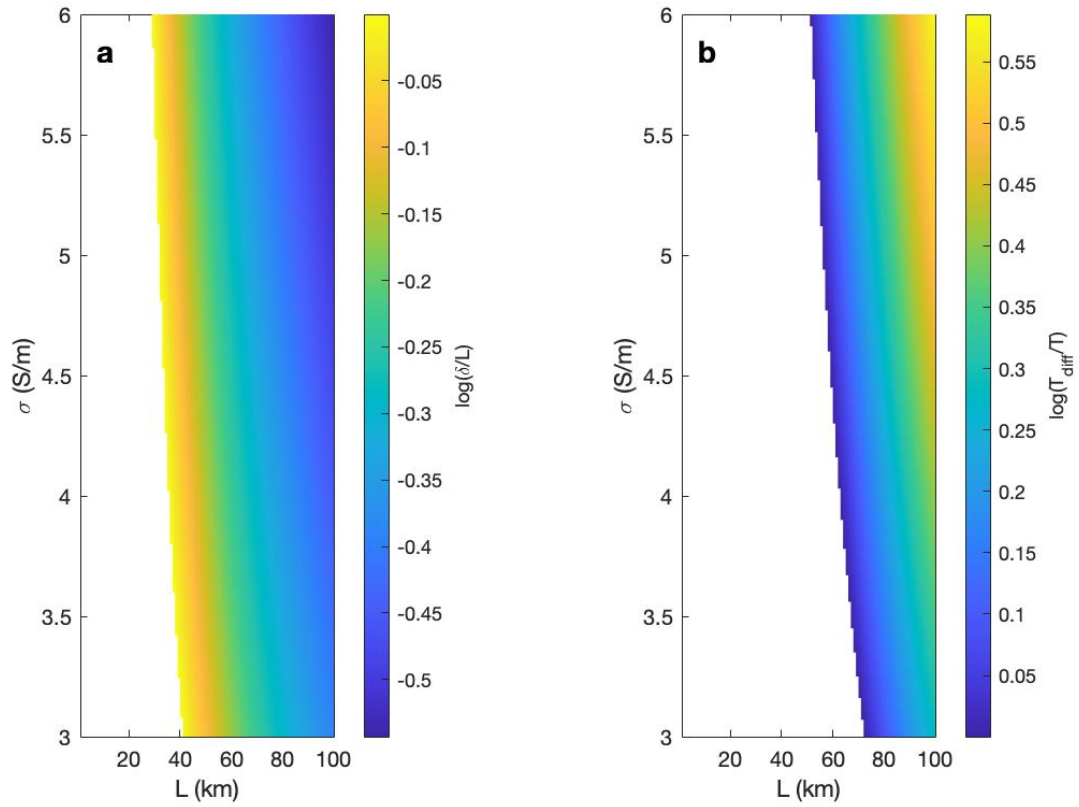


Figure 47: Normalised diffusion depth (a) and time (b) of the C-F magnetic field as contour plots with respect to Ganymede's subsurface ocean thickness and conductivity. In the white region, complete diffusive penetration occurs.

to be modified. Furthermore, the conductivity range is well-sampled, so I evaluate diffusion for every value at 0.01 S/m resolution.

As the result, the C-F magnetic field's diffusion depth and time in the subsurface ocean are shown as contour plots in Figure 47, where normalised lengths and times correspond to colour bars to the right of the sub-plots. Leftmost white regions indicate ocean thicknesses at which normalised diffusion depths and times are greater and lesser than 1 respectively. Here ocean thickness is shown up to only 100 km for clarity of the subplots, since further extension will not change the conclusions drawn. Minimum thicknesses for non-penetration and untimely diffusions are 30 – 40 km and 60 – 70 km respectively – differing because the two parameters are computed independently for each other. Between these thresholds (e.g., $L = 50$ km), diffusion will not penetrate the ocean but will occur in timely manner, hence the lower threshold is taken as the effective barrier to penetration. As expected, increasing ocean thickness monotonically decreases diffusion depth and increases diffusion time.

Since Ganymede's subsurface ocean thickness is unknown, it cannot be said definitively whether diffusive penetration will occur. But if the ocean thickness is significant relative to Ganymede's cross section (e.g., $L > 10$ km), then the C-F magnetic field's diffusion will likely be strongly attenuated even if complete penetration occurs. This means that on balance of probability, a non-negligible inductive response from the ocean should be expected. The secondary magnetic field will share Jovian half-synodic period with the C-F field's oscillation with some unknown phase delay.

6.3. Constraining the Chapman-Ferraro Magnetic Field

If diffusion of the Chapman-Ferraro (C-F) magnetic field is assumed to stop in Ganymede's subsurface ocean, then magnitude of the fluid's secondary inductive response is related to magnitude of the exciting field (Eq. 6.12 – 6.15), which corresponds to the variance of the C-F magnetic field i.e., $\Delta \mathbf{B}_0 = \Delta \mathbf{B}_{CF}$. Therefore, to derive an upper limit for the inductive response, the maximum variance of the C-F field must be obtained. Strength of the C-F magnetic field is proportional to magnetopause current density at the subflow point (Eq. 6.3), which is then proportional to magnetic gradient across the boundary (Eq. 6.1). The difference between Jupiter's and Ganymede's magnetic fields at subflow point can be computed by the idealised analytical model previously used to assess magnetic reconnection and Kelvin-Helmholtz instability (Chapter 3, Section 3.2).

In subsequent analyses, the C-F magnetic field is measured at Ganymede's equatorial surface $r = 1 R_G$. Here I make an initial assumption that the magnetopause remains fixed in location at $r_{MP} = 1.65 R_G$ (default model setting) as the moon traverses up-down inside the Jovian plasma sheet. Figure 48 shows the C-F magnetic field as a function of Jupiter's System-III east longitude (λ_{III}), which approximately correspond to Ganymede's relative position in the plasma sheet, calculated at resolution of $\Delta \lambda_{III} = 1^\circ$ over one Jovian synodic period. The C-F magnetic field notably dips in strength as the moon approaches the plasma sheet's centre ($\lambda_{III} = 68^\circ, 248^\circ$), and therefore does not exhibit monotonic increase (decrease) with the moon's increasing (decreasing) magnetic latitude.

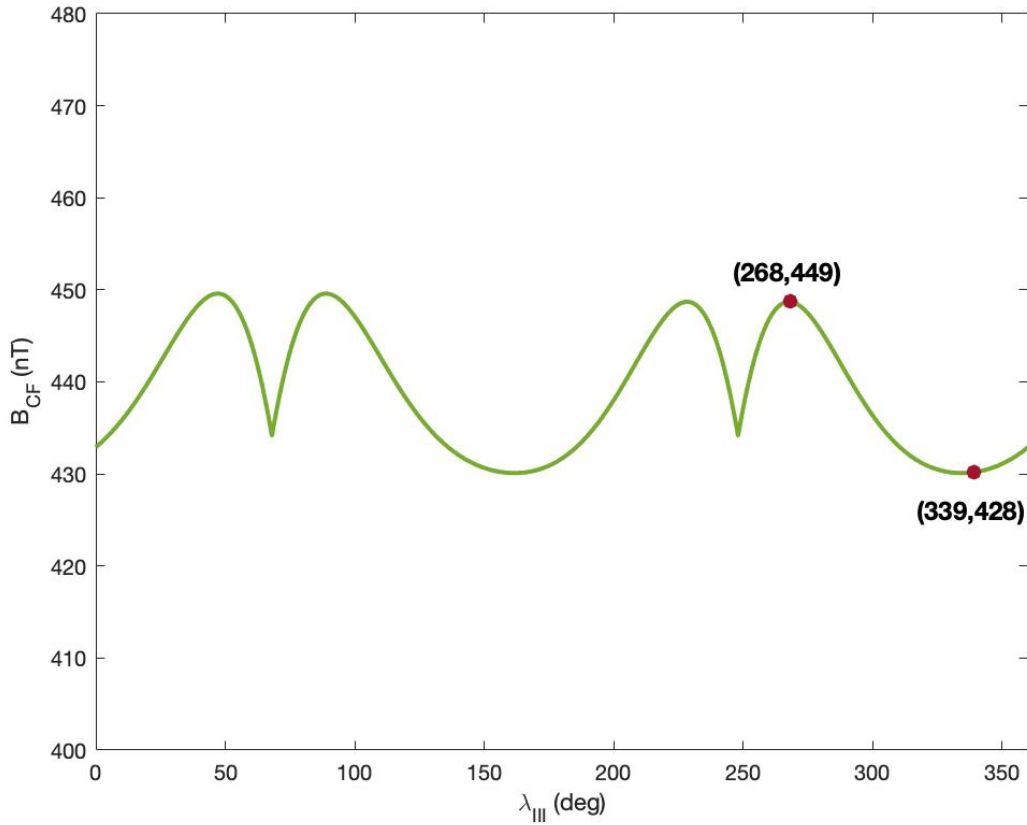


Figure 48: The C-F magnetic field magnitude ($B_{CF} = |\mathbf{B}_{CF}|$) at Ganymede's equatorial surface, estimated from magnetopause current at the subflow point, with respect to System-III Jovian east longitude.

To investigate these troughs, I compute Jupiter's and Ganymede's magnetic fields just inside and outside the magnetopause, alongside the external acting pressure in Figure 49. The Jovian-side field strength (orange) exhibits sinusoidal variations with respect to the Jovian east longitude, increasing in magnitude with Ganymede's absolute distance from the centre of the Jovian plasma sheet, as expected from current theoretical knowledge (Khurana, 1997; Connerney et al., 2018; Connerney et al., 2020). Meanwhile, the Ganymedean-side field strength (blue) exhibits near-centre dips seen in the C-F magnetic field. My analytical model computes B_G as a function of the external acting pressure (black), which exhibits a similar dipping pattern.

The acting pressure is the sum of Jovian magnetic pressure, plasma pressure and (magnetopause-orthogonal component of) dynamic ram pressure. When Ganymede leaves the plasma sheet's centre ($\lambda_{III} = 68^\circ, 248^\circ$) toward highest/lowest in-sheet latitudes ($\lambda_{III} = 158^\circ, 338^\circ$), the magnetic pressure increases as B_j departs from local minima, but the other two pressures decrease due to lowering plasma mass density (ρ_j). Competition between

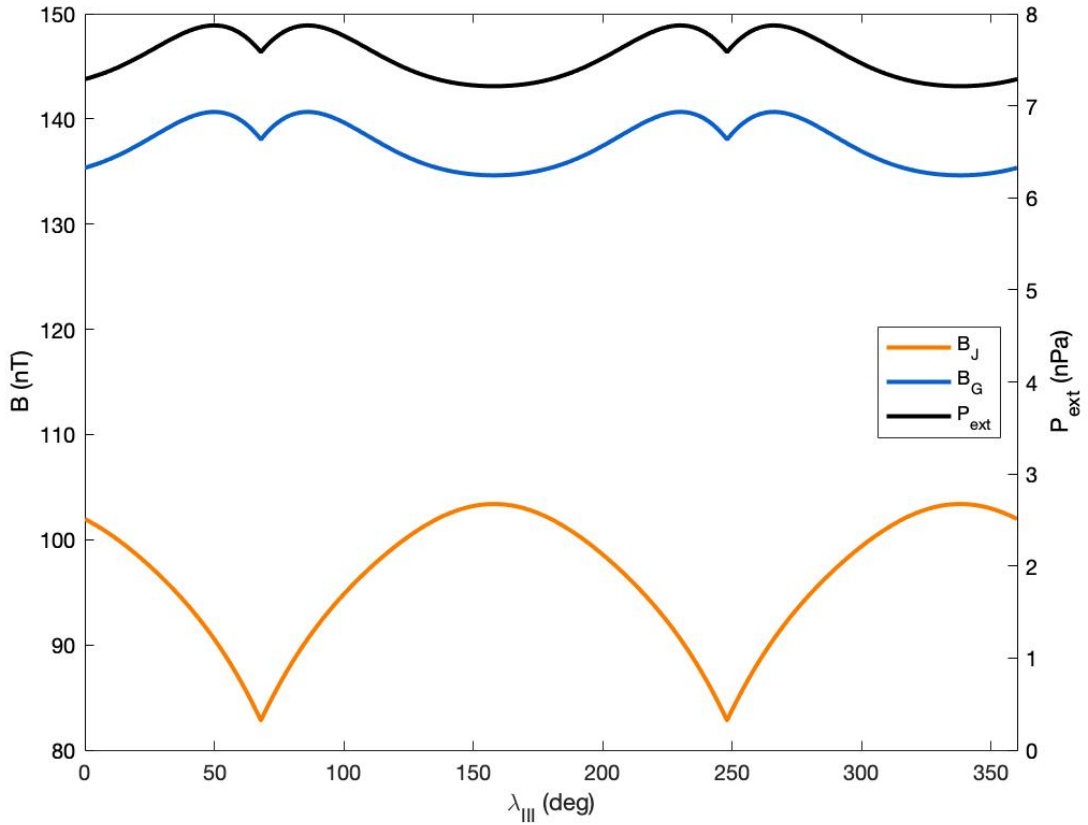


Figure 49: Jovian-side (orange) and Ganymede-side (blue) magnetic field strengths as functions of System-III Jovian east longitude. External acting pressure (black) is also shown as similar function.

these two trends manifest in the acting pressure, where magnetic pressure effect is dominant within $|\Delta\lambda_{\text{III}}| < \sim 20^\circ$ of the central longitudes, leading to net increase in acting pressure. Beyond this limit, plasma and dynamic pressure effects become dominant, leading to net decrease in acting pressure. This double-trough pattern is then passed on to cross-boundary magnetic gradient ($|\delta\mathbf{B}|$) and the C-F magnetic field ($|\mathbf{B}_{\text{CF}}|$) respectively.

Despite its non-monotonic dips in strength, the C-F magnetic field still varies at two periods per one longitude revolution, consistent with Jovian half-synodic period expected from magnetopause plasma-magnetic conditions. The C-F field magnitudes in Figure 48 can be directly added to or subtracted from each other as the field is always parallel/antiparallel to GphiO Z-axis, obtaining maximum variation $|\Delta\mathbf{B}_{\text{CF}}| = -18.61 \text{ nT}$ between $\lambda_{\text{III}} = 268^\circ$ and $\lambda_{\text{III}} = 334^\circ$, assuming the magnetopause standoff distance remains fixed during the transition. In reality, the boundary's movement is unknown, complicated by intermingling influences of the acting pressure and energy-transfer interactions like magnetic reconnection.

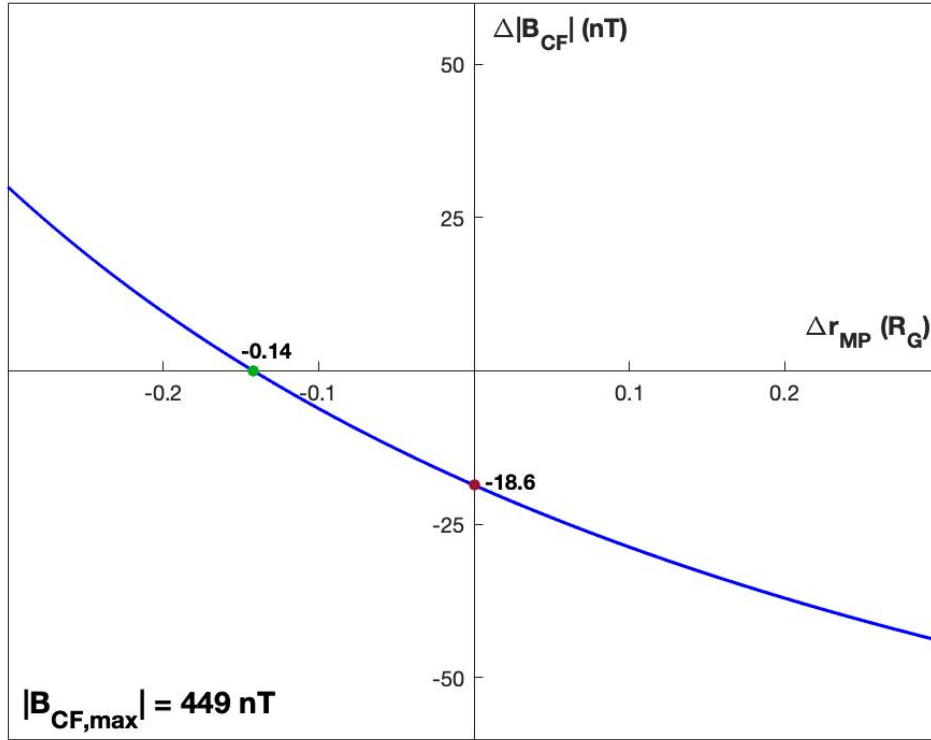


Figure 50: Change in C-F magnetic field strength at Ganymede's equatorial surface as the moon transitions from maximum $\lambda_{III} = 268^\circ$ to minimum $\lambda_{III} = 334^\circ$. The maximum C-F field strength is $|B_{CF,max}| = 449$ nT at $r_{MP} = 1.65 R_G$, shown in the bottom-left corner. The horizontal axis indicates change in magnetopause standoff distance during the transition, where positive values correspond to outward boundary movement away from Ganymede. X and Y intercepts of the curve is marked in green and red respectively.

Magnetopause motions affect the standoff distance and hence the C-F magnetic field measured on Ganymede's surface.

The standoff distance is hence considered as an independent variable within range of $r_{MP} = 1.35 - 1.95 R_G$, centralised by the analytical model's standard $r_{MP} = 1.65 R_G$ with $0.001 R_G$ resolution. For the subsequent comparison, the C-F magnetic field strength at $\lambda_{III} = 268^\circ$ and $r_{MP} = 1.65 R_G$ is taken as the baseline value. Then, I calculate the C-F field strengths at $\lambda_{III} = 334^\circ$ across the prescribed range of standoff distances and compute their deviations from the baseline value. The result is shown in Figure 50 with the standoff distance along the horizontal axis and variance of the C-F magnetic field along the vertical axis. A positive (negative) change in r_{MP} indicates a magnetopause that is farther (closer) to Ganymede's surface than the baseline $r_{MP} = 1.65 R_G$.

If the upstream magnetopause moves outward during Ganymede's transition from $\lambda_{III} = 268^\circ$ to $\lambda_{III} = 334^\circ$, then the C-F field strength at the moon's surface becomes weaker at the end of the transition, therefore the variance $|\Delta \mathbf{B}_{CF}|$ will be more negative than if the magnetopause remains fixed in location. Conversely, an inward magnetopause motion will result in more positive value of $|\Delta \mathbf{B}_{CF}|$. The two trends are reflected together through the monotonic negative gradient seen in Figure 50. The y-intercept is at $|\Delta \mathbf{B}_{CF}| = -18.6$ nT consistent with previous finding. The x-intercept is at $\Delta r_{MP} = -0.14 R_G$, indicating that a slight inward boundary movement is required to negate the exciting magnetic field for this specific Ganymedeian transition.

Because the C-F magnetic field will always vary due to continuous changes in near-magnetopause plasma-magnetic conditions, it is implausible that boundary movement can exactly negate ocean excitation at all times, indicating very high likelihood for secondary induction at Jovian half-synodic period. Given that Ganymede completes one $\Delta \lambda_{III} = 360^\circ$ revolution in 10.53 hours, timeseries of $|\Delta \mathbf{B}_{CF}|$ can be constructed from Figure 48 for any specified starting System-III Jovian east longitude, producing a fixed-boundary form of temporal exciting field function ($\mathbf{B}_0(t)$ in Eq. 6.15). However, accounting for potential boundary motion will be highly complicated, limiting the usefulness of the timeseries. Therefore, we focus on the upper limits of C-F magnetic field variance shown in Figure 50 as the probe into subsurface ocean induction.

Eq. 6.15 asserts that the inductive response from an infinitely conducting ocean will be half of the exciting field's magnitude. Taking $|\Delta \mathbf{B}_{CF}| = -18.6$ nT as example, I obtain $|\mathbf{B}_{ind}^\infty| = 9.3$ nT assuming $r_0 \sim R_G$ (i.e., ocean is close to Ganymede's surface). The fractional response between oceans of infinite and Earth's conductivities is estimated to be $\kappa = 0.84$ (Kivelson et al., 2002), reducing the maximum estimate of Ganymede's oceanic response to $|\mathbf{B}_{ind}| = 7.8$ nT. Given the approximative nature of my methods, the magnetic response of the ocean is generalised to magnitude order range $|\mathbf{B}_{ind}| \sim 1 - 10$ nT, which is well within measurement capability of modern magnetometers. Hence, the signal of Ganymedeian magnetic induction at Jovian half-synodic period should be detectable by future in-situ observations.

Previous to this analysis, magnetic contributions of Ganymede's current system have been studied through resistive MHD simulations (Olsen et al., 2010). The investigation yielded current-associated magnetic field strength ~ 50 nT during G28 Galileo upstream flyby of Ganymede²⁵. Comparison between this value and those in Figure 48 (~ 440 nT) is hampered by the fact that the flyby occurs significantly above and away from the moon's surface equator. Even so, the order of magnitude difference is significant. The discrepancy likely arises partly from inaccuracy in the equatorial ring current model, and from the standoff distance $r_{MP} = 1.65 R_G$ used in the analytical model which is less than the $r_{MP} \sim 2 R_G$ seen in the simulations. Although it is the C-F magnetic field variance that drives magnetic induction, the oscillation amplitude may be proportional to the overall field strength. If so, the inductive response magnitude may lie at the lower end of the $|B_{ind}|$ range above, but it should remain resolvable by future observations.

Narrowing the inductive response range to a single value (for a particular section of Ganymede's orbit) will require improved constraint on Ganymede's magnetopause motion. This can be achieved through adapting numerical simulations to track the magnetopause under variable Jovian plasma-magnetic conditions, which will utilise significant computational resources for relatively inefficient returns. Alternatively, a portion of future in-situ flybys/orbits can be used to finely monitor the boundary's movement. Only after the magnetopause oscillation is constrained that half-synodic response signals will be extractable from magnetometer data. Identifying magnetopause induction will then have to be separated from similar-period signals, including one corresponding to the second harmonic of Jupiter's rotation (Vance et al., 2021).

For the signals to yield insight into the subsurface ocean structure, their associated isolines in the depth-conductivity contour chart must also be calculated. The process will involve direct evaluation of the induction equation (Eq. 6.5) for multiple finitely conductive layers (ionosphere and ocean) model that includes proper characterisation of the C-F magnetic field. Given that a scheme for this task is already available for Ganymede (Vance et al., 2021),

²⁵ The study did not separate the C-F current' contribution from others in the system (e.g., field-aligned currents), but effects from the C-F current is expected to be dominant.

computation of half-synodic isolines presents a logical next step for research on Ganymede's induction response to magnetopause conditions.

Over the last four chapters, I have provided my scientific contributions on Ganymede's magnetic environment, ranging from assessments of magnetopause energy-transfer interactions to the boundary's impacts on both the magnetosphere and the subsurface ocean. My research has largely utilised analytical modelling, aided by limited observations from the Galileo spacecraft. Fortunately, in-situ data availability on Ganymede is expected to greatly increase with advent of new space missions, including an orbiter specifically targeting the largest Jovian moon. In the next and final chapter of my thesis, I will contextualise my research in relation to existing knowledge of Ganymede and discuss how additional observations can revolutionise our understanding of the largest Jovian moon.

Chapter 7: The Future of Ganymede's Research

7.1. Summary of State

Since the discovery of Ganymede's permanent internal magnetic field, the moon magnetosphere has been subjected to numerous data-based and simulation studies, aimed to develop our knowledge on one of the Solar System's most unique magnetic systems. With every passing year, what we know of the magnetosphere increases in detail, from the Alfvén wing structure to the Dungey-like plasma convection and the almost-certain presence of the subsurface liquid ocean.

Yet, while many fundamental questions have been answered, many more remain unfulfilled. Magnetospheric dynamics are thought to be driven by plasma-magnetic interactions in magnetic reconnection and Kelvin-Helmholtz (K-H) instability on the upstream magnetopause. However, little has been predicted of their natures beyond existence, and their influences on the magnetosphere are largely uncovered. Similarly, despite relative certainty of the subsurface ocean's presence, its structure and interactions with surrounding magnetic environment are not known.

Over the course of this thesis, my research seeks to characterise interactions and dynamics in Ganymede's magnetosphere through analytical modelling. Simpler parametrisations lead to less precise results, but they can be obtained using significantly smaller computational resources. Consequently, the analytical model allows evaluation of more fundamental kinetic plasma theory that can only be captured indirectly under current numerical code schemes. Findings from analytical and numerical modelling are therefore complementary, and together they can provide a coherent, theoretically supported picture of Ganymede's magnetosphere.

Chapter 3 provides the first analytical assessment of magnetic reconnection on Ganymede's upstream magnetopause. Reconnection onset conditions are shown to be comfortably satisfied across the entire (closed field) magnetopause at all times, supporting existing prediction of widespread reconnection events from resistive magnetohydrodynamic (MHD) simulations. The associated reconnection electric field is also shown to be large, indicating

that the reconnection structure cannot be constrained at fundamental level. But most importantly, the Ganymede reconnection rate is shown to exhibit variation correlated to Jupiter's rotation, establishing a prediction of direct planetary influence on the moon's magnetosphere. My findings on reconnection onset and rate have been successfully published in *Geophysical Research Letters* (Kaweeyanun et al., 2020).

Chapter 4 then utilises magnetic flux transport to further investigate upstream magnetic reconnection and its impacts on Ganymede's magnetosphere. The survey of transient flux-transfer events (FTEs) provides the first estimate on localisation of reconnection. FTEs are shown to supply a significant, but not dominant, proportion of opened upstream magnetic flux. The finding supports evidence of FTEs' existence from numerical simulations, and shapes expectation for future reconnection observations. The chapter also offers a brief glimpse into Ganymede's scanty known downstream reconnection, showing that its rate must be rapid to maintain a steady-state magnetosphere. A case study of orderly flux transport then illustrates that equal upstream-downstream reconnection will produce Jovian half-synodic variation in Ganymede's aurora latitude, a result that also appears consistent with numerical simulations.

Chapter 5 returns to the upstream magnetopause, providing the first analytical assessment of K-H instability on the boundary. The instability's linear phase onset is found to be viable along Ganymede's magnetopause flanks at all times, forming K-H waves which then propagate downstream parallel to the external Jovian plasma flow. Then, I propose the first picture of the kinetic finite Larmor radius (FLR) effect, in which local ion gyromotion notably enhances K-H instability growth on Ganymede's sub-Jovian magnetopause flank. Finally, the potential for nonlinear K-H vortex growth is evaluated in context of concurring magnetic reconnection. Under strongly antiparallel magnetopause magnetic field configuration, reconnection is expected to disrupt vortex formation, and ultimately suppresses K-H instability's role in cross-boundary plasma and energy transport. The assessment in this chapter has been published in *Journal of Geophysical Research: Space Physics* (Kaweeyanun et al., 2021).

Lastly, Chapter 6 considers magnetic induction by Ganymede's subsurface ocean in response to Jovian half-synodic variation in magnetopause conditions. The study seeks to establish the

Chapman-Ferraro (C-F) magnetic field as the first induction source that arises from Ganymede's internal magnetism. The C-F magnetic signal is shown to reach the subsurface ocean without penetrating through the liquid layer, leading to response whose magnitude depends on yet-unconstrained magnetopause movement, but lies comfortably within resolution range of modern magnetometers. As the result, Ganymede's magnetopause may prove a viable tool for future probing of the ocean's structure and effects.

There is still much to learn about Ganymede's magnetosphere. On the upstream side, the magnetopause oscillation, and thus temporal evolution of magnetospheric compression, remains unconstrained. Little is known about Ganymede's radiation belts beyond their expected existence. Local plasma convection patterns can also be more conclusively established, alongside aurora structures whose interactions with Jupiter are much better studied from the side of Jovian polar caps. On the downstream side, Ganymede's equatorial plasma sheet and downstream reconnection site(s) are currently uncharacterised. Study of downstream reconnection will offer insight into plasma loss processes. This, alongside an evaluation of particle entry and exit through the Alfvén wings, will define the difference between Dungey cycles in super-Alfvénic and sub-Alfvénic environments, the latter of which is unique to Ganymede as of this writing.

Given the estimative nature of my analytical methods, research in the work chapters can be advanced through development of numerical simulations. As computation becomes increasingly powerful, global kinetic simulations becomes more common, allowing particle-level physics to be captured directly without idealised approximations. More efficient models may also allow replacement of steady-state upstream Jovian plasma sheet conditions with variable ones, facilitating better simulation of temporal evolutions experienced by Ganymede in the real environment. Such developments will bridge the gaps between numerical and analytical modelling methods, which will not only strengthen the conclusions from the work chapters, but also answer some of the unknowns discussed in the previous paragraph.

But at the fundamental level, research on Ganymede's magnetic environment is limited by the absence of in-situ observations. Spacecraft data is crucial for verification of results in this thesis, which until then are merely predictions and/or expectations. Including Galileo, the

Ganymede magnetosphere has received six flybys that constitute all available data points. This is insufficient for constraint of basic magnetospheric structure, let alone measurement of temporal dynamics. Hence, our future knowledge of Ganymede correlates to development of space missions, their trajectory plans, and their instruments which will obtain the necessary data.

7.2. JUPITER ICY MOON EXPLORER (JUICE)

As mentioned in Chapter 2, the Juno extended mission provides one flyby of Ganymede on 7 June 2021. The spacecraft entered the moon's magnetosphere from the downstream side and later exited through the upstream magnetopause. Findings from this flyby are currently being published in geophysical journals. So far, Juno's particle instruments have identified whistler-mode waves and electron harmonic emissions in the magnetosphere (Kurth et al., 2022), measured new density of the ionosphere (Buccino et al., 2022), and valued energetic particle fluxes in close proximity of Ganymede (Paranicas et al., 2022). The spacecraft's magnetometer results are currently under peer review and expected to observe magnetopause interactions including magnetic reconnection and/or K-H instability.

While significant, the Juno mission's influence will pale in comparison to that of JUPITER ICY MOON EXPLORER (JUICE) mission, which is set to be launched by the European Space Agency (ESA) in 2023. The JUICE mission possesses two objectives – to understand the habitability potential of Jupiter's icy moons, and to understand the Jovian plasma-magnetic environment and its interactions with the moons. As the sole magnetic moon with a subsurface ocean, Ganymede lies at the centre of both goals, and is therefore the prominent focus of the mission. The timetable of JUICE is listed in Table G. Upon the probe's arrival at Jupiter in 2030, JUICE will first conduct a Jupiter Tour (2030-2032), orbiting Jupiter along the equator and at high latitudes while making flybys of the Galilean moons including Ganymede. In 2032, JUICE will begin insertion into Ganymede's orbit, becoming the first spacecraft to orbit a non-terrestrial moon. The Ganymede Tour is expected to last until approximately 2034, after which JUICE will crash into the moon and terminate the mission.

Table G: Timeline of the JUICE Mission (yr = year, mo = month, d = day)

Phase	Duration	Estimated Year
1. Interplanetary Cruise	8.25 yr	2023-2031
Jupiter Tour		
2. Jupiter's equatorial region (I)	1 yr	2031
3. Europa flybys	36 d	2031
4. Jupiter's high latitude, Callisto flybys	6 mo	2032
5. Jupiter's equatorial region (II), Ganymede flybys	9 mo	2032-2033
Ganymede Tour		
6. Elliptical orbits (I)	30 d	2033
7. High altitude (5,000 km) orbit	90 d	2033
8. Elliptical orbits (II)	30 d	2033
9. Medium altitude (500 km) orbit	102 d	2033-2034
10. Low altitude (200 km) orbit	30 d	2034

Note: Phase timeline information obtained from Grasset et al., 2013 with years adjusted due to delays from the Covid-19 pandemic (Witasse, 2022). The original launch was scheduled for June 2022.

A schematic image of JUICE is shown in Figure 51 with instruments labelled. Furthest away from the spacecraft's centre, attached to an extended boom apparatus, is the magnetometer (J-MAG) for magnetic field strength and orientation observations. The boom is required to prevent contamination of magnetic fields generated by currents inside the other instruments. At peak performance, the magnetometer can resolve magnetic signals down to 0.01 nT. However, due to particle drifts and offsets, the effective resolution of J-MAG is ~ 1 nT (Patrick Brown, personal communication, February 7, 2022).

Complementing the magnetometer are plasma sensors (PEP – Particle Environment Package) used to characterise plasma environment (ESA, 2019), alongside atmospheric ice (MAJIS) and ultraviolet (UV) spectrometers. On the waves' front, JUICE possess radio instruments for explorations of plasma (RPWI), atmosphere and surface (SWI), gravity field (3GM), and subsurface ocean (RIME) environments. Finally, Ganymede's surface can be mapped by a laser altimeter (GALA) and an optical camera (JANUS), the latter of which has 2.4 m resolution.

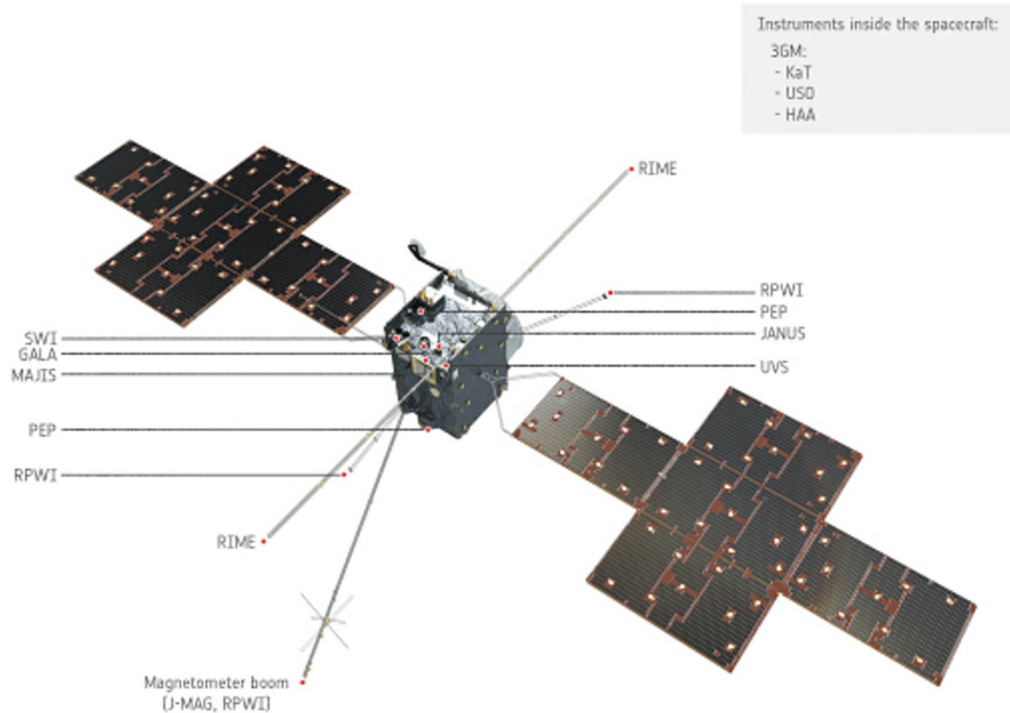


Figure 51: An illustration of Jupiter ICy moon Explorer (JUICE) and its scientific instruments. Image courtesy of ESA, (2019).

The instruments equip JUICE to target a wide array of conditions in and around Ganymede, especially the outer ice shell and the liquid ocean it contains.

In each phase of its trajectory near Ganymede, JUICE targets different aspects of the moon and its magnetosphere (Grasset et al., 2013). Before the orbital insertion, the spacecraft will make flybys of Ganymede in Jupiter's equatorial region (Phase 5). During this time, there will be multiple crossings of the upstream magnetopause, allowing observations of boundary interactions including magnetic reconnection and K-H instability. Orbital insertion will begin with highly elliptical orbits (Phase 6), during which Ganymede will be geologically mapped and its local plasma-magnetic environment will be studied. The magnetosphere will be defined in size dimension. JUICE will also traverse Ganymede's polar caps, studying auroras and particle precipitation along open magnetic flux tubes that connect between Ganymede and Jupiter. These investigations will also take place when JUICE briefly return to elliptical orbits (Phase 8).

The spacecraft's transition to high-altitude (5,000 km) circular orbits (Phase 7) will allow JUICE to focus on global mapping of Ganymede's geological features and surface composition.

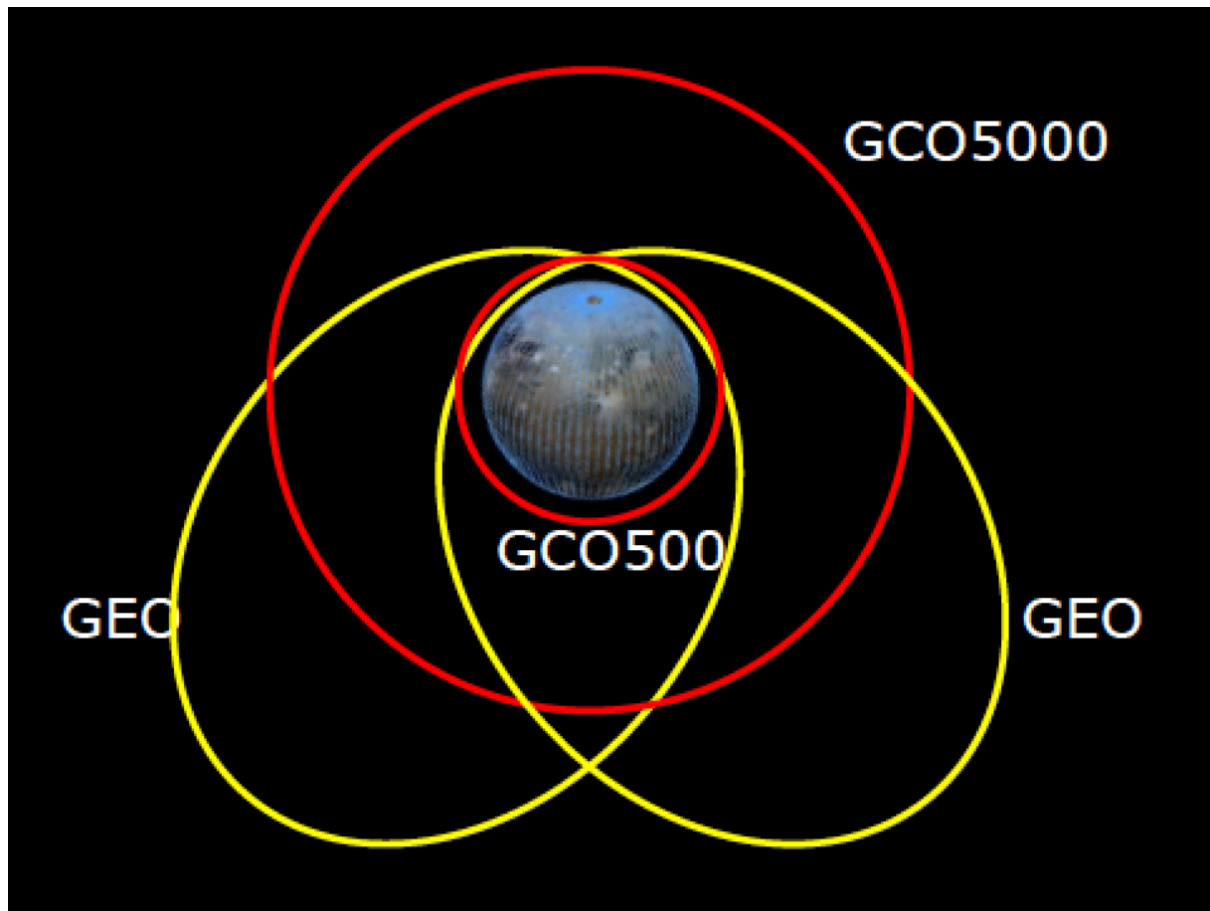


Figure 52: An outline of JUICE's elliptical (GEO, yellow), high-altitude orbit (GCO5000, red), and medium-altitude orbit (GCO500, red) around Ganymede. Image courtesy of Vallat & Witasse, (2018), *Lunar Planet Institute, University Space Research Association*.

Later, medium-altitude (500 km) circular orbits (Phase 9) will provide necessary measurements to resolve the inductive effects from the subsurface ocean. Finally, low-altitude (200 km) circular orbits (Phase 10) will offer detailed constraint on Ganymede's gravity fields, in-situ measurements of its upper atmosphere, and high-resolution imaging of its surface. High- and medium-altitude orbits for JUICE are illustrated in Figure 52.

Note that these trajectory phases are subjected to change prior to the mission's launch. Nevertheless, JUICE shall greatly expand the size of observational samples available on Ganymede and its magnetic environment. The number of magnetopause crossings in Phase 5 is expected to reach double digits as JUICE enters Ganymede's orbit. Then, each elliptical and circular orbit is expected to take ~12 hours, leading to over 400 expected journeys around the moon.

With such data availability, global structures (or lack thereof) for magnetic reconnection (X-line and FTEs) and K-H instability (waves and vortices) can be comfortably established on the upstream magnetopause. The open-closed field boundaries (OCFBs) can be tracked throughout its oscillation (if present), giving a full picture of magnetic flux transport and plasma convection in Ganymede's magnetosphere (aided by direct plasma measurements across the magnetic domain). Following constraint of the magnetopause motion, late-mission orbits will allow measurement of near-surface Ganymede magnetic field, and subsequently separation of multi-period induced magnetic signals over extended time intervals. In all, the information from JUICE will verify the predictions from my research, providing conclusion to the work chapters of this thesis.

We are still in the early days of Ganymede research. With the likely existence of a subsurface liquid ocean, Ganymede harbours possibility for present and future life. It is this potential that powers JUICE and will drive future research on the Solar System's sole internally magnetic moon. As demonstrated through induction, magnetism and subsurface ocean are intertwined, each providing a vital probe toward understanding of the other. Increasingly detailed measurement of Ganymede's magnetic field will yield insights into the ocean's structure, allowing improved definition of its inductive responses and further resolution of the magnetic field.

Ganymede's magnetic field protects the moon from Jupiter's intensive plasma radiation. Understanding of Ganymede's nearby space environment will offer protection for JUICE and other future spacecraft, which may investigate the moon not only for scientific knowledge, but also as a prospective destination for human habitation. If that objective is ever proposed, Ganymede's magnetic field will likely receive attention not dissimilar to Earth's, where space weather is intensively studied. But even without such lofty goal, practical mission considerations will afford attention on the moon's magnetosphere in many years to come.

My work is a small chapter in the long history of Ganymede exploration. The last four years of investigation have been highly rewarding, and taught me so much about a moon that is unlike any other in the Solar System. With its size and magnetism, Ganymede is in many ways comparable to planets rather than moons. I hope my work will prove useful among

researchers of the magnetosphere, and offer insights needed for its future exploration. The road to JUICE will take many years. But with so much still to be learned, the intervening years will not be without activity. I look forward to reading the next chapter of Ganymede, a special satellite that has etched its way into my heart.

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16	Figure 2a: Image of Ganymede from Pioneer 10	Frieden, B.R. & Swindell, W. (1976). Restored Pictures of Ganymede, Moon of Jupiter. <i>Science</i> , 191 (4233), 1237-1241. https://www.jstor.org/stable/1741143	© Science	27 May 2022	Print license available
16	Figure 2b: Image of Ganymede from Voyager 1	Jónsson, B. (1979). Ganymede from Voyager 1. <i>The Planetary Society (NASA-JPL Photojournal)</i> . https://www.planetary.org/space-images/ganymede-from-voyager-1	© Bjohn Jónsson, The Planetary Society	N/A	Permission obtained under Creative Commons Attribution-NonCommercial-NoDerivs 3.0 Unported (CC BY-NC-ND 3.0) license
16	Figure 2c: Voyager 2 image of Ganymede's Nipper Sulcus	Williams, D. & Friedlander, J. (2015a). Close-up of Nipper Sulcus region on Jupiter's moon Ganymede. <i>Ganymede – Voyager 2, NASA</i> . https://nssdc.gsfc.nasa.gov/imgcat/html/object_page/vg2_2063525.html	© NASA	N/A	Image under public domain
16	Figure 2d: Voyager 2 image of Ganymede's Lagash Sulcus	Williams, D. & Friedlander, J. (2015b). Lagash Sulcus and other furrows and ridges on Ganymede. <i>Ganymede –</i>	© NASA	N/A	Image under public domain

		<i>Voyager 2</i> , NASA. https://nssdc.gsfc.nasa.gov/imgcat/html/object_page/vg2_2063729.html			
17	Figure 3: Schematic diagram of scientific instruments aboard Galileo	National Aeronautics and Space Administration (NASA). (2004). <i>Galileo Diagram. Solar System Exploration Multimedia</i> , NASA. https://web.archive.org/web/20041102173235/http://solarsystem.nasa.gov/multimedia/display.cfm?IM_ID=1945	© NASA	N/A	Image under public domain
18	Figure 4: Galileo's magnetometer observation of Ganymede during the G1 flyby.	Kivelson, M.G., Khurana, K.K., Russell, C.T., Walker, R.J., Warnecke, J., Coroniti, F.V., Polanskey, C., Southwood, D.J., Schubert, G. (1996). Discovery of Ganymede's magnetic field by the Galileo spacecraft. <i>Nature</i> , 384 (6609), 537-541. http://doi.org/10.1038/384537a0	© Nature	30/06/2022	Print license available
22	Figure 6: Illustration of Earth's dipole magnetic field	(1st ver.) University Corporation of Atmospheric Research (UCAR). n.d. Earth's Dipole Magnetic Field. <i>Center for Science Education</i> , UCAR. https://scied.ucar.edu/image/earths-dipole-magnetic-field (2nd ver.) Reid, P. (2009). <i>The Earth's</i>	(1st ver.) © UCAR (2nd ver.) Peter Reid	N/A	(1st ver.) Permission request email not returned. Image replaced. (2nd ver.) Permission granted by copyright holder (email).

		Magnetic Field. <i>University of Edinburgh, NASA.</i> https://www.nasa.gov/mission_pages/sunearth/news/gallery/Earths-magneticfieldlines-dipole.html			
24	Figure 7a: The Helix Nebula	National Aeronautics and Space Administration (NASA), Space Telescope Science Institute (STScI), European Space Agency (ESA), National Radio Astronomy Observatory (NRAO). (2020). Helix Nebula, <i>NASA</i> . https://www.nasa.gov/mission_pages/chandra/images/helix-nebula.html	© NASA, STScI, ESA, NRAO	N/A	Image under public domain
24	Figure 7b: Argon plasma lamp	Kummer, J. (2016). Argon – Chemical Elements: A Virtual Museum. <i>Jumk.de Webprojects</i> . https://images-of-elements.com/argon.php	© Jürgen Kummer	N/A	Permission obtained under Creative Commons Attribution-NonCommercial-NoDerivs 3.0 Unported (CC BY-NC-ND 3.0) license
24	Figure 7c: Coronal mass ejection event from the Sun	Gleber, M. (2014). CME Week: The Difference Between Flares and CMEs. <i>NASA</i> . https://www.nasa.gov/content/godda rd/the-difference-between-flares-and-cmes	© NASA	N/A	Image under public domain
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32	Figure 9a-d: Guiding centre drifts under uniform magnetic field, ExB drift, gravitational drift, and magnetic gradient drift.	Wikipedia. (2007). Guiding center. <i>Wikipedia – the Free Encyclopedia</i> . https://en.wikipedia.org/wiki/Guiding_center	© Stannered, Wikipedia	N/A	Permission obtained under Creative Commons 2.5 Generic (CC BY 2.5) license
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		http://silas.psfc.mit.edu/introplasma/chap2.html			
40	Figure 10: Diagram of magnetic reconnection	Eastwood, J.P., Phan, T.D., Øieroset, M., Shay, M.A., Malakit, K., Swisdak, M., Drake, J.F., Masters, A. (2013). Influence of asymmetries and guide fields on the magnetic reconnection diffusion region in collisionless space plasmas. <i>Plasma Phys. Control. Fusion</i> , 55 (12), 1-12. http://doi.org/10.1088/0741-3335/55/12/124001	© Institute of Physics (IOP) Publishing	N/A	Permission obtained under Creative Commons Attribution 3.0 Unported (CC BY 3.0) license
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47	Figure 13: A schematic diagram of Earth's magnetosphere.	Eastwood, J.P., Hietala, H., Tóth, G., Phan, T.D., Fujimoto, M. (2015). What Controls the Structure and Dynamics of Earth's Magnetosphere? <i>Space Sci. Rev.</i> , 188 (1), 251-286. https://doi.org/10.1007/s11214-014-0050-x	© Space Science Review	N/A	Permission obtained under Creative Commons (CC BY) license – open access article
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		512. http://doi.org/10.1038/384511a0			
57	Figure 16: Ganymede's magnetosphere from side and inflow views	Jia, X., Kivelson, M.G., Khurana, K.K., Walker W.J. (2010). Magnetic Fields of the Satellites of Jupiter and Saturn. <i>Space Sci. Rev.</i> , 152 (1), 271-305. http://doi.org/10.1007/s11214-009-9507-8	© Space Science Review	30/06/2022	Print license available
61	Figure 17: Images of Ganymede's aurora taken by the Hubble Space Telescope from	McGrath, M.A., Jia, X., Retherford, K.D., Feldman, P.D., Strobel, D.F., Saur, J. (2013). Aurora on Ganymede. <i>J. Geophys. Res. Space Phys.</i> , 118 (5), 2043-2054. https://doi.org/10.1002/jgra.50122	© Journal of Geophysical Research: Space Physics	30/06/2022	Print license available
63	Figure 18: Comparison between Galileo G8 flyby data and MHD-EPIC simulation results	Tóth, G., Jia, X., Markidis, S., Peng, I.B., Chen, Y., Daldorff, L.K.S., Tenishev, V.M., Boronikov, D., Haiducek, J.D., Gombosi, T.I., Gloer, A., Dorelli, J.C. (2016). Extended magnetohydrodynamics with embedded particle-in-cell simulation of Ganymede's magnetosphere. <i>J. Geophys. Res. Space Phys.</i> , 121 (2), 1273-1293. https://doi.org/10.1002/2015JA021997	© Journal of Geophysical Research: Space Physics	30/06/2022	Print license available
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	reconnection diffusion region	Earth Magnetic Reconnection from In-Situ Observations. <i>Space Sci. Rev.</i> , 160 (1), 95-121. https://doi.org/10.1007/s11214-011-9820-x			
78	Figure 22: Magnetoplasma conditions near idealised Ganymede's magnetopause	Kaweeyanun, N., Masters., A., Jia, X. (2021). Analytical Assessment of Kelvin-Helmholtz Instability Growth at Ganymede's Upstream Magnetopause. <i>J. Geophys. Res. Space Phys.</i> , 126 (8), e2021JA029338. https://doi.org/10.1029/2021JA029338	© Journal of Geophysical Research: Space Physics	30/06/2022	Print license available. Image originally published by myself.
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93	Figure 26: Ganymede's magnetospheric field lines	Jia, X., Walker, R.J., Kivelson, M.G., Khurana, K.K., Linker, J.A. (2008). Three-dimensional MHD simulations of Ganymede's magnetosphere. <i>J. Geophys. Res.</i> , 113 (A6), A06212. https://doi.org/10.1029/2007JA012748	© Journal of Geophysical Research: Space Physics	30/06/2022	Print license available

94	Figure 27: Half-open magnetic flux surface	Zhou, H., Tóth, G., Jia, X., Chen, Y. (2020). Reconnection-Driven Dynamics at Ganymede's Upstream Magnetosphere: 3-D Global Hall MHD and MHD-EPIC Simulations. <i>J. Geophys. Res. Space Phys.</i> , 125 (8), e2020JA028162. https://doi.org/10.1029/2020JA028162	© Journal of Geophysical Research: Space Physics	30/06/2022	Print license available
95	Figure 28: Three popular models of flux-transfer events formation	Paschmann, G., Øieroset, M., Phan, T.D. (2013). In-Situ Observations of Reconnection in Space. <i>Space Sci. Rev.</i> , 178 (2), 385-417. https://doi.org/10.1007/s11214-012-9957-2	© Space Science Review	30/06/2022	Print license available
97	Figure 29a: Flux rope/tube from resistive MHD model	Jia, X., Walker, R.J., Kivelson, M.G., Khurana, K.K., Linker, J.A. (2010). Dynamics of Ganymede's magnetopause: Intermittent reconnection under steady external conditions. <i>J. Geophys. Res.</i> , 115 (A12), A12202. https://doi.org/10.1029/2010JA015771	© Journal of Geophysical Research: Space Physics	30/06/2022	Print license available
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		at Ganymede's Upstream Magnetopause. <i>J. Geophys. Res. Space Phys.</i> , 126 (8), e2021JA029338. https://doi.org/10.1029/2021JA029338			
125	Figure 40: Linear Kelvin-Helmholtz instability assessment on Ganymede's magnetopause	Kaweeyanun, N., Masters., A., Jia, X. (2021). Analytical Assessment of Kelvin-Helmholtz Instability Growth at Ganymede's Upstream Magnetopause. <i>J. Geophys. Res. Space Phys.</i> , 126 (8), e2021JA029338. https://doi.org/10.1029/2021JA029338	© Journal of Geophysical Research: Space Physics	30/06/2022	Print license available. Image originally published by myself.
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137	Figure 42: Equatorial current model for Chapman-Ferraro magnetic field	Glassmeier, K.-H., Auster, H.-U., Motschmann, U. (2007). A feedback dynamo generating Mercury's magnetic field. <i>Geophys. Res. Lett.</i> , 34 (22), L22201. https://doi.org/10.1029/2006GL028001	© Geophysical Research Letters	30/06/2022	Print license available

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160	Figure 51: JUICE and its scientific instruments	European Space Agency (2019). JUICE's instruments. <i>ESA Science & Technology</i> . https://sci.esa.int/web/juice/-/61110-juice-instruments	© ESA	N/A	Image under public domain
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I remember the excitement when Galileo detected Ganymede's magnetic field; it's a fascinating structure, with its little magnetosphere-inside-another-magnetosphere 😊...

Best of luck,
Peter

Peter Reid
Planning & Communications Project Manager
College of Science and Engineering
The University of Edinburgh
Murchison House
10 Max Born Crescent
Edinburgh EH9 3BF

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Asst. Dir. for Planetary Science
Laboratory for Atmospheric & Space Physics
University of Colorado
Fran.Bagenal@lasp.colorado.edu

On Jul 8, 2022, at 9:13 AM, Laura Bloom <Laura.Bloom@lasp.colorado.edu> wrote:

Hi Nawapat,

I am forwarding your email to Fran Bagenal who should be able to help you.

Take care,

Laura Bloom
LASP/University of Colorado
3665 Discovery Drive
BOULDER, CO 80303
Phone: 303-492-6827
Fax: 303-492-1132
laura.bloom@lasp.colorado.edu

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Figure 15

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Figures 22, 39, 40, 41

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Figure 28

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Figure 30

24/10/2022, 15:03

Re: Update on FTE Magnetic Flux Content Collab

Hongyang Zhou <hyzhou@umich.edu>

Fri 2/4/2022 12:47 PM

To: Kaweeyanun, Nawapat <nawapat.kaweeyanun14@imperial.ac.uk>

This email from hyzhou@umich.edu originates from outside Imperial. Do not click on links and attachments unless you recognise the sender. If you trust the sender, add them to your [safe senders list](#) to disable email stamping for this address.

Hi Nawapat,

It's a tough time for all of us, but I hope you are doing well!

Feel free to use any image I provide that will be useful for you. I'm glad to see your research progress at pace! Enjoy writing your thesis and I am glad to read your insights into Ganymede's story.

Best,
Hongyang

On Fri, Feb 4, 2022 at 12:36 PM Kaweeyanun, Nawapat
<nawapat.kaweeyanun14@imperial.ac.uk> wrote:

Dear Hongyang,

I hope life has been treating you well.

I would like to apologise for not getting in touch sooner, and for missing a chance to respond to your query from our last correspondence (the Europlanet Society presentation was a talk, thank you).

I am writing to let you know that since our last correspondence, my research has taken me away from Ganymede's FTE flux content research. However, I am now writing up our work for my Imperial College PhD thesis, and will be including some of the images you very kindly sent me last year. Please let me know if you are alright with this. I would be quite happy to send you the relevant thesis chapter when it is completed (if you are interested, of course).

Looking forward to hearing from you.

Best Wishes
Nawapat

--

Hongyang Zhou
Dept. of Climate and Space Sciences and Engineering
University of Michigan, Ann Arbor, MI 48109-2143

Figure 37

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Figure 42

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Figure 52

24/10/2022, 15:03

RE: Image Permission (Thesis)

Olivier Witasse <Olivier.Witasse@esa.int>

Thu 7/21/2022 8:50 PM

To: Kaweeyanun, Nawapat <nawapat.kaweeyanun14@imperial.ac.uk>

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Hello

No problems.

You can put courtesy ESA.

Good luck with your PhD, send me the final document if it's possible!

Olivier

From: Kaweeyanun, Nawapat <nawapat.kaweeyanun14@imperial.ac.uk>

Sent: 21 July 2022 18:10

To: Olivier Witasse <Olivier.Witasse@esa.int>

Subject: Image Permission (Thesis)

Dear Olivier Witasse,

I am a final year Ph.D. student at Imperial College London, conducting research on the magnetic environment of Ganymede. As part of my thesis draft, I have included a graphic illustrating orbits of the JUICE spacecraft (attached), which appeared under your name in a presentation made available online. Given that my thesis will be publicly published, I am writing to seek your permission for my inclusion of this figure.

I hope this is the appropriate channel for my enquiry. If not, can you let me know the copyright holder whom I should contact?

Thank you in advance for your help.

Best Wishes,

Nawapat Kaweeyanun

Imperial College London