Plasma Transport Driven by the Rayleigh-Taylor Instability

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Plasma transport driven by the Rayleigh-Taylor instability

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Abstract Two important differences between the giant magnetospheres (i.e., Jupiter’s and Saturn’s magnetospheres) and the terrestrial magnetosphere are the internal plasma sources and the fast planetary rotation. Thus, there must be a radially outward flow to transport the plasma to avoid infinite accumulation of plasma. This radial outflow also carries the magnetic flux away from the inner magnetosphere due to the frozen-in condition. As such, there also must be a radial inward flow to refill the magnetic flux in the inner magnetosphere. Due to the similarity between Rayleigh-Taylor (RT) instability and the centrifugal instability, we use a three-dimensional RT instability to demonstrate that an interchange instability can form a convection flow pattern, locally twisting the magnetic flux, consequently forming a pair of high-latitude reconnection sites. This process exchanges a part of the flux tube, thereby transporting the plasma radially outward without requiring significant latitudinal convection of magnetic flux in the ionosphere.

1. Introduction

The giant magnetospheres (i.e., Jupiter’s and Saturn’s magnetospheres) are characterized by several internal plasma sources and fast planetary rotation [Delamere et al., 2015b, and references therein]. A steady state magnetosphere model implies an outward radial transport of this plasma. Here the steady state magnetosphere model refers to an average of the magnetosphere during an interval that is much longer than the typical time scale for the plasma radial transport. The depletion of the inner magnetospheric flux due to this outward flow requires a magnetic flux refilling process to maintain the conservation of magnetic flux. This radial transport process is key to understanding the dynamics of the giant magnetospheres. For a stretched magnetic configuration (i.e., the middle and outer magnetospheres) [Achilleos et al., 2010], reconnection of closed flux can generate lower flux tube content and lower flux tube entropy compared to the ambient flux tube, allowing rapid inward motion due to the stability requirement [Erickson and Wolf, 1980; Vasyliunas, 1983; Birn et al., 2013; Delamere et al., 2015a]. For the inner magnetosphere, the strong dipolar field is not suitable for magnetic reconnection in the equatorial plane.

It is widely believed that the centrifugally driven interchange instability triggers radial plasma transport [Thomsen, 2013; Mauk et al., 2009; Achilleos et al., 2015, and references therein]. This instability is very similar to the Rayleigh-Taylor (RT) instability [Chandrasekhar, 1961; Bateman, 1978], which operates when the flux tube content or the flux tube entropy decreases radially [Southwood and Kivelson, 1987; Achilleos et al., 2015, and references therein].

It is estimated that Io generates neutral gas at a rate of 600–2600 kg s$^{-1}$ [Hill, 1979; Bagenal et al., 1997; Thomas et al., 2004; Bagenal and Delamere, 2011], of which about 50% is ionized and transported radially outward initially at a low speed (i.e., <100 m s$^{-1}$) [Russell and Huddleston, 2000; Delamere and Bagenal, 2003; Bagenal and Delamere, 2011]. Galileo observed several intermittent, short-duration (i.e., ~26 s) events of increased magnetic field magnitude (i.e., ~1%) and with no ion cyclotron waves near Io, which have been interpreted as narrow channels of fast (i.e., ~100 km s$^{-1}$) inward motion of tenuous plasma (i.e., ~100 cm$^{-3}$) [Thorne et al., 1997; Kivelson et al., 1997]. For Saturn, about 70–750 kg s$^{-1}$ of water vapor is ejected from Enceladus’ geysers and about 12–250 kg s$^{-1}$ of that is ionized and transported radially outward [Bagenal and Delamere, 2011; Delamere et al., 2015b]. So-called “injection” events have often been observed by Cassini in Saturn’s inner magnetosphere, characterized by a short interaction with a hot tenuous plasma embedded within an ambient cold and dense plasma [Krimigis et al., 2005; Mauk et al., 2009; Thomsen, 2013, and references therein]. The transition time for Cassini to cross from ambient plasma into injection events is about few seconds [Paranicas et al., 2016], indicating that injection events have a sharp boundary between the ambient plasma.
In this study, the “interchange instability” is referred to any instability that interchanges the position of fluid elements and decrease the total free energy (e.g., the convective instability, the Rayleigh-Taylor instability, and the ballooning instability). This can either be based on a local stability criterion, which is the most common application or on the interchange of entire flux tubes. The latter is more restrictive because it requires global knowledge including ionospheric conductance and is applicable only when M-I coupling is faster than the time scale of the instability. This paper illustrates that local interchange can force high-latitude reconnection, which can transport plasma radially outward without the convection of the flux tube foot points in the ionosphere. By definition, component reconnection generates localized parallel electric fields. We will demonstrate this process by using a RT instability in a simple three-dimensional slab configuration. The next section introduces the basic properties of the RT instability and discusses the difference between centrifugal instability in a realistic magnetospheric configuration and the simple RT instability in a slab configuration. The numerical simulation methods, results, and discussion will be presented in sections 4, and 5, respectively.

2. Rayleigh-Taylor Instability Onset Conditions

It is instructive to review the basic properties of the Rayleigh-Taylor (RT) instability by investigating the simplest case. Assuming an ideal incompressible magnetohydrodynamic (MHD) system, two homogenous regions with different density (left: \( \rho_1 \) and right: \( \rho_2 \)) are vertically separated by an infinitely thin boundary layer (see Figure 1). There is a gravitational acceleration, \( g \), in the horizontal direction toward the left, and the magnetic field components, \( \mathbf{B} \), are only in the vertical plane. Linear theory [Chandrasekhar, 1961; Bateman, 1978] gives the RT growth rate,

\[
q = \sqrt{gk\frac{\rho_2 - \rho_1}{\rho_2 + \rho_1} - \frac{2(\mathbf{k} \cdot \mathbf{B})^2}{\mu_0(\rho_1 + \rho_2)}} \tag{1}
\]

where \( k \) is the RT wave vector and \( \mu_0 \) is the vacuum permeability. Here the first term on the right-hand side of equation (1) indicates that the RT instability driven by the gravitational force operates when the right fluid is denser than the left (i.e., \( \rho_2 > \rho_1 \)) in a hydrodynamic system. The second term shows that the presence of a magnetic field along the wave vector \( \mathbf{k} \) suppresses the RT instability.
More generally, a hydrodynamic system is RT unstable as long as \( \mathbf{g} \cdot \nabla \rho < 0 \) for an incompressible system. For a compressible system with an assumption that the wavelength is much shorter than the pressure scale length, \( \lambda = |\rho(\nabla p)^{-1}| \), the system is RT unstable when \( \mathbf{g} \cdot \nabla s > 0 \), where \( s = p \rho^{-\gamma} \) is the specific entropy, and \( \gamma = 5/3 \) is the ratio of specific heats for the monatomic assumption [Bateman, 1978]. It is well known that some instabilities (e.g., the Kelvin-Helmholtz instability) can be suppressed if the wavelength is comparable to the width of the transition layer. However, linear analysis implies no threshold of the shortest wavelength caused by a finite width (or small gradient) of the transition layer for the inviscid hydrodynamic RT instability, but the growth rate decreases in such conditions [Chandrasekhar, 1961; Bateman, 1978]. Hence, the RT instability tends to operate on a short wavelength mode. The shortest wavelength can be limited by the geometry (i.e., a parallel magnetic field component) as well as the microscopic physics. For instance, it has been demonstrated that the RT mode with a short wavelength (i.e., the wavelength is comparable to the density gradient scale length, \( L_\rho \)) can be easily stabilized when finite Larmor radius effects are taken into account (i.e., ion Larmor radius, \( L_g \ll L_\rho \)) [Huba, 1996; Huba and Winske, 1998].

For the Jovian inner magnetosphere, it is expected that the RT instability is localized inside of the Io plasma torus (i.e., scale height of the torus \( \sim 1 R_J \)), because both the large density gradient and the centrifugal force are localized at low latitudes, and the magnetic field line motion at high latitudes is limited by the Pedersen conductivity. For simplicity we assume that the dipole field is perpendicular to the equatorial plane, and the wave vector of the RT mode is mainly along the azimuthal direction. In this three-dimensional configuration, a RT mode localized along the polar direction indicates that the domain mode has a field-aligned wave vector component. Based on our prior discussion, one can always find a RT unstable mode by decreasing the perpendicular wavelength (i.e., \( k \approx k_L \gg k_\parallel \)) even in the case where the magnetic field energy is much greater than the potential energy.

For an equilibrium with straight magnetic field lines and cold plasma (i.e., the environment of the plasma torus in Jupiter’s inner magnetosphere), the centrifugal instability operates when the centrifugal force opposes the flux content gradient. For a homogeneous magnetic field (e.g., our simulation configuration), the onset criterion is identical to the Rayleigh-Taylor instability [Achilleos et al., 2015]. In spite of the similarities between the RT instability and the centrifugal instability, we note that the presence of magnetic field curvature and nonzero steady state velocity for the centrifugal instability in the real magnetosphere may modify the detailed
dynamics. However, the low latitude vorticities in the nonlinear stage, and consequently high-latitude reconnection, as we will discuss in detail in the next sections, should be insensitive to these two aspects.

3. Numerical Methods

High-latitude double reconnection driven by the nonlinear Rayleigh-Taylor instability is demonstrated in this study by using a three-dimensional MHD simulation. The full set of normalized resistive MHD equations is numerically solved by using the leap-frogs scheme [Fletcher, 1991; Otto, 1990]. All length scales in the results presented below are measured in units of the width of the initial density gradient. Similarly velocities are measured in units of the Alfvén speed, and time is normalized to the Alfvén transit time for the width of the initial density gradient. The simulation is carried out in a rectangular volume with a size of \([-L_x, L_x] \times [0, L_y] \times [0, L_z]\). Here the \(x\) and \(z\) directions represent the radially inward direction and the northward direction, respectively. The \(y\) direction is determined by the right-hand rule. In this configuration, the dynamics significantly evolves in the \(x\) direction and will eventually be influenced by the size in the \(x\) direction. Thus, the dimension in the \(x\) direction is set to be large enough (i.e., \(L_x = 20\)) that the RT instability can still evolve several growth times (i.e., \(\sim q^{-1}\)) in the nonlinear stage before the perturbation is reflected by the boundary in the \(x\) direction. The dimension along the \(y\) direction (i.e., \(L_y = 2\)) determines the largest wavelength in the \(y\) direction, and the \(L_z = 71\) limits the largest wavelength along the background magnetic field (i.e., dipole field). The shortest wavelength is limited by the grid resolution. In this study we use uniform grids in the \(y\) and \(z\) directions with the grid separation of 0.1 and nonuniform grid in the \(x\) direction with a best resolution of 0.1.

In this geometry, the gravitational acceleration in the negative \(x\) direction, \(g = -g(x)\hat{x}\) is set to be zero in the vicinity of the boundaries in the \(x\) direction to simplify the boundary conditions, which is given by \(g(x) = g_0[\tanh(x+L_x)−\tanh(x−L_x)]\). The equilibrium profile used in the simulations is \(\rho(x) = 1 + \delta \rho \tanh(x), B_z = 1\), and the thermal pressure is determined by the force balance, which is \(p(x) = \int_z^{L_z} \rho(x')g(x') dx' + p_0\) (see Figure 1). Here \(g_0 = 0.25, L_x = 0.75L_x, \delta\rho = 0.6, \) and \(p_0 = 0.25\).

In this study the numerical model is only used for demonstrating the transport mechanism. This process should be viewed as a nonlinear preexisting instability (e.g., RT instability and centrifugal instability), triggering a secondary instability (i.e., magnetic reconnection). We emphasize that the secondary instability plays an important role for plasma transport. Each instability associates a set of typical time scales and length scales. For Jupiter, the preexisting instability has a typical length scale, \(L_0\), about one or two Jupiter radii, and a time scale \(\tau_0\) of order few days. The secondary instability, in contrast, has a typical length scale, \(L_s\), only of order few kilometers and a typical time scale, \(\tau_s\), only of order a few seconds. Hence, to numerically simulate a relatively realistic preexisting instability, we face the situation of either insufficiently resolving the secondary instability (if normalized by the first set of dimensions \(L_0\) and \(\tau_0\)), or the radial transport is smeared by numerical diffusion (if normalized by the second set of dimensions \(L_s\) and \(\tau_s\)). The presence of the nonlinear stage of the preexisting instability is evidently demonstrated by the often observed fast inward flow. We, therefore, set the initial condition to be unrealistically unstable for the preexisting instability.

The RT instability is triggered by a velocity perturbation localized in the \(z\) direction, which is given by \(\delta v = f(z)[\nabla \Phi(x, y) \times \hat{z}]\). Here the stream function is \(\Phi = \delta \nu_1 \cos(\pi y/L_y) \cosh^{-1}(x/d_s)\), and the localization function is \(f(z) = 0.5[\tanh((z + L_z)/d_z) − \tanh((z − L_z)/d_z)]\), where \(d_s = 2, L_z = 20, d_z = 2\), and \(\delta \nu_1 = 0.1\).

The simulation employs a current dependent resistivity using a threshold current density and a very small background resistivity. This choice ensures that the resistivity is very small almost everywhere except in the location where the relative electron and ion drift speeds are large [Ma et al., 2014]. We also carried out the simulation with a constant resistivity model, showing high-latitude reconnection sites as well, which indicates that the presence of magnetic reconnection is insensitive to the resistivity model.

It is beneficial to use open boundary conditions in the \(x\) direction (i.e., \(\partial_x = 0\)) as previously discussed. The boundary conditions in the \(y\) direction and \(z = 0\) are determined by symmetry properties of the MHD equations [Otto et al., 2007]. An artificial friction term, \(-v(z)\nu\), on the right-hand side of the momentum equation is applied close to the top boundary (i.e., \(z = L_z\)) to mimic the assumption that the magnetospheric magnetic field is tied to the perfectly conducting ionosphere (i.e., \(v(x, y, L_z) = 0\)) [Ma et al., 2014]. Here friction coefficient is given by \(\nu(z) = 1 + 0.5 \left\{ \tanh \left[ (z − L_z)/15 \right] − \tanh \left[ (z + L_z)/15 \right] \right\} \).
We use a shallow gradient (see Figure 1) because a large gradient of the friction term leads to a sharp kink of magnetic field associated with this gradient, which is equivalent to a thin current layer with a current mostly in an \( xy \) plane perpendicular to the background magnetic field. As such, the frozen-in condition can be broken by resistive diffusion in this plane. This processes can be referred to as “slippage,” because it is akin to the convection of magnetic flux in the ionosphere due to a finite Pedersen conductivity, and it does not involve antiparallel magnetic field components and the parallel electric fields usually associated with reconnection. Therefore, slippage is different from the high-latitude double reconnection process which is described in next section. To eliminate slippage, a shallow gradient of the friction term as well as a shallow gradient of the initial perturbation along the \( z \) direction is used in our simulation.

High-latitude reconnection in our model is component magnetic reconnection in three dimensions, being difficult to describe from a topological view. A good example is the three-dimensional Kelvin-Helmholtz instability [see Faganello et al., 2012, Figure 2]. In this study we use three different methods to demonstrate the presence of high-latitude reconnection, i.e., estimating the field-aligned potential difference, \( \phi \); examining the frozen-in condition by the definition; and examining the change of the flux tube content and flux tube entropy. Here the field-aligned potential difference, \( \phi = \int E_\parallel ds \), is the integral of the parallel electric field, \( E_\parallel \), along the magnetic field line. For the first method, Hesse and Schindler [1988] proved that for \( B \neq 0 \), magnetic reconnection with global effects occurs only if \( \phi \neq 0 \) on a measurable set of field lines in the diffusion region. Hence, the field-aligned potential difference, \( \phi \), is often used as a measurement of the three-dimensional component magnetic reconnection rate. For the second method, we follow the definition of the frozen-in condition, that is, two fluid elements originally along the one field line should remain on the same field line. The large displacement of the magnetic field line footprint at the equatorial plane attests to the violation of the frozen-in condition. The third method relies on the MHD conserved quantities, i.e., flux tube content, \( M = \int (\rho / B) ds \), and flux tube entropy, \( H = \int (\rho^{1/\gamma} / B) ds \), where the integral is taken along the magnetic field lines traced from the top boundary to the equatorial plane. It can be readily shown that the change of these quantities can only be due to the violation of the frozen-in condition through magnetic reconnection from the MHD perspective [Birn et al., 2006].

4. Nonlinear Rayleigh-Taylor Instability and High-Latitude Double Reconnection

Figure 2 (top) presents the plasma density (color index) and the in-plane bulk velocity (black arrows) in the equatorial plane at \( t=18 \), showing the convection flow pattern driven by the nonlinear RT instability. The outward flow (i.e., \( v_x < 0 \)) is associated with the heavy plasma and an enhanced \( B_z \) (not shown in the plot), while the inward flow (i.e., \( v_x > 0 \)) corresponds to the light plasma and diminished \( B_z \) (not shown in the plot). Recall that the inward flow near \( \rho = 0 \) is characterized by an increase of the magnetic field magnitude [Kivelson et al., 1997], while, in contrast, the inward flow in Saturn’s inner magnetosphere (i.e., the “injection” events) typically correspond to a decrease of the magnetic field magnitude [Hill et al., 2005]. The differences are likely due to different profiles of the flux tube entropy, which requires an investigation using a more realistic configuration. We note that in this specific case the outward flow speed is comparable to the inward flow speed, being contrary to the well-known observational feature. However, this is because we use a quasi-eigenmode perturbation of the RT instability. In the real magnetosphere, the width of inward flow channel and the speed of inward flow are determined by magnetic flux conservation.

Figure 3 shows the topological structure of three-dimensional RT instability at the nonlinear stage (\( t = 20 \)). The isosurface of the plasma density at \( \rho = 1 \) (green surface) represents the boundary layer between the heavy and the light plasma, illustrating the localization of the RT mode along the \( z \) direction. The red and blue lines are the magnetic field lines traced from the top boundary on the heavy (i.e., \( x > 0 \)) and the light (i.e., \( x < 0 \)) plasma sides, respectively. With the growth of localized convection flow driven by the RT instability, the magnetic flux is strongly twisted near \( z = 30 \) (i.e., the transition between the RT active region and the unperturbed region), generating antiparallel magnetic field components mainly in the \( xy \) plane. As such, reconnection will be triggered once the width of the thin current layer formed by these antiparallel magnetic field components is on the ion inertial scale. The purple patch inside of the twisted magnetic flux is the isosurface of the parallel electric field at 0.01, indicating the reconnection site. This high-latitude (i.e., \( z = 30 \)) reconnection site is expected to operate at the edge of Io torus (i.e., about \( 2R_I \)) in the real Io environment.

A better illustration of the localization of the parallel electric field is presented in Figure 4, showing the integration of the parallel electric field along the \( y \) direction, \( \phi_y = \int E_\parallel dy \). The bright yellow region indicates the
significant increase of the parallel electric field, and the cone shape (the light purple region) represents the RT active region. It can be readily inferred from the direction of the vortex and the magnetic field that the majority of the parallel electric field is expected to be positive. However, the filamentary structure near the equatorial (black region) plane implies the existence of small-scale current layers. Such small-scale current layers with a typical width being much smaller than the overall structure (i.e., about 0.5 or eighth of the RT wavelength) can be clearly visualized in Figure 5, which shows the parallel current density $j_{\parallel}$ (color index) and in-plane current density $j_{xy}$ (black arrows) in the $z = 5$ plane at $t = 20$. Note the negative $j_{\parallel}$ region (purple region in Figure 5) is associated with the negative $\phi_y$ region (black region in Figure 4), where the parallel electric field, $E_{\parallel} = \eta j_{\parallel}$, and $\eta$ is the resistivity.

Figure 2 (bottom) shows the field-aligned potential difference integrated from the equatorial plane to the top boundary (color index) in the equatorial plane at $t = 18$, corresponding to the strong vortices formed by the convection flow pattern, which is consistent with our argument of reconnection driven by the twisting of magnetic flux. The purple dots are the footprints of the magnetic field line traced from the line $x = 0$ in the top boundary at $t = 18$, while the blue triangles are the fluid particles traced initially from the line $x = 0$ in the equatorial plane. The separation of these two sets of markers indicates the violation of the frozen-in condition, since they are originally on the same magnetic field lines. The correlation between the large separation and the strong field-aligned potential difference demonstrates that the frozen-in condition is broken at the high-latitude reconnection sites.
Figure 3. The topological structure of three-dimensional RT instability at the nonlinear stage ($t=20$). The green and purple surfaces are the isosurface of the plasma density at $\rho = 1$ and the parallel electric field at 0.01, respectively. The blue and red lines are the selected magnetic field lines. The direction of these selected magnetic field lines near the high-latitude reconnection site is indicated by red and blue arrows, respectively. The plasma density in the equatorial plane is represented by the color index (i.e., the red region is heavier than the blue).

Based on the symmetry properties, there is another reconnection site with negative parallel electric field near the $z = -30$ region, which does not show in Figure 3. A magnetic field line traced from the heavy side (i.e., $x > 0$) in the top boundary undergoing reconnection at both sites returns to the heavy side in the $z = -L_z$ plane, which is the same as its initial condition. However, the portion of the flux tube with high plasma content between these two reconnection sites has been replaced by the flux with lower plasma content. This process transports plasma radially outward, without involving the convection of magnetic flux in the ionosphere.

Figure 6 shows the change of flux tube content, $\Delta M = M|_{t=20} - M|_{t=0}$ (top), the change of flux tube entropy, $\Delta H = H|_{t=20} - H|_{t=0}$ (middle), and the field-aligned potential difference, $\phi$ (bottom) in the top boundary at $t = 20$, illustrating the radial transport process. In our slab configuration, the initial flux tube content has a negative gradient along the $x$ direction (i.e., $\partial_x M|_{t=0} < 0$), whereas the entropy has a positive gradient (i.e., $\partial_x H|_{t=0} > 0$). The change of these two quantities show a similar pattern but with different sign, being consistent with the scenario that double reconnection exchanges a portion of the flux tube. Note, an increase of the flux tube content, $\Delta M$, along the $x$ direction in the simulation implies a radial outward transport of the plasma, consistent with expectation. In contrast, a decrease of the flux tube entropy, $\Delta H$, along the $x$ direction in the simulation indeed agrees with the process of flux tube exchange in a slab configuration driven by
the RT mode. In reality, the centrifugally driven interchange instability operates when the flux tube content or the flux tube entropy decreases radially. For Jupiter's inner magnetosphere, the generation of hot plasma (e.g., 100 eV) from pickup in the vicinity of Io's orbit locally increases the flux tube content and entropy, being clearly centrifugal interchange unstable. As such, the interchange instability can transport both flux tube content and entropy radially outward. For Saturn's inner magnetosphere, the cold dense plasma from Enceladus can significantly increase the flux tube content. However, the colder plasma might have a minor influence.

Figure 4. Integration of parallel electric field along $y$ direction at $t = 20$.

Figure 5. Parallel current density, $j_\parallel$ (color index) and in-plane current density, $j_{xy}$ (black arrows) in $z = 5$ at $t = 20$. 
Figure 6. (top) The change of the flux tube content $\Delta M$, (middle) the change of the flux tube entropy $\Delta H$, and (bottom) the parallel electric potential difference along the magnetic field line $\phi$ integrated from the top boundary to the equatorial plane in the $z = L_z$ plane at $t = 20$. on the flux tube entropy. Thus, the centrifugally driven interchange instability operates when the destabilizing effects from a negative gradient of the flux tube content overcome the stabilizing effects from a positive gradient of the flux tube entropy. It is expected that heavy flux tube content can still be transported radially outward by the interchange instability. However, understanding how flux tube entropy is radially transported in such configurations requires a further study by using a more realistic configuration.

5. Summary and Discussion

In this study, it is demonstrated that the RT instability in a three-dimensional slab configuration forms a convection flow pattern, locally twisting the magnetic flux, consequently forming a pair of the high-latitude reconnection sites. This process exchanges a part of the flux tube, thereby transporting the plasma radially outward without involving the convection of magnetic flux in the ionosphere. We argue that in spite of the differences discussed in section 2, instabilities including, but not limited to, the RT instability and the centrifugal instability can lead to the formation of high-latitude reconnection sites in the giant magnetospheres. The onset condition for high-latitude reconnection requires sufficiently large antiparallel magnetic field components driven by the low-latitude localized convectional flow. The growth of any instabilities with an azimuthal wave vector leads to the radial inward or outward motion of footprints of magnetic flux tube in the equatorial region. Such a process will eventually generate a convective flow mainly due to the incompressibility of the plasma with a low plasma beta (i.e., $\beta \ll 1$ for inner magnetosphere), which is insensitive to the detailed process of the instability. However, the unstable region is likely to be localized along the polar ($z$) direction for the following two reasons. First, the large density gradient and centrifugal force are mainly localized at low
latitudes. Second, magnetic flux convection in the ionosphere is limited by the finite Pedersen conductivity. In such conditions, the localized instabilities are likely to generate antiparallel magnetic components. To further trigger magnetic reconnection thin current sheets corresponding to these antiparallel magnetic components are required. Note the sharp boundary (i.e., < 70 km) [see Kivelson et al., 1997, Figure 4] of “interchange” events is an important observational signature, suggesting that the width of the current layer is about the ion inertial scale (i.e., ≈ 20 km) as required with the reconnection onset condition.

Our model suggests a radial transport of plasma by exchanging a part of the neighboring flux tube without involving a motion of the entire flux, being somewhat different from the definition of the “interchange” by Kivelson [2005]. This means that the outward plasma transport process does not have to drag the entire magnetic flux tube from the inner magnetosphere to the outer magnetosphere. In this study, the model does not radially transport magnetic flux due to the symmetry properties, which is a singularity case. More generally, braking the symmetry properties by simply including a magnetic $B_z$ component in the equilibrium state can cause a magnetic field line to undergo reconnection at the high-latitude sites asynchronously. In this case, one would expect an intermittent radial transport of magnetic flux. However, the total net magnetic flux depletion should remain small. As such, even few narrow inward flow channels with a moderate speed may sufficiently refill the magnetic flux.

Due to the lack of in situ measurements of plasma velocity, the inward flow speeds are uncertain. Thorne et al. [1997] estimated an inward speed of ≈ 100 km s$^{-1}$, assuming that the flow travels from 6.3$R_J$ to 6.03$R_J$ and that the particles cannot drift out of the inward channel. Here the source region of the inward flow is identified by the phase space density for energetic S$^+$ ions. However, the estimated inward flow density (i.e., ≈ 100 cm$^{-3}$) by Thorne et al. [1997] is close to the density near 9$R_J$ from Jupiter rather than 6.3$R_J$ [Bagenal and Delamere, 2011]. If that is the case, the inward speed should be close to 1000 km s$^{-1}$, being faster than the Alfvén speed. In contrast, Frank and Paterson [2000] suggested that the radial component of plasma flow was directed toward Jupiter at a speed of 4 km s$^{-1}$. Recently, Yoshikawa et al. [2014] inferred the inward flow velocity of hot plasma to be on the order of 100 km s$^{-1}$ from the Earth-orbiting EUV spectroscope EXCEED (Extreme Ultraviolet Spectroscope for Exospheric Dynamics) on board the Hisaki spacecraft [Yoshikawa et al., 2014].

For a typical Io torus plasma environment, the electron density is about ≈ 2000 cm$^{-3}$ at a temperature of ≈ 5 eV. The in situ observation of the superthermal electrons (i.e., ≈ 100 eV with a density ~1% of the thermal population) leads to a critical unresolved issue of the physical source mechanism for these ubiquitous superthermal electrons. Based on the electron temperature profile observed by Voyager 1 [Sittler and Strobel, 1987], superthermal electrons are suggested to be transported radially inward from about 9$R_J$. However, the thermal equilibration time scale for those superthermal electrons near Io is about 1 h, being comparable to the radial inward traveling time from 9$R_J$ to the Io torus with an assumption of 100 km s$^{-1}$ inward speed. As such, significantly diffused superthermal electrons are unlikely to be observed by Galileo near the Io torus. However, high-latitude double reconnection driven by the interchange instability provides a plausible source of continuous energetic electrons through different mechanisms, which include, but are not limited to, parallel electric fields and wave-particle interactions (e.g., ion-inertia Alfvén wave). Therefore, superthermal electrons may be presented in the Io torus even if the inward flow from 9$R_J$ is 4 km s$^{-1}$.

The parallel electrical field also implies a positive dynamo (i.e., $\mathbf{j} \cdot \mathbf{E} > 0$), meaning the dissipation of Poynting flux, which may have auroral implications. Reconnection at high latitudes also violates the assumption of equipotential between magnetosphere and ionosphere, reducing the coupling of the magnetosphere to the ionosphere and being a possible contributing factor to the subcorotational flow.

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