Stability of the Magnetotail Current Sheet With Normal Magnetic Field and Field-Aligned Plasma Flows

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Abstract One of the most important problems of magnetotail dynamics is the substorm onset and the related instability of the magnetotail current sheet. Since the simplest 2D current sheet configuration with monotonic \( B_z \) was proven to be stable to the tearing mode, the focus of the instability investigation moved to more specific configurations, for example, kinetic current sheets with strong transient ion currents and current sheets with non-monotonic \( B_z \) (local \( B_z \) minima or/and peaks). The stability of the latter current sheet configuration has been studied both within kinetic and fluid approaches, whereas the investigation of the transient ion effects was limited to kinetic models only. This paper aims to provide a detailed analysis of the stability of a multi-fluid current sheet configuration that mimics current sheets with transient ions. Using the system with two field-aligned ion flows that mimic the effect of pressure non-gyrotropy, we construct a 1D current sheet with a finite \( B_z \). This model describes well recent findings of very thin intense magnetotail current sheets. The stability analysis of this two-ion model confirms the stabilizing effect of finite \( B_z \) and shows that the most stable current sheet is the one with exactly counter-streaming ion flows and zero net flow. Such field-aligned flows may substitute the contribution of the pressure tensor nongyrotropy to the stress balance but cannot overtake the stabilizing effect of \( B_z \). Obtained results are discussed in the context of magnetotail dynamical models and spacecraft observations.

1. Introduction

The problem of current sheet stability is key for most theories and models of magnetospheric dynamics because such stability determines magnetic reconnection onset (Gonzalez & Parker, 2016; Yamada et al., 2010) and triggers magnetospheric substorms (Angelopoulos et al., 2013; Baker et al., 1996; Sitnov, Birn, et al., 2019). A short review of investigations of the magnetotail current sheet instability should start from the work by Coppi et al. (1966) and Laval et al. (1966), suggesting that magnetic reconnection results from the tearing mode driven by electron currents. Schindler et al. (1973) and Schindler (1974) showed that in realistic magnetotail configuration the normal magnetic field component \( B_z \) (see sketch in Figure 1) destroys the resonant electron interaction with tearing mode (at least for realistic \( B_z \) magnitudes), so only ions can drive tearing. The inclusion of \( B_z \neq 0 \) into self-consistent current sheet models require either consideration of a specific class of exact 2D solutions (see, e.g., Kan, 1973 and the most recent generalizations in Vasko et al., 2013; Yoon & Lui, 2005) or weakly-2D (with \( \partial^2 / \partial x^2 \ll \partial^2 / \partial z^2 \)) solutions (see, e.g. Schindler, 1972, and the most recent generalizations in Artemyev et al., 2016; Birn, Schindler, & Hesse, 2004; Schindler & Birn, 2002). Further investigation of electron stabilizing effect due to \( B_z \neq 0 \) (Galeev & Sudan, 1985; Galeev & Zelenyi, 1976) reduces the parametric space for instability, whereas the consideration of such weakly-2D current sheet configuration (where \( j_z B_z / c \) tension force is balanced by the plasma pressure gradient \( \partial p / \partial x \), see sketch in Figure 1) demonstrated that such current sheets are stable (Lembege & Pellat, 1982; Pellat et al., 1991; Quest et al., 1996), with a possible exception of unrealistically stretched field line configurations with extremely small \( B_z \) (Goldstein & Schindler, 1982). To resolve the contradiction between well observed magnetic reconnection in the Earth's magnetotail (e.g., Angelopoulos et al., 2008; Nagai & Machida, 1998; Petrukovich et al., 1998; Sergeev et al., 1995) and theoretical current sheet stability, other models have considered the possibility of the reduction of electron stabilizing effect (Kuznetsova & Zelenyi, 1991; Zelenyi et al., 2008) and comprehensive and more general current sheet configurations (Pritchett & Buchner, 1995; Sitnov & Schindler, 2010; Sitnov et al., 2013). The important role in the development of such new current sheet models has been played by a series of MHD (Birn, Dorelli, et al., 2004; Birn et al., 1998) and kinetic...
(particle-in-cell) (Pritchett & Coroniti, 1994, 1995; Pritchett et al., 1991) simulations of the thin current sheet formation in the magnetotail. These simulations show formation of non-monotonic $B_z(x)$ profile with a local $B_z$ minimum and inverse $\partial B_z/\partial x < 0$ gradient. Theoretically the formation of a $B_z$ minimum in the near-Earth plasma sheet may attribute to the steady earthward convection (Hau et al., 1989). Further investigation of a current sheet with $\partial B_z/\partial x < 0$ has shown that such current sheet is tearing unstable (Bessho & Bhattacharjee, 2014; Birn et al., 2018; Merkin et al., 2015; Pritchett, 2015; Sitnov et al., 2014). There is also indirect observational evidence for the formation of such current sheet configurations in the near-Earth magnetotail (Angelopoulos et al., 2020; Sergeev et al., 2018). All such theoretical models of current sheet instability utilize the same class of 2D current sheets based on the stress balance ($j_z B_z/c = \partial p/\partial x$) for (for review of current sheet models of this class, see Baumjohann et al., 2007; Schindler, 2006 and references therein). However, as will be discussed below, modern spacecraft observations suggest that magnetotail current sheets may not always belong to this class.

A finite $H$, field in the middle and distant magnetotail current sheet (where magnetic field line curvature does not contribute to the cross-tail pressure balance) requires the plasma pressure gradient $\partial p/\partial x$ to balance the tension force $j_z B_z/c$, whereas the equatorial plasma pressure $p$ must equal the lobe magnetic field pressure $B^2/8\pi$ (Baumjohann et al., 1990; Petrukovich et al., 1999) that is fitted by a series of empirical models (Nakai et al., 1991; Shukhtina et al., 2004). These models give $B_{lobe} \approx 25nT \cdot (-x/20R_E)^{-1}$, and the corresponding $j_z \approx (eB_{lobe}/4\pi B_z) \cdot (\partial B_{lobe}/\partial x)$ is limited to $\sim 2 \text{nA/m}^2$ for radial distances $x \sim 20R_E$ and realistically small $B_z \sim \text{InT}$. As such current density amplitudes are much smaller than what is typically observed $5 - 10 \text{nA/m}^2$ (Lu et al., 2019; Petrukovich et al., 2015; Runov et al., 2006; Vasko et al., 2015), the stress balance $j_z B_z/c = \partial p/\partial x$ cannot universally describe the magnetotail equilibrium. If we write down the current sheet thickness $L_z \approx cB_{lobe}/4\pi j_z$ and length $L_x \approx (\partial \ln p/\partial x)^{-1}$, this stress balance gives $L_x \approx L_z B_{lobe}/2B_z$, whereas there are no observations of current sheets with $L_x$ much exceeding this limit (Artemyev et al., 2011, 2015). Empirical reconstructions of the magnetotail configuration during the growth phase of a substorm (i.e., before magnetic reconnection) also demonstrates the existence of very long current sheets, likely with $L_x \gg L_z B_{lobe}/2B_z$ (Sitnov et al., 2021; Sitnov, Stephens, et al., 2019). Thus, statistical observations at the middle (and distant) magnetotail of intense currents $j_z \geq 5 - 10 \text{nA/m}^2$, indirect $L_x$ estimates and empirical magnetotail reconstructions suggest that 2D current sheet models with $j_z B_z/c = \partial p/\partial x$ may not be suitable for investigation of the magnetotail current sheet stability (at least for a significant sample of observed configurations). An exploration of models alternative to the class with $j_z B_z/c = \partial p/\partial x$ is therefore necessary. This argumentation is not valid for the near-Earth magnetotail where the curvature force can contribute to the cross-sheet (vertical) pressure balance. Investigation of stability of the near-Earth current sheet requires analysis of a fully 2D magnetotail configurations (e.g., Goldstein & Schindler, 1982; Schindler, 2006).

Plasma anisotropy (if any) may significantly reduce the tension force $(1 - \Lambda)j_z B_z/c = \partial p/\partial x$ with $\Lambda = 4\pi (p_i - p_e)/B^2$ (Cowley, 1978; Rich et al., 1972), that is, current sheets at the limit of fire-hose instability
(Λ → 1) may be almost 1D (Cowley & Pellat, 1979; Francfort & Pellat, 1976). In particular, two-dimensional MHD equilibrium models have shown that both firehose pressure anisotropy and field-aligned flows may yield more stretched magnetotail configurations (Hau, 1993, 1996). However, ions, mostly contributing to the thermal pressure, are mainly isotropic in the magnetotail current sheet (Artemyev et al., 2019; C.-P. Wang et al., 2013), whereas parallel anisotropic electrons (Artemyev et al., 2014; Walsh et al., 2011) do not contribute more than Λ ~ 0.1 – 0.3 for the absolute majority of current sheets (Artemyev, Angelopoulos, Vasko, et al., 2020). More subtle (less measurable and more complex to be described) with respect to anisotropy, another kinetic effect, pressure non-gyrotropy (p_{cz} ≠ 0), allows balancing of 1D current sheet magnetic forces j_z B_z/c = ∂p_{cz}/∂z (Ashour-Abdalla et al., 1996; Burkhart et al., 1992; Eastwood, 1972; Mingalev et al., 2007, 2018; Pritchett & Coroniti, 1992). This effect has been included into the series of kinetic current sheet models (e.g., Sitnov & Merkin, 2016; Sitnov et al., 2003, 2006; Zelenyi et al., 2011 and references therein) that can describe many properties of observed current sheets (Artemyev & Zelenyi, 2013). Although direct spacecraft measurements of ion non-gyrotropy in the magnetotail current sheet are quite challenging (see discussion in Artemyev et al., 2010, 2019; Zhou et al., 2009), such a non-gyrotropy seems to be a prospective solution of the dilemma why magnetotail current sheets are much longer than 2D stress balance limit L_z = L_x B_0/c/2 B_z. Therefore, investigations of stability of 1D current sheets (i.e., current sheet that mimics a p_{cz} ≠ 0 effect) should reveal if they are more unstable to tearing mode than the very stable 2D current sheets. This p_{cz} ≠ 0 effect is a solution for the 1D current sheet balancing alternative to the balance by the cross-sheet plasma flow. Thus, we consider both these mechanisms of 1D current sheet balance: p_{cz} ≠ 0 in absence of cross-sheet bulk flow and strong cross-sheet flow (as in the rotational discontinuity balance).

Although full PIC simulations of 2D current sheet stability (e.g., Liu et al., 2014; Lu et al., 2018; Pritchett et al., 1991, 1997; Sitnov & Swisdak, 2011; Sitnov et al., 2009) have provided many details on the external driver threshold needed to trigger magnetic reconnection, there are almost no investigations of the stability of a 1D non-gyrotropic current sheet. To include such investigations into the more general context of the tearing instability (Birn & Priest, 2007; Biskamp, 2000), it would be reasonable to start with a fluid resistive model that can reveal a B_z role in 1D current sheet instability. To mimic the effect of p_{cz} ≠ 0 in the stress balance of the fluid current sheet model with B_z ≠ 0, we will adopt multi-fluid model proposed by Steinthauer et al. (2008). In this model two counter-streaming ion flows create v_x ∂_x v_x terms that balance j_z B_z/c force with a zero net flow. The generalization of this model would contain imbalanced flows, that is to say, a non-zero net flow. This generalization can be then reduced to a single fluid model with the field-aligned flows balancing j_z B_z/c force, that is, to the classical rotational discontinuity (Hudson, 1970) with the flow velocity equal to the Alfvén velocity. This paper is devoted to the investigation of the stability of such a generalized 1D multi-fluid model.

The formation and instability of current sheets with stretched magnetic field lines is a common problem in both magnetotail and solar physics (Reeves et al., 2008; Terasawa et al., 2000). In the latter case magnetic reconnection is believed to explain charged particle acceleration and magnetic field energy release in eruptive flares (the so called standard CSHKP model, see Carmichael, 1964; Hirayama, 1974; Kopp & Pneuman, 1976; Sturrock, 1966), non-eruptive events (including, e.g., coalescence of magnetic loops, see Sakai & de Jager, 1996), streamers (e.g., Edmondson & Lynch, 2017; Riley & Luhmann, 2012; Réville et al., 2020), and solar wind current sheets (Gosling, 2012; Phan et al., 2006). In contrast to the magnetotail investigations, mostly focused on the stabilization of the tearing mode by B_z field, the theory of tearing instability for solar applications is dominated by models of resistive tearing mode in B_z = 0 sheets (e.g., Del Sarto et al., 2016; Dobrowolny et al., 1983; Loureiro et al., 2012; Ofman et al., 1991; Tenerani et al., 2015). However, the stabilizing effect of B_z has also been discussed in context of the solar physics (Somov & Verneta, 1993; Verneta & Somov, 1987). Therefore, our investigation of the tearing instability in 1D current sheet with B_z ≠ 0 and plasma flows may be of great interest to solar physics applications.

The stress balance in the 1D current sheet model with B_z ≠ 0 is controlled by counter-streaming plasma flows (see Figure 1b), whereas imbalance of these flows results in the net plasma flow across the current sheet. For solar wind plasma, such cross-sheet flow is due to current sheet (rotational discontinuity) motion relative to the solar wind (Hudson, 1970; Tsurutani & Ho, 1999). In the Earth's magnetotail, there are a couple of mechanisms responsible for the formation of such counter-streaming plasma flows. First, the formation of a thin current sheet is often associated with enhanced precipitations of hot plasma sheet electrons.
into the ionosphere, and these precipitations drive the ionospheric outflow consisting of cold oxygen and hydrogen ions (Keika et al., 2013; Kronberg et al., 2015; Maggiolo & Kistler, 2014). Outflow ions shape fast beams moving along magnetic field lines (Artemyev, Angelopoulos, Runov, & Zhang, 2020; Kistler et al., 2005; Sauvaud et al., 2004) and contribute to the stress balance $\rho_z$ (Eastwood, 1972, 1974; Hill, 1975). Although such beams forming in the south and north hemispheres should be generally balanced (i.e., there is a stress balance without a net flow), the precise balance between them is not guaranteed, and there could be a net flow across the current sheet. Second, there are beams of energetic ions moving along the plasma sheet boundary layer from the distant magnetotail and coming back to the plasma sheet after reflection from the Earth’s dipole field (Ashour-Abdalla et al., 1992, 1996; Grigorenko et al., 2011). These are solar wind protons accelerated in the distant magnetotail by convection (Ashour-Abdalla et al., 1993; Cowley & Shull, 1983; Zelenyi, Dolgonosov, et al., 2006) or reconnection electric fields (Ashour-Abdalla et al., 2006; Grigorenko et al., 2009). Such acceleration mechanisms in combination with Speiser motion (Lyons & Speiser, 1982; Speiser, 1965) shape ion field-aligned beams that contribute significantly to the stress balance in 1D current sheet (Burkhart et al., 1992; Mingalev et al., 2007; Pritchett & Coroniti, 1992). The asymptotic solutions (for infinitely small $B_t$) of such 1D current sheets with counter-streaming plasma flows form a class of models developed in Sitnov et al. (2000) and then generalized by Sitnov et al. (2003); Sitnov et al. (2006) and (Zelenyi et al., 2011; Zelenyi, Malova, et al., 2006). Models of this class describe bifurcated and embedded current density profiles with properties similar to current sheet properties in the Earth’s magnetotail (see model/observation comparison in (Artemyev et al., 2008; Sitnov et al., 2006; Zhou et al., 2009). The distinguishing feature of these models is the sufficiently long 1D current sheet with the current density magnitude significantly exceeding estimates for 2D isotropic current sheet equilibria, $c(\partial p/\partial x)/B_z$. In this study, we do not specify a particular mechanism of field-aligned plasma flows and focus on the stability of the 1D current sheet model with such flows.

The paper is organized as follows. In Section 2 we present the equations for the zeroth-order equilibrium and the derivation of equations for the perturbed fields. In Section 3, we show the results, that is, the dispersion relation of the resistive tearing mode, obtained by numerically solving the perturbation equations derived in Section 2. In Section 4 we discuss the results in the context of Earth’s magnetosphere. In Section 5 we conclude this study.

2. Basic Equations

In this section, we present in detail the model of the background field and derivation of the linearized equation set which is solved numerically. The numerical results are presented in Section 3.

2.1. Background Fields

Incompressibility is assumed throughout the paper such that the plasma density is homogeneous and unperturbed: $\rho(x,t) \equiv \rho_0$. The background magnetic field consists of a Harris-type anti-parallel component and a uniform normal component:

$$ \mathbf{B} = B_0 \tanh \left( \frac{z}{a} \right) \hat{e}_z + B_z \hat{e}_z $$  \hspace{1cm} (1)

The current density flows along $y$ across magnetic field lines lying in $(x,z)$ plane, whereas a field-aligned plasma flow does not change this magnetic field geometry. One can show that $\nabla \times (\mathbf{B} \cdot \nabla \mathbf{V}) \neq 0$ with this configuration and thus it is impossible to maintain equilibrium with a scalar pressure $P = P_1$ solely. As already discussed in Section 1, one simple way to establish equilibrium is introducing the plasma flow such that the shear stress of the flow balances the tension force of the magnetic field. In one-fluid MHD model, an Alfvénic flow $\mathbf{V} \equiv \mathbf{B}/\sqrt{4\pi P}$ is required such that $\rho \mathbf{V} \cdot \nabla \mathbf{V} \equiv \mathbf{B} \cdot \nabla \mathbf{B}/4\pi$ and the equilibrium is achieved with a uniform total pressure $P^T = P + B^2/8\pi$.

We can consider a more generalized case where the protons shape two populations, denoted by subscripts “+” and “−” respectively such that the two corresponding momentum equations are:

$$ \rho_{\pm} \mathbf{V}_{\pm} \cdot \nabla \mathbf{V}_{\pm} = -\nabla P_{\pm} + \frac{q}{m_p} \rho_{\pm} \left( \mathbf{E} + \frac{1}{c} \mathbf{V}_{\pm} \times \mathbf{B} \right) $$  \hspace{1cm} (2)
where \( E \) is the electric field, \( q \) is the charge of proton, and \( m_p \) is the proton mass. Sum up the two equations and use the relation \( E \approx -V_e \times B/c \) where \( V_e \) is the electron flow velocity from the massless and cold electron momentum equation

\[
\rho e V_e \cdot \nabla V_e + \rho e V_+ \cdot \nabla V_+ = -\nabla \left( P + \frac{B^2}{8\pi} \right) + \frac{1}{4\pi} B \cdot \nabla B
\]

we get

\[
\rho e V_e \cdot \nabla V_e + \rho e V_+ \cdot \nabla V_+ = -\nabla \left( P + \frac{B^2}{8\pi} \right) + \frac{1}{4\pi} B \cdot \nabla B
\]

where \( P = P_+ + P_- \) and we have used the approximation \( J \approx c \nabla \times B/4\pi \) where \( J \) is the electric current density. For convenience we assume \( \rho_+ = \rho_- = \rho/2 \), i.e., the two ion populations are of the same density. Then we write \( V_z = \pm \alpha \nu_A \) where \( \nu_A = B/\sqrt{4\pi \rho} \) is the magnetic field in Alfvén velocity unit and \( \alpha \) is two constants controlling the speeds of the two ion flows. Apparently we have assumed that both of the two ion populations are streaming along the magnetic field lines. With these assumptions, we get two equilibrium criteria:

\[
\begin{aligned}
\frac{d}{dz} \left( P + \frac{B^2}{8\pi} \right) &= 0 \\
\left[ 1 - \frac{1}{2} \left( \alpha_+^2 + \alpha_-^2 \right) \right] B \frac{dB}{dz} &= 0
\end{aligned}
\]

Here we note that all the background fields are functions of coordinate \( z \) only. The first equation is the pressure balance condition which constrains the thermal pressure \( P(z) \) and the second equation is the tension force balance condition which leads to the requirement

\[
\alpha_+^2 + \alpha_-^2 = 2
\]

Without loss of generality, we assume \( 1 \leq \alpha_+ \leq \sqrt{2} \) and thus \( \alpha_- = \pm \sqrt{2 - \alpha_+^2} \). The positive (+) branch of \( \alpha_+ \) corresponds to the case that the two ion flows are counter-streaming and the negative (−) branch of \( \alpha_- \) corresponds to the case that the two ion flows are of the same direction. We define \( \mu = (\alpha_+ - \alpha_-)/2 \) such that the average flow speed \( V = (V_+ + V_-)/2 = \mu \nu_A \), that is, \( \mu = V/\nu_A \) is the amplitude of the average ion velocity normalized by the Alfvén velocity. In Figure 2, we plot the \( \alpha_+ (\alpha_+) \) and \( \mu (\alpha_+) \) curves for \( \alpha_+ \in [1, \sqrt{2}] \). By varying \( \alpha_+ \) and selecting either the positive or negative branch of \( \alpha_- \), we are able to regulate the average flow speed such that \( 0 \leq \mu \leq 1 \). Specifically, \( \alpha_+ = \alpha_- = 1 \) corresponds to \( \mu = 0 \) as in this case the two ion populations are counter-streaming with the same speed \( \nu_A \). On the other hand, \( \alpha_+ = \pm 1 \) leads to \( \mu = 1 \) as in this case the two ion populations are streaming with the same velocity \( V = \nu_A \) and this model converges to the one-fluid MHD model.

2.2. Perturbation Equations

To linearize the equation set, we write the perturbed, or first-order, fields as \( u_{z,e}, \rho_{z,e}, b, j \) and \( E_z \) respectively and assume all perturbations are in the form:

\[
f(x, z, t) = f(z) \exp(ikt + \gamma t).
\]

The linearized momentum equations are:

\[
\begin{aligned}
\frac{1}{2} \frac{d}{dz} \left( P + \frac{B^2}{8\pi} \right) &= 0 \\
\left[ 1 - \frac{1}{2} \left( \alpha_+^2 + \alpha_-^2 \right) \right] B \frac{dB}{dz} &= 0
\end{aligned}
\]

where \( \frac{d}{dz} \left( P + \frac{B^2}{8\pi} \right) \approx -\nabla \nu_A \). For convenience, we further define

\[
u = \frac{1}{2} \left( u_+ + u_- \right), \quad w = \frac{1}{2} \left( u_+ - u_- \right)
\]

where \( u \) is the perturbation of the ion net flow. By adding up and taking difference between the two equations of Equation 7, we get

\[
\begin{aligned}
\gamma u + f_+ \cdot \nabla \nu_A + V_A \cdot \nabla f_+ &= -\frac{1}{\rho_0} \nabla (p_+ + p_-) + \frac{1}{c \rho_0} (j \times B + J \times b) \\
\gamma w + f_- \cdot \nabla \nu_A + V_A \cdot \nabla f_- &= -\frac{1}{\rho_0} \nabla (p_+ - p_-) + \frac{q}{m_p c} \left[ w \times B + \frac{1}{2} (\alpha_+ + \alpha_-) V_A \times b \right]
\end{aligned}
\]
where \( f_\pm = (\alpha_+ u_+ ± \alpha_- u_-) / 2 \), and we have used the assumption \( A \mathbf{V} \pm = ± \alpha \mathbf{V} \pm \mathbf{A} \).

If we normalize length to the thickness of the current sheet \( A \mathbf{A} \), normalize magnetic field to \( A \mathbf{A} \), normalize speed to the Alfvén speed \( V_{A0} = B_0 / \sqrt{4\pi \rho_0} \), Equation 9b becomes:

\[
\gamma \mathbf{w} + \mathbf{f}_e \cdot \nabla \mathbf{B} + \mathbf{B} \cdot \nabla \mathbf{f}_e = -\nabla (p_+ - p_-) + \frac{a}{d_i} \left[ \mathbf{w} \times \mathbf{B} + \frac{1}{2} (\alpha_+ + \alpha_-) \mathbf{V}_A \times \mathbf{b} \right]
\]

(10)

where \( d_i = c / \sqrt{4\pi ne^2 / m_e} \) is the ion skin depth and \( a \) is the plasma number density. We can see that, in the MHD limit where the typical length scale \( A \mathbf{A} \) is much larger than the ion skin depth \( d_i \), this equation reduces to:

\[
\mathbf{w} \times \mathbf{B} + \frac{1}{2} (\alpha_+ + \alpha_-) \mathbf{V}_A \times \mathbf{b} = 0
\]

(11)

or alternatively

\[
\mathbf{w} = \frac{1}{2} (\alpha_+ + \alpha_-) \mathbf{b}
\]

(12)

which correlates the magnetic field perturbation and the difference between the perturbations of the two ion flow velocities. In the special case \( a_\pm = ±1 \), we have \( \mathbf{w} \equiv 0 \), which is consistent with the fact that the two ion populations merge into one fluid. The induction equation is acquired by inserting \( \mathbf{E}_1 = -\mathbf{u}_e \times \mathbf{B} / c - \mathbf{V}_e \times \mathbf{b} / c + \eta \mathbf{j} \), where \( \eta \) is the resistivity, into the linearized Maxwell-Faraday equation \( \partial \mathbf{b} / \partial t = -c \nabla \times \mathbf{E}_1 \). As we are analyzing the MHD case and hence the Hall effect is neglected, we can write \( \mathbf{E}_1 = -\mathbf{u} \times \mathbf{B} / c - \mathbf{V} \times \mathbf{b} / c + \eta \mathbf{j} \). After normalization, the linearized induction equation is written as:

\[
\gamma \mathbf{b} = \mathbf{b} \cdot \mathbf{V} \cdot \mathbf{V} + \mathbf{b} \cdot \mathbf{V} \mathbf{u} - \mathbf{u} \cdot \nabla \mathbf{b} + \frac{1}{2} \nabla \mathbf{B}^2
\]

(13)

where \( S = aV_{A0} / \eta' \) is the Lundquist number and \( \eta' = c^2 \eta / 4\pi \) is the magnetic diffusivity. Here we have used the divergence-free conditions for both the magnetic field and the flow velocity.
After some algebra starting from Equations 9a, 12 and 13, a two-equation set for \( u_z \) and \( b_z \) is written as

\[
\gamma \left( u''_z - k^2 u'_z \right) + \mu \left\{ B_z \left( u'''_z - k^2 u'_z \right) + ik \left[ B_z \left( u''_z - k^2 u_z \right) - B''_z u_z \right] \right\} = 0
\]
\[
\sigma \left\{ B_z \left( b'''_z - k^2 b'_z \right) + ik \left[ B_z \left( b''_z - k^2 b_z \right) - B''_z b_z \right] \right\}
\]
\[
\gamma b_z = \left( ik B_z u_z + B_z u'_z \right) - \mu \left( ik B_z b_z + B_z b'_z \right) + \frac{1}{S} \left( b''_z - k^2 b_z \right)
\]

where \( \mu = (\alpha_+ - \alpha_-)/2, \sigma = 1 - (\alpha_+ + \alpha_-)^2/4 \) and prime indicates \( d/dz \). From Equation 14, we see that \( B_z \) is a singular parameter as it increases the order of the equation for \( u_z \) from two to three. In addition, compared with the tearing mode with an anti-parallel magnetic field, the growth rate is in general complex in this model rather than purely real, implying propagating perturbations. It is also worth noting that, if \( B_z = 0 \) and \( \alpha_+ = \alpha_- = 0 \) such that \( \mu = 0, \sigma = 1 \), Equation 14 reduces to the classical tearing mode equation:

\[
\gamma \left( u''_z - k^2 u'_z \right) = \left( ik B_z u_z + B_z u'_z \right) - \mu \left( ik B_z b_z + B_z b'_z \right) + \frac{1}{S} \left( b''_z - k^2 b_z \right)
\]

### 3. Stability Analysis

The system of Equation 14 is a boundary-value eigen-problem. The boundary conditions are that \( u_z \) and \( b_z \) vanish far from the current sheet: \( u_z, b_z (z \to \pm \infty) = 0 \). In practice, we cannot set the boundaries to infinity. However, far from the current sheet we have \( B_z \approx \pm B_0 \), i.e., \( B_z \) is approximately constant, so it is observed that \( u_z, b_z \propto \exp(-k|z|) \) are solutions to Equation 14a. Plugging the condition \( u_z, b_z \propto \exp(-k|z|) \) into Equation 14b, we get the ratio between \( u_z \) and \( b_z \) at the boundaries. We use the boundary-value-problem solver implemented in the Python library SciPy (Virtanen et al., 2020) to solve Equation 14. The solver adopts a fourth order collocation algorithm with the control of residuals (Ascher et al., 1994; Kierzenka & Shampine, 2001) and is able to solve the eigenvalue and eigen-functions simultaneously. The solver has been successfully utilized to analyze the linear stability of the oblique tearing mode with a guide field (Shi et al., 2020).

#### 3.1. Effect of \( B_z \neq 0 \)

In Figure 3, we plot the dispersion relations \( \gamma - k \) and \( \omega - k \) in the top (panel (a)) and bottom (panel (b)) panels respectively. We note that, here \( \gamma \) and \( \omega \) are the real and opposite-imaginary parts of the complex \( \gamma \) appearing in Equation 14, that is, \( \gamma_{\text{complex}} = \gamma - i\omega \). For this figure, the Lundquist number \( S \) is fixed at \( 10^4 \) and \( \mu, \sigma \), i.e., the average ion speed normalized to \( V_{\text{Shear}} \), is fixed at 1.0, corresponding to an Alfvénic one-fluid model as discussed in Section 2.1. In each panel, curves of different colors correspond to different values of \( B_z \), ranging from 0.001 to 0.1. To validate the solver and to testify whether the solution to Equation 14 converges as the singular parameter \( B_z \to 0 \), we also solve the problem with \( B_z = 0 \) at which the equation set degenerates to a lower order. The solved \( \gamma \) and \( \omega \) for \( B_z = 0 \) are plotted as dashed curves and we can see that as \( B_z \) approaches to 0 from a finite value, the \( \gamma - k \) curve converges to the dashed one. We note that if \( B_z \) is exactly zero, \( \omega \) is also zero, that is, the mode is not propagating due to the fact that the background shear flow is symmetric in \( z \). From panel (a) of Figure 3, we see that the growth rate is monotonically decreasing with \( B_z \). That is to say, the existence of a finite \( B_z \) quenches the growth of the instability.

In analysis of the tearing instability, how the maximum growth rate scales with the Lundquist number is important. In Figure 4, we show the maximum growth rate \( \gamma_m \), which is the peak value of each \( \gamma - k \) curve such as those shown in Figure 3, as a function of the Lundquist number \( S \) for \( \mu = 1 \) and different values of \( B_z \) in log-log scale. In the figure, squares are the numerical results and solid lines are linear-fittings of the squares. Obviously, \( \gamma_m \) decreases with increasing \( B_z \), consistent with the results shown in Figure 3. For \( B_z = 0 \), a clear power-law relation \( \gamma_m \propto S^{-0.65} \) is observed. We note that in a current sheet without any flow, the fastest growing tearing mode has a growth rate \( \gamma_m \propto S^{-0.5} \). Thus, the scaling \( \gamma_m \propto S^{-0.65} \) indicates that the current sheet is more stable with the Alfvénic background flow than without the flow. As we increase \( B_z \), the \( \gamma_m - S \) line is no longer straight, as can be seen from the curves \( B_z = 0.005 \) and \( B_z = 0.01 \), whose slopes gradually transit to \( -1 \) at large Lundquist numbers. When \( B_z \) is large enough, that is, \( B_z = 0.05 \) and \( B_z = 0.075 \), the whole \( \gamma_m - S \) curve in the regime \( S \gtrsim 10^3 \) is straight with a slope \( -1 \). We note that, a growth rate \( \gamma \propto S^{-1} \) implies that \( \gamma \propto \eta/u^2 \), that is to say, there is actually no “growth” of instability in the system
because the growth rate is supported purely by the diffusion of the background magnetic field. Thus, Figure 4 shows that, a finite $A_A \Delta z$ significantly stabilizes the current sheet: when either $A_A \Delta z$ or $A_A$ is large enough, the tearing instability vanishes and only the diffusion is taking effect in transferring the energy from the background magnetic field to the perturbed fields.

3.2. Effect of $\mu < 1$

In Figure 5, we plot the $\gamma - k$ (panel (a)) and $\omega - k$ (panel (b)) curves for $S = 10^4$, $B_z = 0.01$, and varying $\mu$, that is, the average ion flow speed. From panel (a), we see that the growth rate decreases with $\mu$ in general. As shown in Figure 6, we plot the maximum growth rate $\gamma_m$ as a function of the Lundquist number $S$ for $B_z = 0.01$ and different values of $\mu$ in log-log scale. Apparently, the $\gamma_m - S$ relation is not a single power law and the $\gamma_m - S$ curve steepens with $S$. Similar to what is shown in Figure 4, for very large $S$ the relation converges to $\gamma_m \propto S^{-1}$, that is, a status where the instability is merely a result of diffusion.

From Figures 5 and 6, we can see that as $\mu \to 0$, the growth rate converges to zero. This is the case $\alpha_+ = \alpha_- = 1$ where the two ion populations counter stream at exactly the same speed $V_{\text{io}}$. In this case, $\mu = \sigma = 0$.
since otherwise we are decreases with for we apply a two- and 6 in log-log scale. The squares are the 1986- that decays to zero at 1988 is achieved and varying (Chen et al., and we define the Lundquist number for 7 in this two-proton model is not equivalent to the flow speed in the as a function of and different flow speed as long as its velocity is parallel to the background magnetic field. In Figure 7, the results are calculated based on the one-fluid MHD model, that is, there is only one proton population. Figure 7 shows that, in the MHD case, γn decreases with μ, at least for μ ≤ 1 (Chen et al., 1997). For μ > 1, the shear flow is super-Alfvénic and thus Kelvin-Helmholtz instability will arise (S. Wang et al., 1988). Except for μ = 1.0, γn scales with S as γn ∝ S^{−0.5}. That is to say, although the background flow suppresses the tearing instability, the scaling relation of the maximum growth rate with S does not change. However, as μ approaches unity, the scaling relation changes to γn ∝ S^{−0.65}, implying that the system becomes more stable, consistent with previous studies (Dahlburg et al., 1997; Einaudi & Rubini, 1986) which show that the growth rate scales as S^{−β} with 1/2 < β < 2/3. But we note here that, these studies claim that in resistive-MHD regime the current sheet with Alfvénic plasma flow (i.e., our μ = 1 case) is stable. However, our numerical results indicate that this case is not ideally stable but has finite positive growth rate as shown in Figures 3 and 7. In general, we can write γnτL ∼ S^{−β} where S = aV_{A0}/η and τL = aV_{A0}. We note that this expression is for a 1D current sheet which is infinitely long. In practice, consider a current sheet of finite length L, we usually need to measure the time scale by τL = L/V_{A0} and we define the Lundquist number by SL = LV_{A0}/η. This transforms the scaling relation to γnτL ∼ SL^{−β}(a/L)^{−β−1}. Apparently, the inverse aspect ratio a/L is a key parameter determining the growth rate γnτL. For an arbitrary inverse aspect ratio a/L ∼ SL^{δ}, we get γnτL ∼ SL^{δ+β+1}, implying a threshold δ = β/(β + 1), at which γnτL ∼ O(1) is achieved and this is the so-called “ideal tearing” (Pucci & Velli, 2013; Pucci et al., 2020; Tenerani et al., 2016). For the classical tearing β = 1/2, we have δ = 1/3, that is to say when a macroscopic current thins to a/L ∼ SL^{1/3},

![Figure 4](image-url)
the growth rate of the tearing mode transits from extremely small (assuming $A_L \to \infty$) to unity and thus the current sheet breaks up rapidly. As $\beta$ increases, such as in the MHD case with Alfvénic flow shown in Figure 7 or in the two-ion case shown in Figures 4 and 6, the critical value of $\delta$ also increases, indicating that the current sheet must be thinner in order to achieve fast growth of the tearing instability. For example, for $\beta \approx 2/3$ as in the Alfvénic flow case, we have $\delta = 2/5 > 1/3$ and especially, for $\beta = 1$, we get the threshold $a/L \sim S_t^{-1/2}$, indicating that fast growth of instability happens only when Sweet-Parker type current sheet is formed. This is because, as we discussed before, $\gamma_n \tau_a \sim S_t^{-1}$ implies that the growth of instability is fully supported by diffusion.

### 3.3. Eigen-Functions

In Figure 8, we show the eigen-functions $u_c$ and $b_c$ for three sets of parameters. Panel (a) is $B_z = 0.01$, $\mu = 1.0$, panel (b) is $B_z = 0.05$, $\mu = 1.0$, and panel (c) is $B_z = 0.01$, $\mu = 0.051$. The shapes of the eigen-functions do not differ significantly among the three cases, though the relative amplitudes of $u_c$ and $b_c$ change with...
the parameters. The most outstanding feature of these eigen-functions is the quasi-odd function $A_3(z)$. It is known that for the classical tearing mode, $A_3$ is symmetric in $z$ such that $A_3(z = 0) > 0$, which provides the magnetic flux necessary for the reconnection to happen. Meanwhile, $A_3$ is naturally asymmetric in $z$ as $z = 0$ is the stagnation point of the flow. From Figure 8, we can see that $A_3$ is nearly asymmetric in $z$ so that $A_3(z = 0)$ is very small, limiting the reconnection rate. In addition, as the function $A_3(z)$ is very different from the classical tearing case, plasmoids will not be generated at the center of the current sheet but on the two sides of the current sheet instead.

4. Discussion

In this study we investigate the stability of 1D current sheet with a finite normal component of magnetic field $B_z \neq 0$ and the tension force $j_z B_z / c$ balanced by the plasma flows. Two main results are (a) strong stabilizing effect of $B_z$ in such 1D current sheets and (b) possible destabilization by an imbalance of counter-streaming flows, that is, by the field-aligned net flow. Next, we will discuss these results in the context of the Earth’s magnetotail dynamics.

4.1. Current Sheet Instability at Substorm Onset

As observations suggest the formation of thin (almost 1D) current sheets during the substorm growth phase (Artemyev et al., 2015; Petrukovich et al., 2013; Sitnov et al., 2021), there is a natural interest to investigate the instability of such current sheet configuration. The main idea is that 1D current sheets with nongyrotropic pressure tensor (i.e., with $B_j j_z / c \approx \partial p_c / \partial z$) would be more unstable than classical 2D current sheets (with $B_j j_z / c \approx \partial p / \partial x$), and this resolves the issue of magnetic reconnection onset during weak substorms without strong external drivers (see discussion in Zelenyi et al., 2008). We use multi-fluid current sheet model with the counter-streaming plasma flows generating such non-gyrotropy ($p_{cz} \neq 0$). Our results show that we cannot make current sheet with $B_z \neq 0$ more unstable solely by substituting $\partial p / \partial x$ by $\partial p_c / \partial z$.

![Figure 6. Maximum growth rate $\gamma_m$ as a function of $S$ for $B_z = 0.01$ and varying $\mu$. The squares are the numerical results and the solid lines are linear-fittings of the squares. The black dashed line is $\gamma \propto S^{-1}$ for reference.](image-url)
This result suggests that the investigation of spontaneous (undriven) instability in the magnetotail current sheet requires more specific current sheet configurations. One of the ideas has been proposed by Sitnov and Schindler (2010) for 2D current sheets and confirmed in a set of MHD (Birn et al., 2018; Merkin et al., 2015) and kinetic (Bessho & Bhattacharjee, 2014; Pritchett, 2015; Sitnov et al., 2014) simulations which show that current sheets with non-monotonical $A_A(x)$ profiles can be more unstable. Alternative ideas would be needed to explain spontaneous reconnection in 1D current sheets (with $A_A(x) = const$), and we showed that such ideas should include more than pressure nongyrotropy.

![Figure 7](image.png)

**Figure 7.** Maximum growth rate $\gamma_m$ as a function of $S$ for $B_z = 0$ and varying $\mu$ based on one-fluid MHD model. The squares are the numerical results and the solid lines are linear-fittings of the squares. The black dashed lines are $\gamma \propto S^{-0.5}$ and $\gamma \propto S^{-0.65}$ for reference.

This result suggests that the investigation of spontaneous (undriven) instability in the magnetotail current sheet requires more specific current sheet configurations. One of the ideas has been proposed by Sitnov and Schindler (2010) for 2D current sheets and confirmed in a set of MHD (Birn et al., 2018; Merkin et al., 2015) and kinetic (Bessho & Bhattacharjee, 2014; Pritchett, 2015; Sitnov et al., 2014) simulations which show that current sheets with non-monotonical $B_z(x)$ profiles can be more unstable. Alternative ideas would be needed to explain spontaneous reconnection in 1D current sheets (with $B_z = const$), and we showed that such ideas should include more than pressure nongyrotropy.

![Figure 8](image.png)

**Figure 8.** Numerically solved eigen-functions $u_z$ and $b_z$ for three sets of parameters: (a) $B_z = 0.01$, $\mu = 1.0$. (b) $B_z = 0.05$, $\mu = 1.0$. (c) $B_z = 0.01$, $\mu = 0.051$. Cyan solid curves are the real parts of $b_z$, blue dashed curves are the imaginary parts of $b_z$, coral dashed-dotted curves are the real parts of $u_z$, and red dotted curves are the imaginary parts of $u_z$. 
4.2. Field-Aligned Plasma Flows

We generalize the 1D current sheet model of (Steinhauer et al., 2008) by including a net plasma flow. Such flow is shown to be able to reduce the stabilization effect of $B_z \neq 0$. This is an interesting effect because similar field-aligned flows can be included into 2D current sheet equilibria in application to the magnetotail and solar flares (Birn, 1992; Hau, 1996; Nickeler & Wiegelmann, 2010). Using an approach proposed by Birn (1991); Birn (1992), one can introduce arbitrary flow along magnetic field lines in a single MHD, and such flow may influence the stability of 2D current sheets, i.e., altering the magnitude of inverse $\partial B_z / \partial x$ gradient that is considered to destabilize such current sheets (Birn et al., 2018; Merkin et al., 2015). Note, this destabilizing effect of field-aligned plasma flow differs from the stabilizing effect of cross-field diamagnetic flow considered by Swisdak et al. (2003). Indeed, we focus here on the stability conditions of structures with a finite normal component, whereas the (Swisdak et al., 2003) model and its observational confirmations (Phan et al., 2010, 2014) were obtained for much more unstable tangential discontinuities without a normal component $B_n$. Our results show that current sheets with $B_z \neq 0$ and plasma flow equal to the Alfvén speed are completely stable to tearing, and can therefore survive for a long time, for example, during the long lasting substorm growth phase. The same result explains the absence of magnetic reconnection signatures in multiple rotational discontinuities observed in the near-Sun region (Phan et al., 2020).

Obtained results generalize the previous investigations of the plasma flow effects on the current sheet instability in the absence of $B_z$, that is, for 1D tangential discontinuities with 1D (Chen et al., 1997; Hoshino & Higashimori, 2015; S. Wang et al., 1988) and 2D plasma flows (Ip & Sonnerup, 1996; Phan & Sonnerup, 1991). The effect of decrease of the growth rate for the long wavelength tearing mode due to the sub-Alfvénic plasma flow, found for such tangential discontinuities (Bulanov et al., 1978), has been revealed for the 1D current sheet with $B_z \neq 0$. However, in contrast to the generally stabilizing effect of the sub-Alfvénic plasma flow on the tearing more in tangential discontinuity (Dahlburg et al., 1997; Einaudi & Rubini, 1986), in the current sheet with $B_z \neq 0$ the field-aligned flow increases the peak growth rate, that is, the most stable current sheet with $B_z \neq 0$ is formed by two counter-streaming flows, whereas any imbalance between these two flows provides an additional free energy for the tearing mode. This effect resembles the effect of super-Alfvénic flow driving the tearing mode in the tangential discontinuities (S. Wang et al., 1988).

4.3. Speculation on Possible Drivers of Magnetic Reconnection in the Magnetotail

Our results show that the 1D current sheet, such as the ones formed at the late stage of the substorm growth phase, is well stable relative to the tearing mode. Together with results of 2D current sheet stability (see discussion in Sitnov et al., 2002), our results rise the question of initialization of substorm reconnection in absence of a strong external driver (i.e., spontaneous reconnection). One of the solutions of this issue has been proposed by Sitnov and Schindler (2010) and Sitnov et al. (2013) for 2D current sheets with non-monotonical $B_z(x)$ profiles. This solution has been shown to work for current sheets with $\partial p / \partial x = B_z j_z / c$, and one of the prospective directions of further investigation is the merging of models of such 2D current sheets and current sheet with field-aligned plasma flows. Inclusion of such flows (in multi-fluid approach) will relax $\partial p / \partial x = B_z j_z / c$ condition by addition $\partial p_{\perp i} / \partial z$ term, and thus realistic quasi-1D current sheets with a weak $\partial p / \partial x$ can be described. From other point of view, even weak $\partial p / \partial x$ allows inclusion of a non-monotonical $B_z(x)$ profiles as a possible source of free energy for the tearing mode. Another prospective approach consists of further generalization of multi-fluid models with the inclusion of realistic anisotropic electron components (see discussion in Artemyev, Angelopoulos, Vasko, et al., 2020). Intense electron currents in high-beta plasma of thin current sheets may provide an additional free energy source for the tearing mode excitation.

5. Conclusions

In this work, we carried out linear stability analysis of the 1D MHD current sheet with a uniform normal magnetic field ($B_z$ in the magnetotail coordinates). In our model, the magnetic tension force is balanced by the field-aligned ion flow, which is generalized by assuming that ions shape two populations of the same density but with different field-aligned velocities. The main results of our calculation are
The existence of a finite \( B \), significantly stabilizes the tearing instability. With a finite \( B_0 \), the scaling relation between the maximum growth rate of the instability (\( \gamma_m \)) and the Lundquist number (\( S \)) converges to \( \gamma_m \propto S^{-1} \) in the limit \( S \to \infty \), indicating the instability is purely diffusion-supported.

2. In this two-ion model, the most stable case is when the two ion populations counter-stream at the same speed \( V = B/\sqrt{4\pi \rho} \) such that the net flow velocity is zero. In this case, the instability does not grow at all. An imbalance between the two ion populations, which leads to a non-zero net flow, destabilizes the current sheet. The most unstable case is when the two ion populations stream in the same direction at the same speed \( V = B/\sqrt{4\pi \rho} \). This is the one-fluid MHD case where the average flow velocity is \( B/\sqrt{4\pi \rho} \).

Our results show that, in the magnetotail, a 1D, that is, very long and thin, the current sheet which has a finite \( B \) component is very stable to the tearing instability. To answer the question of the onset of fast spontaneous reconnection in the magnetotail, more ingredients are required, such as the ion/electron kinetic effects.

**Data Availability Statement**

The boundary-value-problem solver used in this study can be found at [https://docs.scipy.org/doc/scipy](https://docs.scipy.org/doc/scipy) and is described in Virtanen et al. (2020).

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