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The scaling of collisionless magnetic reconnection in an electron–positron plasma with non-scalar pressure

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Abstract. Collisionless magnetic reconnection via tearing instability in non-relativistic electron–positron (pair) plasma with an anisotropic pressure is investigated. The equilibrium magnetic field is considered to be sheared force-free, and a set of linearized collisionless Magnetohydrodynamics equations describes the evolution of reconnection dynamics. A linear analytical analysis, based on scaling, demonstrates that in such a pair plasma, breaking the frozen in flow constraint for field lines can be mainly provided by the non-gyrotropic pressure of electrons and positrons (rather than the particle bulk inertia) when the current sheet width is smaller than the particle Larmor radius ($\Delta x < r_L$). This condition is satisfied when $\beta > d^2$ ($d = c/\omega_p$ is the particle skin-depth with the electron/positron frequency ω_p and $\beta = 8\pi P^{(0)}/B_0^2 \ll 1$). Meanwhile, on top of the Lorentz force and in the absence of the reconnection facilitating mechanism of the Hall effect, non-scalar pressure force can accelerate bulk plasma into the diffusion region at the scale lengths of the order of d. Therefore, the respective regime of tearing instability proceeds much faster compared with the case of an isotropic pressure with a new dimensionless growth rate of $(\gamma\tau_A) \sim d$.

1. Introduction

Magnetic reconnection via tearing instability plays an instrumental role in numerous laboratory and astrophysical plasma physics phenomena. The stored magnetic energy converts explosively into the kinetic and thermal energy of plasma in short timescales (Priest and Forbes 2000; Zweibel and Yamada 2009; Yamada et al. 2010; Lewis et al. 2012). The main problem is the observed rate of reconnection, which typically exceeds by far theoretical predictions based on the single-fluid magnetohydrodynamic (MHD) models with classical resistivity (Furth et al. 1963). Two-fluid description of a plasma brings the Hall effect and the electron bulk inertia. The former effect facilitates reconnection process, and the latter one is an additional (on top of plasma resistivity) mechanism of the magnetic field line breaking, usually called collisionless reconnection.

There is, however, one special case in which the Hall effect is absent from the outset, which is an electron–positron plasma: equal masses of the negatively and positively charged particles eliminate the Hall effect (see Sec. 2). Although a pair plasma does not support the Hall effect, fast reconnection in pair plasmas has been reported in several simulations and theoretical works (Bessho and Bhattacharjee 2007, 2010; Daughton and Karimabadi 2007). In fact a pair plasma is an ideal environment to study the physics of collisionless reconnection. On the other hand, the electron–positron plasma is also of interest on its own in a number of astrophysical objects such as jets from active galactic nuclei (Lesch and Birk 1998; Wardle et al. 1998; Larrabee et al. 2003), plasma winds (Lyubarsky and Kirk 2001), gamma ray bursts (Drenkhahn 2002), as well as in laboratory experiments (Greaves and Surko 1995; Pedersen et al. 2012).

In a collisionless pair plasma, when the velocity distribution function of particles is not restored to an isotropic one by sufficiently infrequent collisions, the non-scalar pressure effects become important in reconnection process. Some publications have discussed magnetic reconnection in an electron–positron plasma (Shukla et al. 1986, 1996; Stenflo et al. 1986; Yu et al. 1986; Pritchett 1995; Zenitani et al. 2001; Bessho and Bhattacharjee 2005, 2007, 2010; Hosseinpour and Vekstein 2008; Karlicky 2008; Cai et al. 2009; Zenitani et al. 2009), and, in particular, some of them considered the anisotropic pressure effects (Bessho and Bhattacharjee 2005, 2007, 2010; Cai et al. 2009). Particle-in-cell simulations of Bessho and Bhattacharjee (2005, 2007, 2010) showed that the fast reconnection can happen in pair plasmas. Meanwhile, divergence of the electron–positron pressure tensor essentially plays the role of an effective collisionless resistivity in fast reconnection process. Also, Cai et al. (2009) investigated the tearing mode with pressure gradient effects in pair plasmas, and concluded that in collisionless regime, where the electron inertia is dominant, effects of pressure gradient are different for small and large pressure gradients. Small pressure gradients can enhance the growth rate of tearing mode, while the large pressure gradients would reduce the growth rate. On the other hand, Daughton and Karimabadi (2007) discussed that in the large-scale electron–positron plasmas, the reconnection rate is determined by the system size, the reconnection is highly dynamic, and a steady-state reconnection is never established. Furthermore, Shukla et al. (1996) by assuming an isotropic pressure and a sheared magnetic field found that the particle inertia can cause tearing instability in a collisionless electron–positron plasma even with negative values of the stability parameter, $\Delta' < 0$.

In this study, by taking the electron and positron anisotropic pressure effects into account, we analytically analyze the linear evolution of reconnection dynamics (i.e., studying evolution of the in-plane perturbed magnetic fields and the plasma flow) due to tearing instability in a sheared force-free magnetic field. We are interested in a reconnection regime, for which the off-diagonal components of the non-scalar pressure break the frozen in flow condition rather than the bulk inertia. On the other hand, the force exerted by the non-scalar pressure, besides the Lorentz force, can accelerate plasma flow toward the reconnection site, and therefore the corresponding growth rate of this regime of tearing mode so that the respective reconnection rate could modify. Here the main aim is to investigate the instability growth rate of tearing mode in the presence of non-scalar electron/positron pressure tensor in governing MHD equations (e.g., the generalized Ohm's law and the equation of plasma motion). To do so, we have made some simplifying assumptions in our analysis (see Sec. 2), the major of which is the constant-ψ approximation. In Sec. 2, we sketched the model and basic equations. Then Sec. 3 gives analytical analysis of MHD equations and discussion of this analysis, and a brief summary is followed in Sec. 4.

2. The model and basic equations

We consider a tearing unstable slab of a highly conductive electron–positron plasma, which is embedded in a sheared force-free magnetic field,

$$\mathbf{B}^{(0)}(x) = \hat{z}B_z^{(0)}(x) + \hat{y}B_y^{(0)}(x), \tag{2.1}$$

with

$$B_y^{(0)} = B_0 f(x), \;\; B_z^{(0)} = B_0[1 - f^2(x)]^{1/2}, \tag{2.2}$$

and surrounded by two perfectly conducting walls at $x = \pm l$. It is assumed that $f(x)$ is an odd function of x, so for the tearing perturbation of the form $exp(iky + \gamma t)$ reconnection occurs at the resonant surface, $x = 0$, where the poloidal field component $B_y^{(0)}$ changes its sign. The dynamics of an electron–positron plasma follows from the equation of motion for these two species,

$$nm\left(\frac{\partial \mathbf{V}_e}{\partial t} + \mathbf{V}_e \cdot \nabla \mathbf{V}_e\right) = -ne\left(\mathbf{E} + \frac{1}{c}\mathbf{V}_e \times \mathbf{B}\right) - \nabla \cdot \mathbf{P}_e, \tag{2.3}$$

$$nm\left(\frac{\partial \mathbf{V}_p}{\partial t} + \mathbf{V}_p \cdot \nabla \mathbf{V}_p\right) = ne\left(\mathbf{E} + \frac{1}{c}\mathbf{V}_p \times \mathbf{B}\right) - \nabla \cdot \mathbf{P}_p. \tag{2.4}$$

By summing up (2.3) and (2.4), one gets the equation of motion of a pair plasma,

$$\rho\left[\frac{d\mathbf{V}}{dt} + \frac{1}{4n^2e^2}(\mathbf{j} \cdot \nabla)\mathbf{j}\right] = \frac{1}{c}\mathbf{j} \times \mathbf{B} - \nabla \cdot (\mathbf{P}_e + \mathbf{P}_p), \tag{2.5}$$

while their subtraction yields the generalized Ohm's law,

$$\mathbf{E} + \frac{1}{c}\mathbf{V} \times \mathbf{B} = \frac{m}{2ne^2}\left[\frac{\partial \mathbf{j}}{\partial t} + (\mathbf{V} \cdot \nabla)\mathbf{V} + (\mathbf{j} \cdot \nabla)\mathbf{V}\right] + \frac{1}{2ne}\nabla \cdot (\mathbf{P}_p - \mathbf{P}_e). \tag{2.6}$$

Here $\mathbf{V} = (\mathbf{V}_e + \mathbf{V}_p)/2$ is the plasma bulk velocity and $\mathbf{j} = ne(\mathbf{V}_p - \mathbf{V}_e)$ is the current density, which represents the rotation of magnetic field via the Ampère's law, $\mathbf{j} = (c/4\pi)\nabla \times \mathbf{B}$. Together with the induction equation

$$\frac{\partial \mathbf{B}}{\partial t} = -c\nabla \times \mathbf{E}, \tag{2.7}$$

(2.5)–(2.7) provide a coupled set of equations describing collisionless reconnection in a non-relativistic pair plasma. As seen from (2.6), a pair plasma does not support the Hall effect (i.e., the absence of $\mathbf{j} \times \mathbf{B}$ term): The magnetic field lines remain frozen into bulk plasma flow unless the electron/positron bulk inertia or non-scalar pressure tensor are taken into account. The Hall effect cancellation means that there is no perturbation in the z component of magnetic field (quadruple guide field). Furthermore, plasma motion in the $(x - y)$ plane can be assumed incompressible. The electron/positron thermal pressure is defined as

$$\mathbf{P}^{e,p} = m\int (\mathbf{u} - \mathbf{V})_{e,p}(\mathbf{u} - \mathbf{V})_{e,p} f(\mathbf{u}, \mathbf{r}, t)_{e,p} d^3\mathbf{u}_{e,p}, \tag{2.8}$$

where $\mathbf{V}_{e,p} = \int \mathbf{u}_{e,p} f(\mathbf{u}, \mathbf{r}, t)_{e,p} d^3\mathbf{u}_{e,p} / \int f(\mathbf{u}, \mathbf{r}, t)_{e,p} d^3\mathbf{u}_{e,p}$, and $f(\mathbf{u}, \mathbf{r}, t)_{e,p}$ is the electron/positron anisotropic distribution function. Temporal evolution equation for the pressure tensor of particles can be derived by taking the second moment of the Vlasov equation in the center of mass system of electrons/positrons, which yields

$$\frac{\partial P_{\alpha\beta}^{e,p}}{\partial t} = -\frac{\partial}{\partial x_l}\left(Q_{\alpha\beta l}^{e,p} + V_{l,ep}P_{\alpha\beta}^{e,p}\right) - P_{\alpha l}^{e,p}\frac{\partial V_{\beta,ep}}{\partial x_l} - P_{\beta l}^{e,p}\frac{\partial V_{\alpha,ep}}{\partial x_l} - \frac{e}{mc}B_l\left(\varepsilon_{\alpha jl}P_{j\beta}^{e,p} + \varepsilon_{\beta jl}P_{j\alpha}^{e,p}\right), \tag{2.9}$$

where the electron heat flux tensor is

$$\mathbf{Q}^{e,p} = m\int (\mathbf{u} - \mathbf{V})_{e,p}(\mathbf{u} - \mathbf{V})_{e,p}(\mathbf{u} - \mathbf{V})_{e,p} f(\mathbf{u}, \mathbf{r}, t)_{e,p} d^3\mathbf{u}_{e,p}. \tag{2.10}$$

Our present analysis concentrates on a closure at the level of the second-order moment and, therefore, we ignore the electron/positron heat flux tensor contribution for the sake of simplicity. One can assume that in the vicinity of the resonant surface, $x = 0$, the initial magnetic field $\mathbf{B}^{(0)}$ has a uniform shear: $f(x) \approx \alpha x$, i.e., $B_y^{(0)} \approx B^{(0)}\alpha x$, $B_z^{(0)} \approx B^{(0)}$, where α^{-1} is the magnetic shear length. An initially almost uniform plasma with a gyrotropic pressure tensor ($P_{\alpha\beta}^{(0),e,p} = \delta_{\alpha\beta}P^{(0)}$) is assumed. In what follows, it is useful to introduce two non-dimensional parameters: $\epsilon \equiv \alpha l$, which is a measure of the magnetic shear inside the reconnection current sheet, and $d \equiv d/l$, the scaled inertial skin-depth of particles, which is always small ($d \ll 1$) for all applications of interest. It is convenient to express the magnetic field as

$$\mathbf{B}(x, y, t) = \nabla\psi(x, y, t) \times \hat{z} + B_z\hat{z}, \tag{2.11}$$

where the poloidal flux function is

$$\psi(x, y, t) = \psi_0(x) + \psi_1(x)\, exp\,(iky + \gamma t) \tag{2.12}$$

$[\psi_0 = -B_0 \int f(x)dx \approx -B_0\alpha x^2/2]$. In addition, the bulk velocity of incompressible plasma $\mathbf{V}$ is represented as

$$\mathbf{V} = (\nabla\phi \times \hat{z}) + V_z\hat{z}, \tag{2.13}$$

where the stream function $\phi(x, y, t) = \phi(x)\, exp\,(iky + \gamma t)$ corresponds to the vortical component of the poloidal plasma flow. Note that initially the equilibrium plasma flow is zero.

The linearized version of (2.5) and (2.7) reads as

$$\gamma\psi_1 = -cE_z^{(1)} = ikB_0\alpha x\phi + \frac{\gamma mc^2}{4\pi n^2}\nabla^2\psi_1 + \frac{c}{2ne}\left(\frac{\partial P_{xz}^{(1),p}}{\partial x} + \frac{\partial P_{xy}^{(1),p}}{\partial y} - \frac{\partial P_{xz}^{(1),e}}{\partial x} - \frac{\partial P_{xy}^{(1),e}}{\partial y}\right), \tag{2.14}$$

$$\gamma\phi_1 = \frac{icB_0\alpha x}{8\pi nmk}\nabla^2\psi_1 - \frac{1}{2nm}\left(\frac{\partial P_{xx}^{(1),p}}{\partial x} + \frac{\partial P_{yx}^{(1),p}}{\partial y} + \frac{\partial P_{xx}^{(1),e}}{\partial x} + \frac{\partial P_{yx}^{(1),e}}{\partial y}\right). \tag{2.15}$$

Equations (2.14) and (2.15) are the x components of (2.7) and (2.5) respectively. The first term on the right-hand side of (2.14) is due to the advective bulk plasma flow, the second term is due to the particle bulk inertia, and the last term is related to the non-scalar pressure of electrons and positrons. A contribution made by the advective part of the inertial term $(\mathbf{V} \cdot \nabla)\mathbf{V}$ is of the order of ϵ^2 and, therefore, can be neglected here. Similar to this, the corresponding contributions due to $(\mathbf{j} \cdot \nabla)\mathbf{j}$ and $(\mathbf{j} \cdot \nabla)\mathbf{V}$ terms in (2.14) and (2.15) are also of the order of ϵ^2. Furthermore, the two terms on the right-hand side of (2.15) are related, respectively, to the Lorentz force and the non-gyrotropic pressure gradients. In order to analyze (2.14) and (2.15) one needs to find out the pressure tensor perturbations, $P_{\alpha\beta}^{(1),ep}$. This can be done by linearizing (2.9) about the initial equilibrium. We assume that the tensor perturbation components have the form $P_{\alpha\beta}^{(1),ep} = P_{\alpha\beta}^{(1),ep}(x)\, exp\,(iky + \gamma t)$. Inside the current sheet, the initial out-of-plane field $B_z^{(0)}$ is dominant, so one can put there $\mathbf{B}^{(0)} = B_0\hat{z}$. The linearized (2.9) for $P_{\alpha\beta}^{(1),ep}$ takes the form

$$\gamma P_{\alpha\beta}^{(1),ep} = -P^{(0)}\left(\frac{\partial V_{\beta,ep}^{(1)}}{\partial x_\alpha} + \frac{\partial V_{\alpha,ep}^{(1)}}{\partial x_\beta}\right) - \omega_B\left(\varepsilon_{\alpha jz}P_{j\beta}^{(1),ep} + \varepsilon_{\beta jz}P_{j\alpha}^{(1),ep}\right), \tag{2.16}$$

with $\omega_B = eB_0/mc$. As seen from (2.16), the pressure tensor perturbation is caused by the non-uniformity of the bulk flow of particles, which is counter-balanced by the isotropization effect of perpendicular (out-of-plane) magnetic field. A straightforward solution of (2.16) yields the following expressions for $P_{\alpha\beta}^{(1),ep}$:

$$P_{xx}^{(1),ep} = -P_{yy}^{(1),ep} = \frac{P^{(0)}}{2\omega_B}\left(\frac{\partial V_{y,ep}^{(1)}}{\partial x} + \frac{\partial V_{x,ep}^{(1)}}{\partial y}\right) - \frac{\gamma P^{(0)}}{2\omega_B^2}\frac{\partial V_{x,ep}^{(1)}}{\partial x}, \tag{2.17}$$

$$P_{xy}^{(1),ep} = -\frac{P^{(0)}}{\omega_B}\frac{\partial V_{x,ep}^{(1)}}{\partial x} - \frac{\gamma P^{(0)}}{4\omega_B^2}\left(\frac{\partial V_{y,ep}^{(1)}}{\partial x} + \frac{\partial V_{x,ep}^{(1)}}{\partial y}\right), \tag{2.18}$$

$$P_{xz}^{(1),ep} = \frac{P^{(0)}}{\omega_B}\frac{\partial V_{z,ep}^{(1)}}{\partial y} - \frac{\gamma P^{(0)}}{\omega_B^2}\frac{\partial V_{z,ep}^{(1)}}{\partial x}, \tag{2.19}$$

$$P_{yz}^{(1),ep} = -\frac{P^{(0)}}{\omega_B}\frac{\partial V_{z,ep}^{(1)}}{\partial x} - \frac{\gamma P^{(0)}}{\omega_B^2}\frac{\partial V_{z,ep}^{(1)}}{\partial y}, \tag{2.20}$$

where the electron and positron bulk velocity perturbations, $\mathbf{V}_{e,p}^{(1)} = \mathbf{V} \mp \mathbf{j}/2ne$, are

$$V_{x,ep}^{(1)} = \frac{\partial\phi}{\partial y}, \quad V_{y,ep}^{(1)} = -\frac{\partial\phi}{\partial x}, \quad V_{z,ep}^{(1)} = -\frac{c}{4\pi ne}\left(\frac{\partial^2\psi_1}{\partial x^2} - k^2\psi_1\right). \tag{2.21}$$

After inserting (2.17)–(2.21) into (2.14) and (2.15), the closed system of equations is obtained for the magnetic field perturbation ψ_1 and the stream function ϕ. By introducing non-dimensional variables with length normalized by l, ψ_1 by $B_0 l$, ϕ by γl^2, these equations then become

$$\psi_1(0) = ik\epsilon x\phi + d^2\psi_1'' - \frac{1}{4}\beta\, d^4\psi_1^{IV}, \tag{2.22}$$

$$\phi = -\frac{i\epsilon x}{k(\gamma\tau_A)^2}\psi_1'' + \frac{1}{4}\beta\, d^2\phi''. \tag{2.23}$$

Here, $\tau_A = l/V_A = l(4\pi nm)^{1/2}/B_0$ is the Alfvèn transit time, $\beta = 8\pi P^{(0)}/B_0^2$. It is also assumed that $\nabla^2 \approx d^2/dx^2$ since $k(\Delta x) \ll 1$, where (Δx) is the current sheet width.

We are restricted to the case of a moderate magnitude of the stability parameter Δ' so that the 'constant-ψ' approximation can be used. To simplify analytical analysis in the next section, we put $k \sim \epsilon \sim 1$. It is worth noting the symmetry of the functions involved in (2.22) and (2.23). The poloidal flux function $\psi(x)$ has even parity, while the stream function $\phi(x)$ is an odd function of x. Therefore, terms with 'wrong' parity are omitted.

3. Analytical analysis and discussion

Three terms on the right-hand side of (2.22), as mentioned above, represent different effects involved in the magnetic field evolution: advection by the bulk flow of plasma (the first term), bulk inertia of electron (the second term), and contribution of the particle pressure tensor (the third term). The last two effects actually can break the frozen in flow condition. Similar to this, the two terms on the right-hand side of (2.23) are related to the Lorentz force and the non-gyrotropic electron/positron pressure force, which can accelerate the plasma flow. We restrict discussion to the case where the non-scalar particle pressure is the main mechanism for breaking the frozen-in flow constraint rather than the bulk inertia. It means that the following inequality must be satisfied in (2.22):

$$d^2\psi_1'' \ll \beta\, d^4\psi_1^{IV} \approx \beta\, d^4\frac{\psi_1''}{(\Delta x)^2},$$

or

$$(\Delta x) \ll \beta^{1/2}\, d = r_L, \tag{3.1}$$

where r_L is the particle gyroradius. Hence, if the current sheet width (Δx) is smaller than the particle gyroradius r_L, then the bulk inertial effect is not significant compared to the non-gyrotropic pressure effect, and so can be ignored in (2.22). The following analysis, basically, considers such small scale-length processes.

Now let us consider a case for which the non-scalar particle pressure effect is as significant as the Lorentz force in the reconnection dynamics. Thus, by putting equal the two terms on the right-hand side of (2.23)

$$\frac{\Delta x}{(\gamma\tau_A)^2}\psi_1'' \sim \beta\, d^2\phi'', \tag{3.2}$$

one gets

$$\phi \sim \frac{(\Delta x)^3}{\beta\, d^2(\gamma\tau_A)^2}\psi_1''. \tag{3.3}$$

By omitting the inertial term (the second term) in (2.22), the first and third terms can balance each other as

$$(\Delta x)\phi \sim \frac{(\Delta x)^4}{\beta\, d^2(\gamma\tau_A)^2}\psi_1'' \sim \beta\, d^4\psi_1^{IV}, \tag{3.4}$$

with ϕ being replaced from (3.3), and it yields

$$(\gamma\tau_A) \sim \frac{(\Delta x)^3}{\beta\, d^3}. \tag{3.5}$$

In order to obtain the respective instability growth rate, one needs to estimate the corresponding current sheet width (Δx) as follows: The last term on the right-hand side of (2.22), which can be written as $\beta\, d^4\psi_1^{IV} \sim \beta\, d^4\psi_1''/(\Delta x)^2 \sim \beta\, d^4\psi_1(0)\Delta'/(\Delta x)^3$ (note that $\Delta' = \psi_1^{(-1)}(0)\int \psi_1'' dx \approx \psi_1''\Delta x/\psi_1(0)$) provides the reconnected poloidal flux. Therefore, this term balances the left-hand side of this equation

$$\beta\, d^4\psi_1(0)\Delta'/(\Delta x)^3 \sim 1,$$

i.e.,

$$(\Delta x) \sim \beta^{1/3} d^{4/3}. \tag{3.6}$$

Now, after inserting (3.6) into (3.5), the following scaling for the growth rate of tearing instability obtains

$$(\gamma\tau_A) \sim d. \tag{3.7}$$

Note that in the absence of non-gyrotropic pressure of electrons and positrons, the bulk inertial effect becomes important. In this case, the growth rate of collisionless tearing instability is $(\gamma\tau_A) \sim d^3$ for $\epsilon \sim 1$ and $k \sim 1$ (Shukla et al. 1996). By reminding that $d \ll 1$, the growth rate of collisionless tearing mode in the presence of the non-scalar electron/positron pressure is much greater than the case with a scalar pressure, i.e., $(\gamma)_{\text{nonscalar}}/(\gamma)_{\text{scalar}} \approx d/d^3 = d^{-2} \gg 1$.

4. Summary

We investigated linear analytical analysis of the collisionless tearing instability in an incompressible and non-relativistic electron–positron plasma (pair plasma). A planar slab of a uniform plasma is considered with a sheared force-free equilibrium magnetic field. The non-scalar electron and positron pressure tensors are kept in governing MHD equations: the generalized Ohm's law and the equation of motion. Due to cancellation of the Hall effect, there would not be a perturbed guide magnetic field and, therefore, the reconnection dynamics is described by the evolution of the poloidal magnetic field and the plasma flow. Note that the initial out-of-plane magnetic field tends to isotropize the plasma. Under the parametric condition of $\beta > d^2$, which leads to a current sheet width smaller than the particle Larmor radius, one can ignore of the particle bulk inertial effect in violation of the frozen in constraint. Instead, this role can be played by the non-gyrotropic pressure of particles. On top of the Lorentz force, non-gyrotropies in the electron and positron pressure, at the scale length of the particle inertial skin-depth, can force plasma flow to accelerate. Subsequently, this flow advects field lines to the diffusion region (reconnection site), where the non-gyrotropic pressure effects support the reconnection mechanism. For this regime of tearing mode, our scaling resulted in a new dimensionless growth rate of $(\gamma\tau_A) \sim d$, which is, in fact, much faster than the case where the non-gyrotropic effects are insignificant, or basically the plasma pressure is scalar. In other words, the bulk inertia-driven tearing mode grows with the rate of $\sim d^3$ (Shukla et al. 1996), which is clearly much smaller than the growth rate of the non-scalar pressure driven mode, $\sim d$, since $d \ll 1$.

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