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Mechanisms of Spontaneous Reconnection: From Magnetospheric to Fusion Plasma

Lev Zelenyi · Anton Artemyev

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Abstract Very often space plasma is treated as collisionless. We check the validity of this paradigm considering various regimes of tearing mode (spontaneous reconnection) including effects of particle collisions and shear of magnetic field. We briefly describe Pitaevskii's effect of effective modification of collision frequency due to the finite particle Larmor radius in the presence of magnetic field. This effect results in a significant increase of the role of collisionality, especially in a weakly magnetized systems. Another popular paradigm is related with application of MHD description to collisionless or weakly collisional systems. We show, that for current sheets observed in the Earth magnetotail and magnetopause as well as for current sheets formed in Solar corona and in laboratory devices most appropriate is the kinetic semi-collisional tearing regime. Role of "collisions" could play usual Coulomb pair collisions of electrons and ions (e.g. in Solar corona) or effective collisions (scattering) of electrons with the microturbulence wave modes. Transition to real MHD modes requires either very large collisions frequencies and/or very large amplitudes of the magnetic field shear. The largest domain in the parameter space is occupied by the kinetic regimes of tearing mode growth where dissipation is provided either by Landau damping or by real (or effective) collisions.

Keywords Tearing instability · Magnetic reconnection

1 Introduction

Starting from the original paper by Giovanelli (1947) reconnection of magnetic field lines is considered as a main mechanism of magnetic energy dissipation. First MHD models describing quasi-stationary magnetic reconnection already included main elements of this process: diffusion region (Sweet 1958), slow shock waves (Petschek 1964) and thin

L. Zelenyi · A. Artemyev (✉)
Space Research Institute, RAS, Moscow, Russian Federation
e-mail: Ante0226@yandex.ru

L. Zelenyi
e-mail: lzelenyi@iki.rssi.ru

current sheet (Syrovatskii 1966). There are several comprehensive books (Parker 1994; Priest and Forbes 2000; Biskamp 2000; Birn and Priest 2007) devoted to various aspects of magnetic reconnection and related charged particle acceleration. The most of theories can be attributed to the one of two possible approaches: kinetic collisionless approach and fluid resistive approach.

Magnetic reconnection plays an important role in various plasma systems starting from rarefied collisionless plasma of interplanetary medium and planetary magnetospheres and going to a weakly collisional plasma of Solar corona and then to collision dominated plasma of laboratory devices (see review by Yamada et al. 2010). Initialization of the magnetic reconnection corresponds to the instability of current sheet separating magnetic fields with opposite polarities. Therefore, the problem of relationship between collision and collisionless reconnection regimes can be reformulated as a problem of current sheet instabilities in presence of collisions with arbitrary frequency ν_{eff} and of an magnetic shear (also of an arbitrary intensity). It should be noted that as was shown by Coppi et al. (1966b), heuristically even the case with $\nu_{eff} = 0$ could be reduced to collisional if one will take into account that Landau damping (providing necessary dissipation for the case with $\nu_{eff} \to 0$) could be roughly considered as supporting effective scattering of electrons with frequency $\nu_{eff} \sim v_{T_e}/\lambda$, where λ is the wavelength of the mode and v_{T_e} is electron thermal velocity.

First paper devoted to instability of current sheet relative to the tearing mode was written by Furth (1962). In this paper the stability of neutral current sheet with magnetic field reversal was considered relative to a periodical fluctuation of a normal component of magnetic field. The further development of theory of the tearing instability includes effects of electric field perturbation and effect of magnetic field shear (Laval et al. 1966). Investigation of the tearing mode based on energy variation principle was developed for current sheets (Schindler and Soop 1968) and generalized in the recent monograph by Schindler (2006).

Application of the tearing instability to collisionless plasma of the Earth magnetotail was done in pioneering work by Coppi et al. (1966b). Further investigations have shown that the principal role for this instability is played by the finite normal component of magnetic field, which magnetizes electrons and destroy corresponding Landau resonant damping (Schindler 1974; Galeev and Zelenyi 1976). Magnetized electrons provide the effect of tearing stabilization due to combination of the frozen-in condition and condition of quasi-neutrality. This effect could be so strong that the spontaneous reconnection mode will be stable for the entire parameter range (Pellat et al. 1991).

Stabilization of the magnetotail current sheet contradicts to numerous observations of magnetic reconnection (see, e.g., Angelopoulos et al. 2008). This problem can be solved by choice of proper initial equilibrium, which describes magnetotail current sheet with a number of additional realistic effects. For example it was shown that embedded thin current sheets often observed in the downtail are unstable relative to the tearing mode (Zelenyi et al. 2008, 2010). Alternative idea corresponds to current sheet with the reversed longitudinal gradient of the normal component of magnetic field. Such current sheets could become unstable relative to the tearing mode (Sitnov and Schindler 2010) or to a more exotic instabilities also resulting in magnetic reconnection (Pritchett and Coroniti 2011).

Although, stability problem for the magnetotail current sheet is of primary importance, in this review we consider mainly the stability of current sheet without normal component of magnetic field, but in presence of a shear magnetic field component having an arbitrary intensity. Therefore our consideration deals with current sheet of planetary magnetopause. Collisions (which is the primary goal of our paper) weakly influence the properties of equilibrium current sheet solutions, but once the system unstable—collisions strongly control the rate of instability growth. We will combine in our analysis effects of usual Coulomb collisions and effects of particle scattering at microturbulent fluctuations ("effective" collisions). These effects correspond to plasma microturbulence and can be important for current sheets in Solar corona and laboratory experiments.

2 Initial Equilibria

Harris current sheet (Harris 1962) can be considered as the simplest kinetic equilibrium model describing the basic properties of space current sheets (in absence of the normal component of magnetic field). This model corresponds to the velocity distribution f_α for particles with mass m_α and charge q_α introduced as shifted Maxwellian distribution:

$$f_\alpha = C_{0\alpha}\exp\bigl(-(H_\alpha - v_{D_\alpha}P_y)/T_\alpha\bigr) = C_\alpha N(z)\exp\bigl(-\bigl(v_x^2 + v_z^2 + (v_y - v_{D_\alpha})^2\bigr)/v_{T_\alpha}\bigr)$$

where $H_\alpha = m_\alpha(v_x^2 + v_y^2 + v_z^2)/2$ is particle energy, $P_y = m_\alpha v_y + (q_\alpha/c)A_y$ is generalized momentum, v_{D_α} is a constant particle drift velocity, $T_\alpha = m_\alpha v_{T_\alpha}^2/2$ is particle temperature, $C_\alpha = n_0(2\pi v_{T_\alpha})^{-3/2}$ is the constant of normalization with particle density in the central region of current sheet, n_0. Distribution of particle density is defined by function $N(z) = \cosh^{-2}(z/L)$, where L is the current sheet thickness. Here α denotes type of particles: $\alpha = i$ for ions and $\alpha = e$ for electrons.

Density of cross-tail current supported by distribution function f_α is $j_y(z) = (4\pi \times B_0/Lc)\cosh^{-2}(z/L)$ and resulting magnetic field acquires the simple form $B_x = B_0\tanh(z/L)$. Therefore, in the central region of current sheet $z \sim 0$ magnetic field B_x changes sign. This region is filled by particles crossing $z = 0$ and oscillating in nonlinear potential. Corresponding equation of particle motion across current sheet has the form $\ddot{z} \approx -z(const - z^2)$ (Sonnerup 1971). Unmagnetized particles are trapped inside the region $|z| < R_\alpha$, where $R_\alpha = \sqrt{L\rho_\alpha}$ with $\rho_\alpha = v_{T_\alpha}/\Omega_\alpha$ and $\Omega_\alpha = |q_\alpha|B_0/m_\alpha c$ (Dobrowolny 1968). Reflecting from magnetic "walls" $z = \pm R_\alpha$, these particles move along current sheet plane and can therefore interact with unstable waves accordingly to the Landau mechanism (see review by Galeev 1979, and references therein). Several recent investigation of tearing and drift instabilities were devoted to the precise calculations of the impact of these resonant particles (see, e.g., Lapenta and Brackbill 1997; Daughton 1999; Daughton and Karimabadi 2005; Karimabadi et al. 2005).

It can be noticed, that the velocity distribution of particles in the Harris current sheet remains the same even if the guide component of magnetic field $B_y = const$ is applied to the sheet. To take into account inhomogeneous $B_y = B_y(z)$ one needs to consider $j_x(z)$ current and corresponding modification of the velocity distribution (see review by Roth et al. 1996, and references therein). However, for simplified geometry with $B_y = const$ we can restrict our analysis by the modified Harris equilibrium distribution.

We consider below normal component $B_z \neq 0$ only for illustrating effect of modification of collision frequency (so called Pitaevskii effect). This effect could be very important for weakly magnetized and weakly collisional plasmas. One should note, however, that presence of $B_z \neq 0$ brings principal topological change to the system configuration and results in appearance of very different plasma equilibria (see models of 1D current sheet with $B_z \neq 0$ in Kropotkin et al. 1997; Sitnov et al. 2000; Zelenyi et al. 2000). Modification of the current sheet and corresponding accompanying effects are described in review by Zelenyi et al. (2011).

3 Dissipative Effects

Besides the Landau kinetic mechanism of dissipation providing the growth of certain wave modes in the current sheets, there are more standard effects of direct energy dissipation due to effective (turbulent conductivity) or Coulomb particle collisions. For magnetospheric plasma system certain role in generation of effective conductivity can be played by whistler waves (see, e.g., Deng and Matsumoto 2001), by Alfven-whistler mode (Huang and et al. 2012), and by lower-hybrid waves (see Huba et al. 1977; Fujimoto et al. 2011). Independently of dispersion of waves forming turbulence, their interaction with particles can be approximately considered as effective collisions with certain collision frequency ν_{eff}. For example, collision frequency provided by weak lower-hybrid turbulence can be described by expression

$$\nu_{eff} = \omega_{LH}\left(\frac{1}{2}\rho_i |d\ln N/dz|\right)^3 (T_i/T_e)^2$$

where ω_{LH} is lower-hybrid frequency.

To take into account this effect in kinetic model of current sheet instability we use collision integral in Bhatnagar-Gross-Krook (BGK) form (Bhatnagar et al. 1954), which was originally designed for Coulomb collisions. So in the consideration below ν_{eff} could have the meaning either of effective or Coulomb collision frequency. We consider only effect of collisions on instability and neglect by their influence on the initial equilibrium. Collision integral for perturbation of the velocity distribution $f_{1\alpha}$ can be written as

$$\mathrm{St} f_{1\alpha} = -\sum_\sigma \nu_{\alpha\sigma}\left(f_{1\alpha} - \frac{f_{0\alpha}}{\int f_{0\alpha} d\mathbf{v}}\left(\int f_{1\alpha} d\mathbf{v} + \frac{m_\alpha}{T_\alpha}\mathbf{v}\int \mathbf{v} f_{1\alpha} d\mathbf{v}\right)\right)$$

where $\nu_{\alpha\sigma}$ is frequency of collisions of α-type and σ-type particles: $\nu_{ie} = \mu\nu_{ei}$ with $\mu = m_e/m_i$. Linearized Vlasov equation gives following expression for $f_{1\alpha}$ (Zelenyi and Taktakishvili 1981):

$$\begin{aligned} f_{1\alpha} = {} & \frac{q_\alpha}{cT_\alpha} f_{0\alpha}\left(A_{1y} v_{D_\alpha} - c\varphi_1 + i\int_{-\infty}^0 \left(\left(\omega v_y(\tau) - v_{D_\alpha}\nu_\alpha\right)A_{1y} - \omega_\alpha c\varphi_1\right)\varepsilon_\alpha(\tau)d\tau\right) \\ & + \sum_\sigma \nu_{\alpha\sigma} f_{0\alpha}\int_{-\infty}^0 \left(n_{1\alpha} + \mathbf{v}(\tau)\int \mathbf{v} f_{1\alpha} d\mathbf{v}\right) n^{-1}\varepsilon_\alpha(\tau)d\tau \end{aligned} \tag{1}$$

where $n = n_0 N(z)$, $n_{1\alpha} = \int f_{1\alpha} d\mathbf{v}$, $\omega_\alpha = \omega + i\nu_\alpha$, $\nu_\alpha = \sum_\sigma \nu_{\alpha\sigma}$, and

$$\varepsilon_\alpha = \exp\left(-i\omega_\alpha\tau + ik\left(x(\tau) - x\right)\right)$$

Wavenumber of perturbation is k. Condition of quasineutrality $n_{1e} = n_{1i}$ gives the perturbation of the scalar potential φ_1 as a function of the perturbation of the vector potential A_{1y}. Substituting (1) into Maxwell equation we obtain

$$\begin{aligned} & \frac{d^2 A_{1y}}{dz^2} - \left(k^2 + V_0(z) + V^<(z)\right)A_{1y} = 0 \\ & V_0(z) + V^<(z) = -\frac{4\pi q_\alpha}{c}\sum_\alpha \int v_y f_{1\alpha} d\mathbf{v} \end{aligned} \tag{2}$$

where we separate adiabatic impact $V_0 = -2L^{-2}\cosh^{-2}(z/L)$ and resonant impact $V^<$. Term $V^<$ can be taken as zero for $|z| > R_\alpha$. The corresponding dispersion relation could be obtained by matching solutions of (2) in the internal ($|z| < R_\alpha$) and external ($|z| > R_\alpha$) regions. This technique is straightforward and described in details in may early publications (Dobrowolny 1968; Galeev 1979; Zelenyi and Taktakishvili 1981)

$$\frac{L}{2}\int_{-\infty}^{+\infty}\left(V_e^<(z) + V_i^<(z)\right)dx = \frac{1-(kL)^2}{kL} \tag{3}$$

Dispersion relation (3) can be rewritten as

$$\frac{\omega_{pe}^2 L^2}{c^2}\frac{\bar{\gamma}}{\bar{\nu}_e}\left(\sqrt{\frac{\rho_e}{L}}\frac{\mathrm{A}_e(1-\mathrm{A}_i)+\mathrm{A}_i(1-\mathrm{A}_e)}{1-\mathrm{A}_e\mathrm{A}_i} + \sqrt{\frac{\rho_i}{L}}\mathrm{A}_i\right) = \frac{1-(kL)^2}{kL} \tag{4}$$

where we neglect electron current density in the region $R_e < |z| < R_i$ in comparison with ion current and introduce the following notations: ω_{pe} is plasma frequency, $\bar{\gamma} = \gamma/kv_{T_e}$, $\bar{\nu}_e = \nu_e/kv_{T_e}$, and

$$\mathrm{A}_\alpha = -i\nu_{\alpha\sigma}\frac{Z_{0\alpha}}{kv_{T_\alpha}}\left(1 + i\nu_{\alpha\alpha}\frac{Z_{0\alpha}}{kv_{T_\alpha}}\right)^{-1}, \quad \alpha \neq \sigma \tag{5}$$

Plasma integral $Z_{n\alpha}$ (Kramp function) has a form

$$Z_{n\alpha}(\omega_\alpha/kv_{T_\alpha}) = \frac{1}{\sqrt{\pi}}\int_{-\infty}^{+\infty}\frac{x^n e^{-x^2}dx}{x-(\omega_\alpha/kv_{T_\alpha}) - i\vartheta\,\mathrm{sign}k}, \quad \vartheta \to 0$$

There are several cases when analytical solutions of (4) could be obtained.

For very weak collisions $\bar{\nu}_e \sim \nu_i/kv_{T_i} \ll 1$ we have $\omega_e \ll kv_{T_i} \ll kv_{T_e}$. Expansion of $Z_{0\alpha}$ function around zero value of $\omega_\alpha/kv_{T_\alpha}$ gives the solution of (4):

$$\bar{\gamma} = \frac{2}{\sqrt{\pi}}\left(\frac{\rho_e}{L}\right)^{3/2}\frac{1-(kL)^2}{kL} = \bar{\gamma}_{0e} \tag{6}$$

This is classical growth rate of the electron tearing mode in absence of collisions and shear (see, e.g., Galeev and Zelenyi 1976).

For very strong collisions $\bar{\nu}_e \sim \nu_i/kv_{T_i} \gg 1$ we can expand $Z_{0\alpha}$ with $\omega_\alpha/kv_{T_\alpha} \gg 1$ and obtain equation for $\bar{\gamma}$

$$\bar{\gamma} = \bar{\gamma}_{0e}\bar{\nu}_e\left(1 + \frac{\Delta_i}{R_e}\left(1 + \bar{\nu}_e^{-2}\mu^{-1/2} + \bar{\gamma}\mu^{-1}\bar{\nu}_e^{-1}\right)^{-1}\right)^{-1} \tag{7}$$

where $\Delta_i = R_i$ if the mean free path $\lambda_{ei} = v_{T_e}/\nu_{ei}$ is larger than R_i, while $\Delta_i \sim \bar{\nu}_e\mu^{-1/2}\rho_e > R_i$, if the mean free path λ_{ei} is smaller than R_i. Here, therefore, we have three solutions of (7). Together with (6) we have

$$\bar{\gamma} = \bar{\gamma}_{0e}\begin{cases} 1, & \bar{\nu}_e \ll 1 \\ \bar{\nu}_e, & 1 \ll \bar{\nu}_e \ll \mu^{-1/8}, \quad \rho_e/L < \mu^{3/4} \\ \mu^{-1/4}/\bar{\nu}_e, & \mu^{-1/8} \ll \bar{\nu}_e \ll \mu^{-1/4}, \quad \rho_e/L < \mu \\ (\mu L/\rho_e)^{1/2}, & \mu^{1/4}(L/\rho_e)^{1/2} < \bar{\nu}_e \end{cases} \tag{8}$$

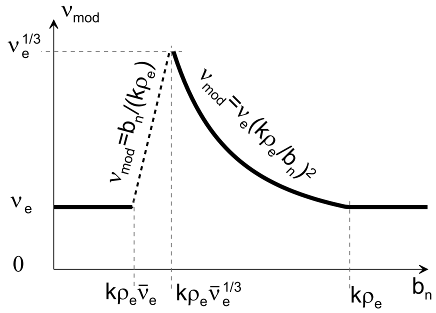

Fig. 1 Modified collision frequency as function of $b_n = B_z/B_0$. Figure is adopted from (Zelenyi and Taktakishvili 1981)

Equation (8) demonstrates how Coulomb (of effective) dissipation replaces the weak and sensitive Landau damping as a driving mechanism of spontaneous reconnection. This effect is especially important if Landau damping is switched off by the influence of a weak normal component of magnetic field and tearing mode becomes linearly stable (Schindler 1974; Galeev and Zelenyi 1976; Pellat et al. 1991; Quest et al. 1996). In this case the presence of collisions provides relatively slow, but persistent growth of reconnecting modes. Resistive mode (similar to the one described by (8) for $B_z = 0$ case, $\bar{\gamma} \sim \bar{\nu}_e$) emerges in this case, even for the modes stable in collisionless regime. These dissipative modes are clearly seen in numerical simulations (Lipatov and Zelenyi 1982), where numerical dissipation due to "shot noise" effects is unfortunately unavoidable.

4 Pitaevskii Effect

It is necessary to note, that the presence of finite normal component of magnetic field in the vicinity of the neutral plane $B_x \approx 0$ results in modification of the collision frequency (so called Pitaevskii effect, see Pitaevskii 1963). The nature of this effect is related with the kinetic character of collision process, which is missed both in τ-approximation and BGK-approach. Strictly speaking, collisions should be described by Landau collisional operator (Pitaevskii and Lifshitz 1981), where all details of distribution function become important (especially gradients of distribution function in a phase space). For short wavelength modes $k\rho_{en} > 1$ ($\rho_{en} = v_{T_e}/\Omega_{en}$ is electron Larmor radius in the vicinity of current sheet neutral plane with electron gyrofrequency in B_z field, Ω_{en}) perturbed distribution $f_{1\alpha} \sim \exp(ikx - i\omega t) \sim \exp(ik\rho_{en})\sin(\Omega_{en}t)$ becomes very inhomogeneous in the phase space and collisions act much more effectively to smoothen it. Pitaevskii took this effect into account and have shown that it could in a first approximation be reduced to the corresponding increase of collision frequency:

$$\bar{\nu}_e \to \bar{\nu}_{mod} = \bar{\nu}_e \times \begin{cases} k^2\rho_{en}^2, & k^2\rho_{en}^2 > 1 \\ 1, & k^2\rho_{en}^2 < 1 \end{cases}$$

Here we take into account that for current sheet geometry $\rho_{en} = \rho_e/b_n$, where $b_n = B_z/B_0$ is the dimensional value of the normal component of the magnetic field. The corresponding modification of ν_e can be found in Fig. 1 and have non-monotonous form.

Modified collision frequency ν_{mod} equals to ν_e for unmagnetized electrons, when $\nu_e < \Omega_{en}$. For strongly magnetized electrons $k\rho_{en} < 1$ (i.e. when the wavelength of perturbations becomes larger than electron gyroradius in B_z field) modified frequency also equals to ν_e. Effect of Pitaevskii starts working in the region with $k\rho_{en} > 1$. In this region modified collision frequency increases with decrease of b_n as $\nu_{mod} \sim b_n^{-2}$ until ν_{mod} is smaller than Ω_{en}. When $\nu_{mod} = \Omega_{en}$ (i.e. $b_n = k\rho_e \bar{\nu}_e^{1/3}$), effect of Pitaevskii stops working. We decided to remind about this effect, because it could be important even for a weakly collisional plasma in a weak magnetic field characteristic for interplanetary (interstellar) space and planetary magnetotails. Simple estimates of Reinolds numbers without taking into account the kinetics of collision process could significantly underestimate their role.

5 Effects of Magnetic Field Shear and Collisional Dissipation

Very often plasma configurations in space and laboratory have the additional component of magnetic field along the current supporting configuration. This component (toroidal in fusion devices) could have small (magnetotail) or large (magnetopause) values, so we will consider the general case when it could have an arbitrary value. In presence of a finite magnetic field B_y motion of particles in the neutral plane can become magnetized by this component. The critical value of B_y for such "magnetization" is defined as: $B_y^* = B_0\rho_\alpha/R_\alpha = B_0\sqrt{\rho_\alpha/L}$. If $B_y < B_y^*$ particles can be considered as unmagnetized, because Larmor radius in B_y is larger than the thickness of the central region of current sheet $|z| < R_\alpha$. For such weak B_y component its influence on system properties could be neglected. For $B_y > B_y^*$ all particles are magnetized (Galeev and Zelenyi 1978; Karimabadi et al. 2005). For these two regimes mechanisms of tearing mode growth are principally different (Drake and Lee 1977; Zelenyi and Taktakishvili 1987). When electrons get magnetized a finite dissipation due to Landau resonance interaction is replaced by the dissipation produced by electron inertia ($m_e \neq 0$).

Here we introduce dimensionless parameter $b_y = B_y/B_0$ and consider regimes of the tearing mode for various values of b_y. In contrast to the system with $B_y = 0$, tearing mode in the current sheet with magnetic field shear is very sensitive to any perturbation of the scalar potential φ (Coppi 1965; Galeev et al. 1986; Daughton and Karimabadi 2005). The spatial domain can be separated into two regions: (1) central region in the vicinity of so called singular surface which is the layer with $k_\parallel = k_x B_x(z)/|\mathbf{B}| = 0$. In this region perturbations of the electrostatic field $-\nabla_\parallel \varphi = -k_\parallel \varphi$ are small and can not compensate perturbation of the inductive field $-c^{-1}\partial A_\parallel/\partial t$, where $A_\parallel = A_{1y}(z)B_y/|\mathbf{B}|$. As a result, a finite electric field $E_\parallel = -\nabla_\parallel \varphi - c^{-1}\partial A_\parallel/\partial t$ exists in the vicinity of singular surface and frozen-in condition breaks down. (2) Outer region where inductive and potential parts of $E_\parallel$ compensate each other ($E_\parallel = 0$) and single fluid-approximation can be used. In the vicinity of the layer with $k_\parallel = 0$ equations for perturbed vector and scalar potentials for the general case with both shear ($b_y \neq 0$) and collisional effects ($\nu \neq 0$) taken into account can be written as (Zelenyi and Taktakishvili 1987)

$$\frac{d^2\varphi}{dz^2} = G(z), \qquad \frac{d^2 A_\parallel}{dz^2} = \frac{\rho_i^2}{2R_i^2 b_y^2}\frac{\omega}{k_\parallel c}G(z) \tag{9}$$

where ω is frequency of perturbation and

$$G(z) = \left(\varphi(z) - \frac{\omega A_\parallel}{k_\parallel c}\right)\frac{2R_i^2}{b_y^2\rho_i^2}\sum_\alpha \frac{Z_{1\alpha}(1+X_\sigma)}{R_\alpha^2 D_\alpha}, \quad \alpha \neq \sigma$$

$$X_\sigma = Z_{1\alpha}\frac{2i\nu_{\alpha\sigma}\omega}{(k_\parallel v_{T_\alpha})^2 D_\alpha}, \quad \alpha \neq \sigma$$

$$D_\alpha = 1 + \frac{i\nu_\alpha Z_{0\alpha}}{k_\parallel v_{T_\alpha}} + \frac{2i\nu_{\alpha\sigma}\omega Z_{1\alpha}}{(k_a v_{T_\alpha})^2}$$

Argument of $Z_{n\alpha}$ function is $\omega/k_\parallel v_{T_\alpha}$. To derive system of (9) we took into account several assumptions: (1) we neglect particle drift v_{D_α} in the vicinity of the layer $k_\parallel = 0$; (2) we assume that $d^2/dz^2 \gg k^2$ and $d^2/dz^2 \gg V_0(z)$; (3) in the vicinity of the layer $k_\parallel = 0$ we assume $A_\parallel \approx const \neq 0$, but we keep terms $d^2A_\parallel/dz^2$; (4) we assume that $\rho_i/b_y \gg R_i$, R_e ($B_y < B_y^*$), i.e. $b_y \ll \sqrt{\rho_i/L}$ (but $b_y \gg \sqrt{\rho_e/L}$ for an external solution).

One can get good physical insight to the problem considering the Doppler-shift of perturbation for collisionless regime with $\gamma > \nu_e$: $\omega' = \omega - \omega_D$, where $\omega_D = k_\parallel v_\parallel \approx k_\parallel v_{T_e}$. While for collisional case $\gamma < \nu_e$ particles motion resembles the diffusion along magnetic field lines. In this case Doppler-shift can be written as $\omega_D \approx k_\parallel^2 v_{T_e}^2/\nu_e$ (Drake and Lee 1977). If the value of Doppler-shift is much smaller than time scale of electric field variation ($\omega_D \ll \omega$), particle can be accelerated by $E_\parallel$ in the vicinity of the layer $k_\parallel = 0$. Condition $\omega_D \ll \omega$ defines the width of the singular region:

$$\Delta_s = \begin{cases} \Delta_s^0 = \bar\gamma b_y L, & \gamma \gg \nu_e \\ \Delta_s^c = \sqrt{\bar\gamma \bar\nu_e} b_y L, & \gamma \ll \nu_e \end{cases}$$

System (9) determines the dispersion relation valid for $\Delta_s \ll L$:

$$\frac{1-(kL)^2}{kL} = \frac{L}{A_\parallel}\int_{-\infty}^{+\infty} \frac{d^2A_\parallel}{dz^2} dz \tag{10}$$

This dispersion relation determines the growth rates for all regimes of the tearing mode depending on b_y and ν_e. The important role is played by relation between scales Δ_s and δ_φ, where δ_φ defines the scale of φ variation, i.e. perturbations of vector potential $A_\parallel$ are not compensated by perturbations of the scalar potential in the domain $|z| < \delta_\varphi$. For systems with $\Delta_s < \delta_\varphi$ kinetic regime of tearing mode is provided by resonant collisionless or collisional interaction with particles in the region $|z| < \Delta_s$ (see left panel of the scheme in Fig. 2). Electrostatic effects become important already outside the region of strong interaction of waves with electrons $|z| < \Delta_s$. For $|z| > \Delta_s$ Doppler shift $\sim k_\parallel(z)v_\parallel \sim k_\parallel(z)v_{T_e}$ strongly reduces the resulting value of the perturbed current $d^2A_\parallel/dz^2 \sim j_\parallel$.

For the opposite case $\Delta_s > \delta_\phi$ (see right panel of the scheme in Fig. 2) electrostatic effects control the evolution of the system because the width of interaction region depends on the width of the domain, where the frozen in condition $E_\parallel = i\omega A_\parallel/c - ik_\parallel\varphi = 0$ is violated. For $|z| > \delta_\varphi$ $E_\parallel \to 0$ and interaction for the cases with $\Delta_s > \delta_\phi$ occurs in MHD regime, when the dissipation could be provided either by collisional ($\sim\nu_e$) or inertial ($\sim m_e$) resistivities.

Below we consider two different regimes of the tearing mode: (1) MHD regime, when inertia or resistivity produce perturbation of current density with spatial scale exceeding ion Larmor radius $\rho_{yi} = \rho_i/b_y$ ($\delta_\varphi \sim \rho_{yi} < \Delta_s$); (2) kinetic regime, when spatial scales of current perturbation are smaller than ion Larmor radius ($\delta_\varphi \sim \rho_{yi} > \Delta_s$).

5.1 Collisionless Systems

If collisional frequency is small ($\nu_e \ll \gamma$) one can neglect the real part of the frequency of perturbation ($\mathrm{Re}\omega = 0$, $\mathrm{Im}\omega = \gamma$) and consider only electron input to the growth of perturbations. In this case system (9) takes a form

$$\frac{d^2\varphi}{dz^2} = \frac{2T_i}{\rho_{yi}^2 T_e}\left(\varphi - \frac{\omega}{k_\parallel c}A_\parallel\right)Z_{1e}, \qquad \frac{d^2A_\parallel}{dz^2} = \frac{T_i}{R_i^2 T_e}\frac{\omega}{k_\parallel c}\left(\varphi - \frac{\omega}{k_\parallel c}A_\parallel\right)Z_{1e} \tag{11}$$

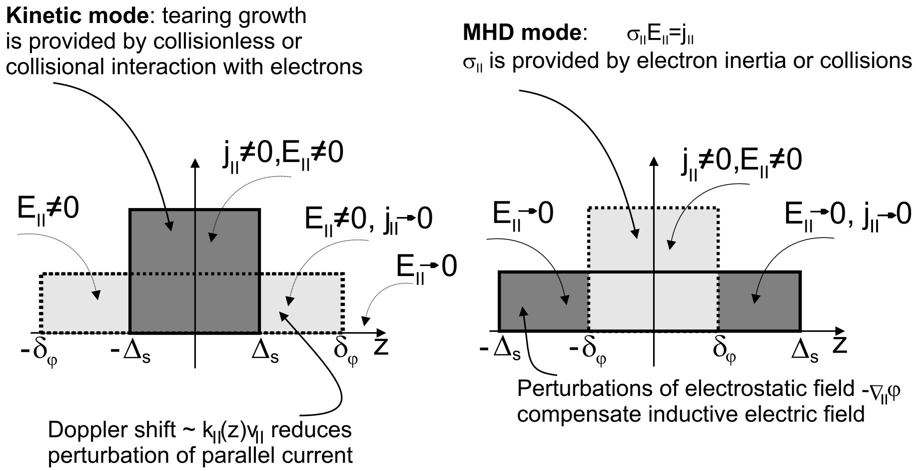

Fig. 2 Schematic view of relation between spatial scales Δ_s and δ_ϕ in kinetic and MHD regimes

For kinetic regime ($\varphi \sim 0$) we could substitute expression for $d^2A_\parallel/dz^2$ (11) into dispersion relation (10) with additional condition $Z_{1e} \approx 1$ valid for $|z| \gg \Delta_s^0$ to obtain growth rate (see Laval et al. 1966)

$$\bar{\gamma}_{01} = \frac{1}{\sqrt{\pi}}\left(1+\frac{T_i}{T_e}\right)\frac{1-(kL)^2}{kL}\left(\frac{\rho_e}{L}\right)^2\frac{1}{b_y}$$

First equation of system (11) gives estimate of the spatial scale $\delta_\varphi \approx \rho_{yi}\sqrt{T_i/2T_e}$. Then limit of this regime is defined by equation $\delta_\varphi \approx \Delta_s$: $b_y < (L/\rho_e)\sqrt{1/2\mu}$.

For MHD regime $\delta_\varphi < \Delta_s$ we can use expansion $Z_{1e} \approx (1/2)(z/\Delta_s)^2$ in the region $|z| < \Delta_s^0$. Substitution of expression for $d^2A_\parallel/dz^2$ (11) into dispersion relation (10) gives (Zelenyi and Taktakishvili 1987)

$$\bar{\gamma}_{02} = \sqrt{\mu}\left(\frac{\rho_e}{L}\right)^3\left(4\left(1+\frac{T_i}{T_e}\right)\frac{T_i}{T_e I}\frac{1-(kL)^2}{kL}\right)^2 \tag{12}$$

with $I = 2\pi\Gamma(3/4)/\Gamma(1/4)$. It is worth to notice, that these two regimes match at demagnetization point b_y^*, i.e. the ratio $\bar{\gamma}_{01}/\bar{\gamma}_{01}$ is some constant around unity, when $b_y = b_y^* = (L/\rho_e)\sqrt{1/2\mu}$ (Galeev and Zelenyi 1977). For the first time this instability (inertial MHD tearing mode) was found in the early paper by Coppi (1965). One can see that the such mode could exist only in the very exotic case $b_y > \sqrt{m_i/m_e}$ ($L \sim \rho_i$) or equivalently for extremely small plasma beta $\beta < m_e/m_i$.

5.2 Collisional Systems

For the case with strong collisions $\nu_\alpha/k_\parallel v_{T_\alpha} \gg 1$ system (9) can be rewritten as

$$\begin{aligned}\rho_{yi}^2\frac{d^2\varphi}{dz^2} &= \frac{2T_i}{T_e}\left(\varphi - \frac{\omega}{k_\parallel c}A_\parallel\right)\frac{k_\parallel^2 v_{T_e}^2}{2\gamma\nu_e + k_\parallel^2 v_{T_e}^2}\\ R_e^2\frac{d^2A_\parallel}{dz^2} &= \frac{\omega}{k_\parallel c}\left(\varphi - \frac{\omega}{k_\parallel c}A_\parallel\right)\frac{k_\parallel^2 v_{T_e}^2}{2\gamma\nu_e + k_\parallel^2 v_{T_e}^2}\end{aligned} \tag{13}$$

For kinetic regime ($\delta_\varphi > \Delta_s$ and $\varphi \to 0$) the solution of the first equation of system (13) in the region $|z| > \Delta_s^0$ is the same as the solution of the first equation of collisionless system (11). As a result, we obtain estimates of $\delta_\varphi \approx \rho_{yi}\sqrt{T_i/2T_e}$. Second equation of system (13) after substitution into (10) with $\varphi = 0$ gives the growth rate for so called semi-collision mode (Drake and Lee 1977):

$$\bar{\gamma}_{sc} = \bar{\nu}_e^{1/3}\left(\frac{\rho_e}{L}\right)^{4/3}\frac{1}{\pi^{1/3}b_y^{2/3}}\left(\left(1+\frac{T_i}{T_e}\right)\frac{1-(kL)^2}{kL}\right)^{2/3}$$

Growth rate $\bar{\gamma}_{sc}$ matches $\bar{\gamma}_{01}$ at $\bar{\nu}_e = (\rho_e/L)^2/b_y$.

For MHD regime ($\delta_\varphi < \Delta_s$) we can obtain the solution of the first equation of system (13):

$$\varphi(z) = \frac{\omega}{k_\parallel c}A_\parallel\frac{z^2}{4\delta_\varphi\Delta_s^c}\int_0^{\pi/2}\sqrt{\sin\theta}\exp\left(-\frac{z^2\cos\theta}{4\delta_\varphi\Delta_s^c}\right)d\theta$$

Substituting $\varphi(z)$ into the second equation of system (13) we obtain the expression for $A_\parallel$, which can be substituted into (10). As a result, we obtain growth rate of the well-known resistive Furth-Killeen-Rosenbluth mode (Furth et al. 1963):

$$\bar{\gamma}_{FKR} = \bar{\nu}_e^{3/5}\left(\frac{\rho_e}{L}\right)^{6/5}\mu^{1/5}\left(\frac{1-(kL)^2}{kL}\frac{2(T_i+T_e)}{T_eI}\right)^{4/5}$$

with $I = 2\pi\Gamma(3/4)/\Gamma(1/4)$. Growth rates γ_{sc} and γ_{FKR} match at $\Delta_s^c \sim \rho_{yi}$.

5.3 Role of B_y

Collisionless growth rate for resonant tearing mode in neutral current sheet with $B_y = 0$ was estimates as $\bar{\gamma}_{0e} \sim (\rho_e/L)^{3/2}$ (Coppi et al. 1966b). Therefore, estimates of collisionless growth rate for inertial mode $\bar{\gamma}_{01} \sim (\rho_e/L)^2/b_y$ with $B_y \neq 0$ becomes equal to $\bar{\gamma}_{0e}$ for $b_y = (\rho_e/L)^{1/2} \ll 1$. This value of b_y corresponds to magnetization of electrons ($B_y = B_y^*$), i.e. $\rho_e/b_y = R_e$ for $b_y = (\rho_e/L)^{1/2}$. For system with $b_y > (\rho_e/L)^{1/2}$ growth rate is determined by electron inertial resistivity and described by expression $\gamma_{01} \sim (\rho_e/L)^2/b_y$ until frequency of collisions is small enough $\bar{\nu}_e < \bar{\gamma}_{01}$. When $\bar{\nu}_e = \bar{\gamma}_{01}$ and $b_y > (\rho_e/L)^{1/2}$ (i.e. $\bar{\nu}_e > (\rho_e/L)^{3/2}$) the semi-collisional regime, where inertial resistivity is replaced by the collisional one (Drake and Lee 1977) establishes with $\bar{\gamma}_{sc} \sim \bar{\nu}_e^{1/3}/b_y^{2/3}$. As we mention above for very large $b_y > (L/\rho_e)\sqrt{1/2\mu}$ collisionless kinetic inertial mode transforms to MHD inertial mode and growth rates also match quite well at $B_y = B_y^*$. Therefore, we have dependence of the growth rate on b_y, where growth rate for intermediate regime $b_y \sim (\rho_e/L)^{1/2}$ can be obtained only by numerical solution of the corresponding dispersion equation (see Zelenyi and Taktakishvili 1987). So, we see that the interplay of different mictroscales of spontaneous reconnection process (scales of resonant or collisional electron interaction Δ_s, scale of the violation of the frozen in condition, δ_ϕ) determines the real modes and mechanisms of its operating.

6 Discussion and Conclusions

The general character of the growth rates for spontaneous reconnection modes as function of ν_e and b_y is shown in Fig. 3, where we also indicate parameter regions for various current sheets in different conditions existing in space and laboratory plasmas. Magnetotail current sheet is typically characterized by small value of magnetic shear $b_y \ll 1$ (Petrukovich 2011). The level of high-frequency wave activity responsible for effective collisions is also weak in this region (Coroniti 1985; Eastwood et al. 2009; Fujimoto et al. 2011). Therefore "MT" domain is situated at a lower left corner of Fig. 3 in $\bar{\nu}_e \ll 1$, $b_y \ll 1$ region. However, this domain could be extended by including current sheets observed under active conditions. For example, in the vicinity of the reconnection regions, where secondary X-lines can be formed due to the tearing instability of current sheet located in the outflow region, component B_y could be relatively strong ($B_y \sim B_0$, see, e.g., Nakamura and et al. 2008; Wang and et al. 2012). Moreover, in case of strong B_y electric field fluctuations related to flows of accelerated particles are often observed in outflow region supporting increase of effective collisions ν_e (Huang and et al. 2012) in agreement with theoretical estimates (Yoon and Lui 2006). However, for weak values of B_y the intensity of these effective collisions is low to be responsible for reconnection (Eastwood et al. 2009), and corresponding MT domain can be expanded to semicollisional region only for sufficiently large B_y. Here we also can mention alternative source of effective conductivity corresponding to stochastic ion and electron motion in the current sheets (Horton and Tajima 1990; Numata and Yoshida 2002). For magnetopause current sheet shear of magnetic field is often strong enough to provide $b_y > 1$ (Berchem and Russell 1982; Panov et al. 2008). Therefore, electrons and ions are magnetized in the vicinity of the neutral plane by B_y. In this case "MP" domain corresponds to the inertial mode with $\gamma = \gamma_{01}$. However, similar to "MT" domain effective collisions due to lower-hybrid and/or ion-cyclotron turbulence (Labelle and Treumann 1988; Panov et al. 2006) could expand "MP" domain up to semi-collisional regime of the tearing mode.

Current sheets detected in Solar corona ("SC" domain) correspond to strong (but finite) shear $m_i/m_e > b_y > 1$ and weak, but finite, collisions (Priest and Forbes 2000; Uzdensky 2003; Birn and Priest 2007). Development of the semi-collision tearing mode in these current sheets results in spontaneous magnetic reconnection (initiating the onset of Solar flares) and the subsequent electron acceleration. Here effective collisions due to high-frequency turbulence could also contribute to the growth rate and help to destabilize current sheet (Büchner 2007). We emphasize that the process of spontaneous reconnection in Solar corona and upper Solar atmosphere is mostly kinetic. Neither the value of shear component b_y, nor the degree of collisionality are strong enough to support this process to be accomplished in MHD regime.

Laboratory devices with relatively cool plasma, where magnetic field configurations with current sheet are produced, are located in Fig. 3 in the domain with strong electron collisions (Frank 2010; Yamada et al. 2010; Frank et al. 2011) enhanced by effective collisions (Ji et al. 2004). Magnetic reconnection in laboratory current sheets due to growth of collisional Furth-Killeen-Rosenbluth tearing mode (Furth et al. 1963) are often observed and described in details (see, e.g., Frank 2010). Moreover, laboratory devices can operate with relatively strong magnetic shear $B_y \sim B_0$ induced initially and growing with development of current sheet (Frank et al. 2005).

Finally tokamaks ("TK" domain) with high-temperature plasma are characterized by strong toroidal field (Wesson 2004; Steinhauer 2011) (strong shear $b_y \gg 1$ in our notations) and moderately strong collisions. Although the degree of collisionality could be enhanced by turbulence (see Budaev et al. 2011), tokamak domain most probably is located as semicollisional regime at Fig. 3 especially for future devices for real hot fusion plasma confinement. For tokamaks tearing instability plays important, although undesirable, role of destruction of magnetic surface (see review Boozer 2012b, and references therein).

Although, the analysis of reconnection mechanisms presented above is more appropriate for description of reconnection in space plasma, here we would like to discuss briefly the comparison with reconnection processes occurring in laboratory and tokamak experiments. The theory of collisional reconnection in tokamaks is originated from paper by Kadomtsev (1975) (see also review Kadomtsev 1987), where the model of resistive reconnection was developed. This model predicts reconnection time $\sim \nu_e^{-1/2}$ and is unable to explain the powerful sawtooth instability related to fast reconnection (von Goeler et al. 1974). Physics of this fast tokamak reconnection is essentially similar to models with large B_y presented in this review. It is believed that $m = 1$ modes in tokamaks results in formation of magnetic islands in the vicinity of a singular layer $\mathbf{k} \cdot \mathbf{B} = 0$ (Rutherford 1973; Rosenbluth et al. 1973; Zakharov 1980). Here the principal role is played by the electron inertia (Wesson 1990), because plasma beta is small enough ($\beta < m_i/m_e$) due to the large shear component of the magnetic field. Weak level of collisions in tokamak plasma leads to the dominance of this inertial mode, which can describe fast sawtooth reconnection (Porcelli 1991). Inertial $m = 1$ mode in tokamaks in principal corresponds to the current sheet thickness of the order of electron inertial length $d \sim m_e^{1/2}$ and develops with the growth rate $\gamma_{m=1} \sim \sqrt{\mu}(\rho_i/d) \sim \sqrt{\mu}$. The same estimates can be obtained for collisionless inertial tearing regime $\gamma_{02} \sim \sqrt{\mu}$ for $L \sim d \sim \sqrt{m_e}$ (see (12)).

Important role of additional effective collisions (or anomalous diffusion) for inertial $m = 1$ mode was considered by Drake and Kleva (1991). Authors have shown that stabilization of $m = 1$ mode due to diamagnetic drifts (effect similar to Doppler-shift effect for collisionless mode, see discussion above and in Drake and Lee (1977)) results in significal reduction of the growth rate. At the same time effective diffusion due to drift instability can provide the increase of reconnection rate. The stabilization effect of diamagnetic drifts for $m = 1$ mode was confirmed by experimental observation (Levinton et al. 1994) and numerical modeling (Zakharov et al. 1993). Other possible candidate for the increase of the tearing growth rate in tokamaks reconnection with large shear magnetic component is the gradient of electron pressure along field lines (Aydemir 1992; Grasso et al. 1999). Presence of the finite electron compressibility results in appearance of nonvanishing parallel electric field in the vicinity of the singular layer. In this case the structure of reconnection region resembles the one shown in Fig. 2 (left panel), where $j_\parallel \neq 0$ domain is embedded into $E_\parallel \neq 0$ domain (Kleva et al. 1995).

General model of two-fluid magnetic reconnection in tokamaks with two limits ($\beta < m_e/m_i$ and $\beta \sim 1$) can be found in Biskamp et al. (1997). In case of small plasma beta electron inertia plays the most important role, while large-β regime corresponds to separation of electron and ion motions and Hall reconnection. The comprehensive review by Porcelli et al. (2002) can be used to obtain more detailed information about inertial and Hall modes of magnetic reconnection in tokamaks, while papers by Park et al. (2006b, 2006a); Igochine et al. (2007) contain comparison of theoretical predictions and experimented observations.

Substantial difference between tearing modes developed in space and laboratory plasmas is provided by the difference of boundary conditions. The traditional approach to growth rate calculations consists in matching of solutions of perturbed Vlasov-Maxwell equations at the boundary separating inner region around the singular layer $\mathbf{k} \cdot \mathbf{B} = 0$ and outer region, where resonant wave-particle interaction or inertial effects can be neglected. Therefore, to determine solutions in the outer region one needs to introduce the certain external boundary conditions. The most appropriate approach for space systems consists in consideration of infinitely distant boundaries with corresponding solutions quickly decreasing with distance from the singular layer (an example of alternative approach can

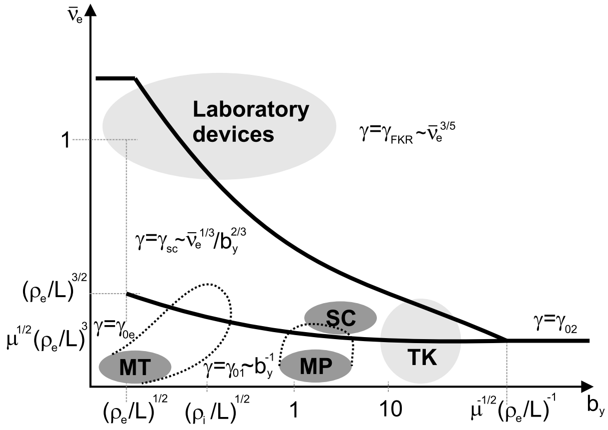

Fig. 3 Regimes of tearing mode: MT denotes magnetotail, MP denotes magnetopause, SC denotes solar corona, and TK denotes tokamaks. Figure is adopted from Zelenyi and Taktakishvili (1981)

be found in Zelenyi and Kuznetsova 1984). Situation is different for tokamaks configurations, where outer boundaries are accessible and have well defined physical properties like infinite conductivity (Coppi et al. 1966a; Wesson 1966). In this case, the set of external solutions is fully controlled by system geometry (see, e.g., Mikhailovskii 1978; Pegoraro and Schep 1986, and references therein). Additionally, characteristic cylindrical-like geometry of the tokamak system corresponds to appearance of local singularities of solutions (Newcomb 1960). Such singularities are absent in simplified plane geometry typical for space systems. These two problems are not encountered in major of space-plasma systems. Therefore, further comparison between spacecraft observations (and corresponding theories) with tokamak and laboratory reconnections requires accurate consideration of the geometry issue (see discussion in Boozer 2012a).

In conclusion, we can mention that for the major part of observed current sheets the semi-collisions regime of spontaneous reconnection seems to play the most important role. This regime principally cannot be described in a frame of MHD approach, until shear of magnetic field becomes unrealistically strong. On the other hand, unlikely that pure kinetic mode with unmagnetized electrons could be realized in realistic systems due to electron magnetization by even very weak magnetic fields. As a result, it is principal that the regime of current sheet destruction in the course of magnetic reconnection should be described in a frame of kinetic models with careful taking into account effects of collisions, which also exist in many seemingly collisionless configurations as effective collisions due to scattering of electrons at microturbulence fluctuations. In addition it should be kept in mind that weak magnetization ($k\rho_\alpha > 1$) of particle trajectories could substantially enhance collision frequencies formally defined in simplified τ- or BGK descriptions due to kinetic properties of the exact collisional operator. This effect known as Pitaevskii one could occur in a wide parameter range.

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References

V. Angelopoulos, J.P. McFadden, D. Larson, et al., Tail reconnection triggering substorm onset. Science **321**, 931–935 (2008). doi:10.1126/science.1160495

A.Y. Aydemir, Nonlinear studies of $m = 1$ modes in high-temperature plasmas. Phys. Fluids B **4**, 3469–3472 (1992). doi:10.1063/1.860355

J. Berchem, C.T. Russell, Magnetic field rotation through the magnetopause—ISEE 1 and 2 observations. J. Geophys. Res. **87**, 8139–8148 (1982). doi:10.1029/JA087iA10p08139

P.L. Bhatnagar, E.P. Gross, M. Krook, A model for collision processes in gases. I. Small amplitude processes in charged and neutral one-component systems. Phys. Rev. **94**, 511–525 (1954). doi:10.1103/PhysRev.94.511

J. Birn, E.R. Priest, *Reconnection of Magnetic Fields: Magnetohydrodynamics and Collisionless Theory and Observations* (2007)

D. Biskamp, *Magnetic Reconnection in Plasmas* (2000)

D. Biskamp, E. Schwarz, J.F. Drake, Two-fluid theory of collisionless magnetic reconnection. Phys. Plasmas **4**, 1002–1009 (1997). doi:10.1063/1.872211

A.H. Boozer, Magnetic reconnection in space. Phys. Plasmas **19**(9), 092902 (2012a). doi:10.1063/1.4754715

A.H. Boozer, Theory of Tokamak disruptions. Phys. Plasmas **19**(5), 058101 (2012b). doi:10.1063/1.3703327

J. Büchner, Astrophysical reconnection and collisionless dissipation. Plasma Phys. Control. Fusion **49**, 325–339 (2007). doi:10.1088/0741-3335/49/12B/S30

V.P. Budaev, S.P. Savin, L.M. Zelenyi, Investigation of intermittency and generalized self-similarity of turbulent boundary layers in laboratory and magnetospheric plasmas: towards a quantitative definition of plasma transport features. Phys. Usp. **54**, 875–918 (2011). doi:10.3367/UFNe.0181.201109a.0905

B. Coppi, Current-driven instabilities in configurations with sheared magnetic fields. Phys. Fluids **8**, 2273–2280 (1965). doi:10.1063/1.1761195

B. Coppi, J.M. Greene, J.L. Johnson, Resistive instabilities in a diffuse linear pinch. Nucl. Fusion **6**, 101–117 (1966a)

B. Coppi, G. Laval, R. Pellat, Dynamics of the geomagnetic tail. Phys. Rev. Lett. **16**, 1207–1210 (1966b). doi:10.1103/PhysRevLett.16.1207

F.V. Coroniti, Space plasma turbulent dissipation—reality or myth? Space Sci. Rev. **42**, 399–410 (1985). doi:10.1007/BF00214995

W. Daughton, The unstable eigenmodes of a neutral sheet. Phys. Plasmas **6**, 1329–1343 (1999). doi:10.1063/1.873374

W. Daughton, H. Karimabadi, Kinetic theory of collisionless tearing at the magnetopause. J. Geophys. Res. (Space Phys.) **110**, 3217 (2005). doi:10.1029/2004JA010751

X.H. Deng, H. Matsumoto, Rapid magnetic reconnection in the Earth's magnetosphere mediated by whistler waves. Nature **410**, 557–560 (2001). doi:10.1038/35069018

M. Dobrowolny, Instability of a neutral sheet. Nuovo Cim., B **55**, 427–442 (1968). doi:10.1007/BF02711653

J.F. Drake, R.G. Kleva, Collisionless reconnection and the sawtooth crash. Phys. Rev. Lett. **66**, 1458–1461 (1991). doi:10.1103/PhysRevLett.66.1458

J.F. Drake, Y.C. Lee, Kinetic theory of tearing instabilities. Phys. Fluids **20**, 1341–1353 (1977). doi:10.1063/1.862017

J.P. Eastwood, T.D. Phan, S.D. Bale, A. Tjulin, Observations of turbulence generated by magnetic reconnection. Phys. Rev. Lett. **102**(3), 035001 (2009). doi:10.1103/PhysRevLett.102.035001

A.G. Frank, Dynamics of current sheets underlying flare-type events in magnetized plasmas. Phys. Usp. **53**, 941–947 (2010). doi:10.3367/UFNe.0180.201009h.0982

A.G. Frank, N.P. Kyrie, S.N. Satunin, Plasma dynamics in laboratory-produced current sheets. Phys. Plasmas **18**(11), 111209 (2011). doi:10.1063/1.3647576

A.G. Frank, S.Y. Bogdanov, V.S. Markov, G.V. Ostrovskaya, G.V. Dreiden, Experimental study of plasma compression into the sheet in three-dimensional magnetic fields with singular X lines. Phys. Plasmas **12**(5), 052316 (2005). doi:10.1063/1.1896376

M. Fujimoto, I. Shinohara, H. Kojima, Reconnection and waves: a review with a perspective. Space Sci. Rev. **160**, 123–143 (2011). doi:10.1007/s11214-011-9807-7

H.P. Furth, The 'mirror instability' for finite particle gyro-radius. Nucl. Fusion **1**, 169–174 (1962)

H.P. Furth, J. Killeen, M.N. Rosenbluth, Finite-resistivity instabilities of a sheet pinch. Phys. Fluids **6**, 459–485 (1963)

A.A. Galeev, Reconnection in the magnetotail. Space Sci. Rev. **23**, 411–425 (1979). doi:10.1007/BF00172248

A.A. Galeev, L.M. Zelenyi, Tearing instability in plasma configurations. Sov. Phys. JETP **43**, 1113 (1976)

A.A. Galeev, L.M. Zelenyi, Model of magnetic-field reconnection in a plane layer of collisionless plasma. JETP Lett. **25**, 380 (1977)

A.A. Galeev, L.M. Zelenyi, Magnetic reconnection in a space plasma, in *Theoretical and Computational Plasma Physics* (1978), pp. 93–116
A.A. Galeev, M.M. Kuznetsova, L.M. Zelenyi, Magnetopause stability threshold for patchy reconnection. Space Sci. Rev. **44**, 1–41 (1986). doi:10.1007/BF00227227
R.G. Giovanelli, Magnetic and electric phenomena in the Sun's atmosphere associated with sunspots. Mon. Not. R. Astron. Soc. **107**, 338 (1947)
D. Grasso, F. Pegoraro, F. Porcelli, F. Califano, Hamiltonian magnetic reconnection. Plasma Phys. Control. Fusion **41**, 1497–1515 (1999). doi:10.1088/0741-3335/41/12/306
E.G. Harris, On a plasma sheet separating regions of oppositely directed magnetic field. Nuovo Cimento **23**, 115–123 (1962)
W. Horton, T. Tajima, Decay of correlations and the collisionless conductivity in the geomagnetic tail. Geophys. Res. Lett. **17**, 123–126 (1990). doi:10.1029/GL017i002p00123
S.Y. Huang, et al., Observations of turbulence within reconnection jet in the presence of guide field. Geophys. Res. Lett. **39**, 11104 (2012). doi:10.1029/2012GL052210
J.D. Huba, N.T. Gladd, K. Papadopoulos, The lower-hybrid-drift instability as a source of anomalous resistivity for magnetic field line reconnection. Geophys. Res. Lett. **4**, 125–126 (1977). doi:10.1029/GL004i003p00125
V. Igochine, O. Dumbrajs, H. Zohm, A. Flaws, ASDEX Upgrade Team, Stochastic sawtooth reconnection in ASDEX upgrade. Nucl. Fusion **47**, 23–32 (2007). doi:10.1088/0029-5515/47/1/004
H. Ji, S. Terry, M. Yamada, R. Kulsrud, A. Kuritsyn, Y. Ren, Electromagnetic fluctuations during fast reconnection in a laboratory plasma. Phys. Rev. Lett. **92**(11), 115001 (2004). doi:10.1103/PhysRevLett.92.115001
B.B. Kadomtsev, Disruptive instability in tokamaks. Sov. J. Plasma Phys. **1**, 710–715 (1975)
B.B. Kadomtsev, Review article: magnetic field line reconnection. Rep. Prog. Phys. **50**, 115–143 (1987). doi:10.1088/0034-4885/50/2/001
H. Karimabadi, W. Daughton, K.B. Quest, Physics of saturation of collisionless tearing mode as a function of guide field. J. Geophys. Res. **110**, 3214 (2005). doi:10.1029/2004JA010749
R.G. Kleva, J.F. Drake, F.L. Waelbroeck, Fast reconnection in high temperature plasmas. Phys. Plasmas **2**, 23–34 (1995). doi:10.1063/1.871095
A.P. Kropotkin, H.V. Malova, M.I. Sitnov, Self-consistent structure of a thin anisotropic current sheet. J. Geophys. Res. **102**, 22099–22106 (1997). doi:10.1029/97JA01316
J. Labelle, R.A. Treumann, Plasma waves at the dayside magnetopause. Space Sci. Rev. **47**, 175–202 (1988). doi:10.1007/BF00223240
G. Lapenta, J.U. Brackbill, A kinetic theory for the drift-kink instability. J. Geophys. Res. **102**, 27099–27108 (1997). doi:10.1029/97JA02140
G. Laval, R. Pellat, M. Vuillemin, Instabilités Électromagnétiques des Plasmas Sans Collisions (cn-21/71), in *Plasma Physics and Controlled Nuclear Fusion Research, vol. II* (1966), pp. 259–277
F.M. Levinton, L. Zakharov, S.H. Batha, J. Manickam, M.C. Zarnstorff, Stabilization and onset of sawteeth in TFTR. Phys. Rev. Lett. **72**, 2895–2898 (1994). doi:10.1103/PhysRevLett.72.2895
A.S. Lipatov, L.M. Zelenyi, The study of magnetic islands dynamics. Plasma Phys. **24**, 1082–1089 (1982)
A.B. Mikhailovskii, Review of instability theory for high pressure Tokamak plasma, Technical report, 1978
R. Nakamura, et al., Cluster observations of an ion-scale current sheet in the magnetotail under the presence of a guide field. J. Geophys. Res. **113**, 7 (2008). doi:10.1029/2007JA012760
W.A. Newcomb, Hydromagnetic stability of a diffuse linear pinch. Ann. Phys. **10**, 232–267 (1960). doi:10.1016/0003-4916(60)90023-3
R. Numata, Z. Yoshida, Chaos-induced resistivity in collisionless magnetic reconnection. Phys. Rev. Lett. **88**(4), 045003 (2002). doi:10.1103/PhysRevLett.88.045003
E.V. Panov, et al., CLUSTER observation of collisionless transport at the magnetopause. Geophys. Res. Lett. **33**, 15109 (2006). doi:10.1029/2006GL026556
E.V. Panov, J. Büchner, M. Fränz, et al., High-latitude earth's magnetopause outside the cusp: cluster observations. J. Geophys. Res. **113**, 1220 (2008). doi:10.1029/2006JA012123
H.K. Park, A.J.H. Donné, N.C. Luhmann Jr., I.G.J. Classen, C.W. Domier, E. Mazzucato, T. Munsat, M.J. van de Pol, Z. Xia, Comparison study of 2D images of temperature fluctuations during sawtooth oscillation with theoretical models. Phys. Rev. Lett. **96**(19), 195004 (2006a). doi:10.1103/PhysRevLett.96.195004
H.K. Park, N.C. Luhmann Jr., A.J.H. Donné, I.G.J. Classen, C.W. Domier, E. Mazzucato, T. Munsat, M.J. van de Pol, Z. Xia, Observation of high-field-side crash and heat transfer during sawtooth oscillation in magnetically confined plasmas. Phys. Rev. Lett. **96**(19), 195003 (2006b). doi:10.1103/PhysRevLett.96.195003
E.N. Parker, Spontaneous current sheets in magnetic fields: with applications to stellar x-rays, in *Spontaneous Current Sheets in Magnetic Fields: With Applications to Stellar X-rays. International Series in Astronomy and Astrophysics*, vol. 1 (Oxford University Press, New York, 1994)

F. Pegoraro, T.J. Schep, Theory of resistive modes in the ballooning representation. Plasma Phys. Control. Fusion **28**, 647–667 (1986). doi:10.1088/0741-3335/28/4/003
R. Pellat, F.V. Coroniti, P.L. Pritchett, Does ion tearing exist? Geophys. Res. Lett. **18**, 143–146 (1991). doi: 10.1029/91GL00123
A.A. Petrukovich, Origins of plasma sheet B_y. J. Geophys. Res. **116**, 7217 (2011). doi:10.1029/2010JA016386
H.E. Petschek, Magnetic field annihilation. NASA Spec. Publ. **50**, 425 (1964)
L.P. Pitaevskii, Effect of collisions on perturbation of body rotating in plasma. J. Exp. Theor. Phys. **44**, 969–979 (1963) (in Russian)
L.P. Pitaevskii, E.M. Lifshitz, *Physical Kinetics*, Course of Theoretical Physics, vol. 10 (Pergamon Press, New York, 1981)
F. Porcelli, Collisionless $m = 1$ tearing mode. Phys. Rev. Lett. **66**, 425–428 (1991). doi:10.1103/PhysRevLett.66.425
F. Porcelli, D. Borgogno, F. Califano, D. Grasso, M. Ottaviani, F. Pegoraro, Recent advances in collisionless magnetic reconnection. Plasma Phys. Control. Fusion **44**, 389 (2002)
E. Priest, T. Forbes, *Magnetic Reconnection* (2000)
P.L. Pritchett, F.V. Coroniti, Plasma sheet disruption by interchange-generated flow intrusions. Geophys. Res. Lett. **381**, 10102 (2011). doi:10.1029/2011GL047527
K.B. Quest, H. Karimabadi, M. Brittnacher, Consequences of particle conservation along a flux surface for magnetotail tearing. J. Geophys. Res. **101**, 179–184 (1996). doi:10.1029/95JA02986
M.N. Rosenbluth, R.Y. Dagazian, P.H. Rutherford, Nonlinear properties of the internal $m = 1$ kink instability in the cylindrical tokamak. Phys. Fluids **16**, 1894–1902 (1973). doi:10.1063/1.1694231
M. Roth, J. de Keyser, M.M. Kuznetsova, Vlasov theory of the equilibrium structure of tangential discontinuities in space plasmas. Space Sci. Rev. **76**, 251–317 (1996). doi:10.1007/BF00197842
P.H. Rutherford, Nonlinear growth of the tearing mode. Phys. Fluids **16**, 1903–1908 (1973). doi:10.1063/1.1694232
K. Schindler, A theory of the substorm mechanism. J. Geophys. Res. **79**, 2803–2810 (1974). doi:10.1029/JA079i019p02803
K. Schindler, *Physics of Space Plasma Activity* (Cambridge University Press, Cambridge, 2006). doi: 10.2277/0521858976
K. Schindler, M. Soop, Stability of plasma sheaths. Phys. Fluids **11**, 1192–1195 (1968). doi: 10.1063/1.1692083
M.I. Sitnov, K. Schindler, Tearing stability of a multiscale magnetotail current sheet. Geophys. Res. Lett. **37**, 8102 (2010). doi:10.1029/2010GL042961
M.I. Sitnov, L.M. Zelenyi, H.V. Malova, A.S. Sharma, Thin current sheet embedded within a thicker plasma sheet: self-consistent kinetic theory. J. Geophys. Res. **105**, 13029–13044 (2000). doi: 10.1029/1999JA000431
B.U.Ö. Sonnerup, Adiabatic particle orbits in a magnetic null sheet. J. Geophys. Res. **76**, 8211–8222 (1971). doi:10.1029/JA076i034p08211
L.C. Steinhauer, Review of field-reversed configurations. Phys. Plasmas **18**(7), 070501 (2011). doi:10.1063/1.3613680
P.A. Sweet, The neutral point theory of solar flares, in *Electromagnetic Phenomena in Cosmical Physics*, ed. by B. Lehnert. IAU Symposium, vol. 6 (1958), p. 123
S.I. Syrovatskii, Dynamic dissipation of a magnetic field and particle acceleration. Sov. Astron. **10**, 270 (1966)
D.A. Uzdensky, Petschek-like reconnection with current-driven anomalous resistivity and its application to solar flares. Astrophys. J. **587**, 450–457 (2003). doi:10.1086/368075
S. von Goeler, W. Stodiek, N. Sauthoff, Studies of internal disruptions and $m = 1$ oscillations in tokamak discharges with soft-X-ray techniques. Phys. Rev. Lett. **33**, 1201–1203 (1974). doi:10.1103/PhysRevLett.33.1201
R. Wang, et al., Asymmetry in the current sheet and secondary magnetic flux ropes during guide field magnetic reconnection. J. Geophys. Res. **117**, 7223 (2012). doi:10.1029/2011JA017384
J. Wesson, Finite resistive instabilities of a sheet pinch. Nucl. Fusion **6**, 130–134 (1966)
J.A. Wesson, Sawtooth reconnection. Nucl. Fusion **30**, 2545–2549 (1990)
J. Wesson, *Tokamaks*, 3rd edn. (Oxford University Press, London, 2004). doi:10.1017/S0022377804003058
M. Yamada, R. Kulsrud, H. Ji, Magnetic reconnection. Rev. Mod. Phys. **82**, 603–664 (2010). doi:10.1103/RevModPhys.82.603
P.H. Yoon, A.T.Y. Lui, Quasi-linear theory of anomalous resistivity. J. Geophys. Res. **111**, 2203 (2006). doi: 10.1029/2005JA011482
L. Zakharov, B. Rogers, S. Migliuolo, The theory of the early nonlinear stage of $m = 1$ reconnection in tokamaks. Phys. Fluids B **5**, 2498–2505 (1993). doi:10.1063/1.860735

L.E. Zakharov, On the nature of disruptive instability in a tokamak. JETP Lett. **31**, 714 (1980)
L.M. Zelenyi, M.M. Kuznetsova, Large-scale instabilities of the plasma sheet driven by particle fluxes at the boundary of the magnetosphere. Sov. J. Plasma Phys. **10**, 326–334 (1984)
L.M. Zelenyi, A.L. Taktakishvili, The influence of dissipative processes on the development of the tearing mode in current sheets. Sov. J. Plasma Phys. **7**, 1064–1075 (1981)
L.M. Zelenyi, A.L. Taktakishvili, Spontaneous magnetic reconnection mechanisms in plasma. Astrophys. Space Sci. **134**, 185–196 (1987). doi:10.1007/BF00636466
L.M. Zelenyi, M.I. Sitnov, H.V. Malova, A.S. Sharma, Thin and superthin ion current sheets. Quasi-adiabatic and nonadiabatic models. Nonlinear Process. Geophys. **7**, 127–139 (2000)
L.M. Zelenyi, A.V. Artemyev, H.V. Malova, V.Y. Popov, Marginal stability of thin current sheets in the Earth's magnetotail. J. Atmos. Sol.-Terr. Phys. **70**, 325–333 (2008). doi:10.1016/j.jastp.2007.08.019
L.M. Zelenyi, A.V. Artemyev, K.V. Malova, A.A. Petrukovich, R. Nakamura, Metastability of current sheets. Phys. Usp. **53**, 933–941 (2010). doi:10.3367/UFNe.0180.201009g.0973
L.M. Zelenyi, H.V. Malova, A.V. Artemyev, V.Y. Popov, A.A. Petrukovich, Thin current sheets in collisionless plasma: equilibrium structure, plasma instabilities, and particle acceleration. Plasma Phys. Rep. **37**, 118–160 (2011). doi:10.1134/S1063780X1102005X