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Quasi-neutral sheet tearing instability induced by electron preferential acceleration from stochasticity

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Abstract. The role of tearing instability in substorm expansion onset is examined. From a consideration of observational constraints on the onset mechanism we conclude that this instability must have a considerable initial stage when the equilibrium magnetic field topology (without neutral lines) is still conserved. Moreover, this instability is shown to have no linear stage. Instead, either explicitly nonlinear or pseudolinear instability of negative energy eigenmodes can develop. The latter possibility may arise from electron preferential acceleration. This acceleration occurs for the electrons near the equilibrium chaos/adiabatic boundary that interact with the perturbation electric field and provide a means of dissipation to restore positive feedback for the instability development. We compute a set of single-particle orbits with realistic pitch angle distributions in a given tearing field model and evaluate statistical averages over the ensemble of particles and wave phases to verify the preferential acceleration of electrons. The resulting instability is analyzed by using a quasi-hydrodynamical multifluid model. Its amplitude can be less than that required for neutral line formation. It has a fast growth rate well within the timescales for substorm expansion onset. Its excitation leads to earthward acceleration of electrons and ions as well as local current amplification in the near-Earth region for excitation of cross-field current instability. In the theme of substorm as a low-dimension catastrophe for the current sheet, this nonlinear tearing instability approximates a fold catastrophe near the upper energy state of the system, while the unstable whistler waves from the cross-field current instability correspond to positive energy oscillations near a new lower-energy state of the current sheet achieved in the catastrophe. The nonlinear saturation of these whistler waves provides the irreversible relaxation of this state.

1. Introduction

The substorm concept [Akasofu, 1964] reflects the nonlinearity of the magnetosphere in extracting energy from the solar wind for immediate and accumulative releases. Accumulative energy release in substorm expansion is explosive in nature and is associated with abrupt reconfiguration of the magnetosphere and its main energy reservoir, the magnetotail, which extends over hundreds of Earth radii downstream.

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Until recently, the near-Earth neutral line (NENL) model [e.g., Hones, 1979] has been the dominant paradigm for substorms (for the review of substorm models see, e.g., Lui [1992a]), despite the enormous difficulty in demonstrating the viability of the underlying spontaneous reconnection process. Theoretical investigation persistently indicates that the tearing mode, proposed as the main reconnection mechanism in the tail current sheet, is linearly stable because of the presence of a significant magnetic field component normal to the tail current sheet.

Tearing instability for spontaneous reconnection in the tail was first proposed by Coppi et al. [1966]. It was based on a very simplified tail model without a magnetic field component normal to the current sheet and could be viewed as arising from coupling between a negative energy wave and a dissipative process. Landau resonance of electrons inside the zero magnetic field

region was proposed to provide the appropriate dissipation. Subsequently, Schindler [1974] showed that nonzero magnetic field component B_n normal to the current sheet magnetizes the electrons and restricts them from providing the necessary dissipation. As a result the required dissipation relies upon the ions that are still unmagnetized. However, Galeev and Zelenyi [1976] found that when electrons are magnetized, the basic characteristic of tearing perturbation is altered from a negative energy wave to a positive one, thus making ion dissipation invalid as a driver for the instability. The small domain of sheet parameters in which the tearing mode might be unstable, according to Galeev and Zelenyi, is shown to be unattainable on the basis of the work of Lembege and Pellat [1982]. They have demonstrated that in the case of adiabatic electrons the tearing mode might grow only for very large wavelengths, too large to be relevant. The suggestion by Coroniti [1980] that the tearing instability would be restored because of pitch angle scattering of electrons from turbulence arising from some microinstability is later considered to be invalid [Pellat et al., 1991].

In the last few years, new hopes for restoring tearing instability emerged in connection with the deterministic chaos phenomenon, i.e., stochastic motions of charged particles in the current sheet with significant field line curvature providing electron pitch angle diffusion without any field turbulence. However, Pellat et al. [1991] showed that the stochastic version of tearing instability [Büchner and Zelenyi, 1987] repeats essentially the same error as that in Coroniti's calculation. The attempt of Kuznetsova and Zelenyi [1991] to revise the conclusion by Pellat et al. [1991] was proved to be incorrect by Brittnacher et al. [1994].

The linear stability of the tearing mode, together with the extremely rare occurrence of southward magnetic field at the neutral sheet during substorms, undermined the basis of the NENL paradigm. Moreover, recent observation [e.g., Baumjohann et al., 1989] now confirmed the earlier result [Lui et al., 1977] on the dominance of earthward hot plasma flows during substorm expansion within $\sim 20\ R_E$ in the tail, implying that the near-Earth neutral line is rarely if ever formed during the initial phase of a substorm [Lui, 1992b]. These conflicts have forced a revision of the traditional NENL model to gain a better agreement with observations [e.g., Baumjohann, 1992].

Along with this discord came the finding at the inner magnetotail of another key substorm element, namely, the phenomenon of current disruption (CD) initiated by a sudden local amplification of current density in the near-Earth edge of the tail current sheet [Lui et al., 1992, and references therein]. CD manifests itself as a sudden fluctuation of the magnetic field with amplitudes comparable to or even greater than the equilibrium values and is associated with a rapid relaxation of a tail-like magnetic field toward a dipole-like field. Such a process leads to earthward plasma injection and a disturbance front expanding tailward, as suggested by Kropotkin [1972] and Chao et al. [1977]. The CD mechanism is conceived to be triggered by the local increase of the ion bulk velocity along the dawn-dusk direction [Lui et al., 1991]. Excited waves are essentially whistlers in a broadband frequency range from near the ion gyrofrequency to near the lower hybrid frequency and have mixed polarization and oblique propagation. The corresponding instability (referred to below as cross-field current instability, or CCI) is basically related to two simplified instabilities: modified two-stream instability and ion Weibel instability. Because of near-Earth plasma sheet thinning prior to substorm onset the diamagnetic current associated with a density gradient is increased and the lower hybrid drift instability may also be excited. Near the neutral sheet the increase of the plasma beta tends to reduce its growth [Huba and Papadopoulos, 1978], but the instability is shown not to be totally suppressed on the basis of a more realistic description of the tail just prior to CD [Yoon et al., 1994]. The process of nonlinear saturation of CCI leads to the deceleration and heating of the primary beam, as described in the case of quasi-linear relaxation of ion Weibel instability by Lui et al. [1993].

In many aspects, CD appears as a separate independent phenomenon induced by intense current buildup that leads to plasma turbulence and anomalous resistivity. It is then possible to construct a substorm model with CD as the primary substorm initiation process. The near-Earth neutral line is perceived as a subsequent development of CD, possibly after the tailward propagation of a rarefaction wave from the CD region [Lui, 1991]. Naturally, this newly proposed model requires further development and evaluation. One of the important aspects of the CCI-CD model is the mechanism of the formation of a strong current density prior to the CCI onset. This is presently under investigation [e.g., Ganguli et al., 1995]. The current should be built up in a timescale no less than the corresponding growth timescale ($\sim$1 min) of CCI. This sudden current buildup has indeed been reported for the late growth phase [Ohtani et al., 1992]. Another problem to be elaborated relates to the global consequences from localized CD activities and in particular the mechanism for merging of the main B_x component of the magnetic field. Some ideas, such as resistive tearing instability with CD as dissipation [Lui et al., 1991; Hesse and Birn, 1993], are proposed to reconcile the NENL model with the new CD observations.

In this paper we explore a new direction in combining the NENL model with the CD model, an approach that may resolve some of their problems mentioned above. It is based on the revision of the statement of the stability problem to take explicitly into account the triggering nature of substorm instabilities. Our reasoning in short is as follows. The substorm onset is known to be

preceded by the quasi-stationary evolution of a linearly stable current sheet in the tail. However, any quasi-stationary transition from linear stability to instability includes the stage where the linear stability/instability timescales and the timescale of the quasi-stationary evolution are comparable (the linear growth rate passes through zero). At this stage the completely nonlinear treatment of the stability problem becomes unavoidable. As is shown in section 4 on the basis of the theory of catastrophes, one can distinguish between the explicitly nonlinear transition to instability near the point of inflection of the potential energy curve of the system (the catastrophe itself) and the excitation of the system by a finite amplitude perturbation before the catastrophe. We explore in our paper the model of tearing instability in the latter situation. Previous nonlinear models of such instabilities [Galeev, 1984, and references therein] needed the formation of the neutral lines to operate. However, such a condition seems too restrictive in view of observations in the dominance of the northward component of the magnetic field and the earthward hot plasma flow during substorms that was mentioned above. Thus we deal with the problem of the instability of a finite amplitude initial tearing perturbation, which is, however, sufficiently small to conserve the equilibrium topology of the magnetic field.

The dynamical chaos is used as the particular catalyst of the instability. Previous attempts to use this phenomenon, which conserves both the particle energy and the dawn-dusk component of the canonical momentum, are thought to have failed because of the linear treatment. To transform this special sort of chaos into ordinary energy dissipation and the effective particle losses from the system, we utilize preferential stochastic acceleration proposed previously. Finally, we elaborate on the relation of the instability model to the CCI-CD phenomena. The validity of the description of substorm events in the tail as a catastrophe in a low-dimensional system with coupling of the nonlinear tearing instability and CCI are discussed.

2. Preferential Stochastic Acceleration

We search for a mechanism of electron dissipation that would be (1) nonlinear in tearing amplitude, but (2) working even at very small amplitudes when the neutral lines are not yet created and the electrons are still magnetized. Let us take into account in particular that (3) in sufficiently thin current sheets the single-particle dynamics may become unstable. In the case of a parabolic magnetic field model $\mathbf{B} = (B_0(z/L), 0, B_n)$, which is most convenient for the study of single-particle dynamics, the particle orbit becomes stochastic when the parameter $\kappa = (B_n/B_0)(L/\rho_0)^{1/2}$ (ρ_0 is the particle Larmor radius in the field B_0) is less than or of the order of unity [e.g., Büchner and Zelenyi, 1989]. It is known also that (4) because of the nonzero equilibrium field B_n the tearing perturbations have an electrostatic component [Galeev and Zelenyi, 1976; Coroniti, 1980]. The mechanism of dissipation, which involves all the features (1)—(4), has been proposed recently by Sitnov [1991] and then applied to tearing studies by Malova and Sitnov [1993]. In this study we pursue the investigation of the mechanism to take into account more realistic pitch angle distributions and the features of tearing eigenmodes.

A simple physical picture of this preferential stochastic acceleration (PSA) is as follows. Let us consider a single particle at a large pitch angle trapped within a current sheet with sharply curved magnetic field lines. As the particle moves along the field line and approaches the sharp curvature region, where its gyroradius is comparable to the curvature radius, the particle orbit becomes stochastic, and its pitch angle can be described in terms of random scattering. This scattering tends to diminish the pitch angle of the particle and allows its mirror point to be irreversibly displaced in the x direction toward Earth (Figure 1a). A weak electric field E_x associated with a tearing perturbation can then modify the energy of the particle. For deterministic

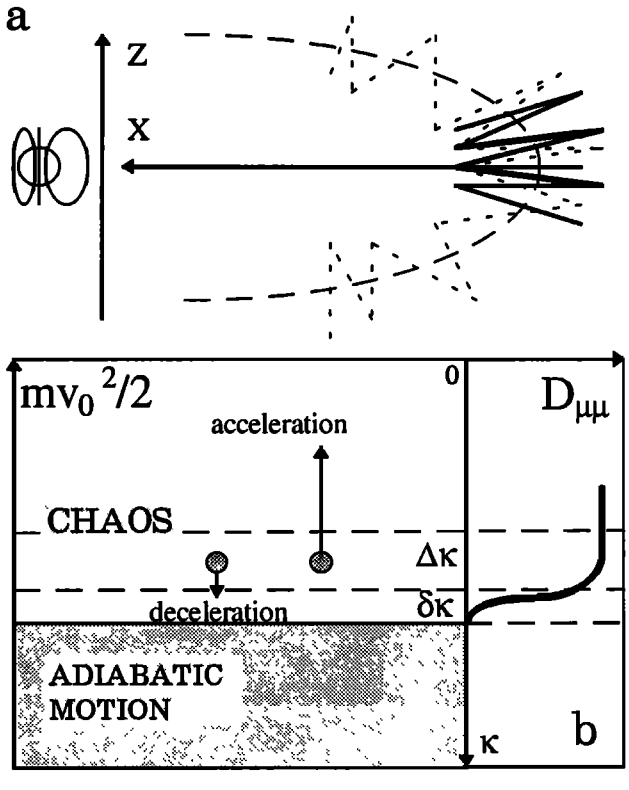

Figure 1. Qualitative picture of the effect of acceleration domination: (a) irreversible earthward particle displacement due to stochastic scatterring (solid curve corresponds to initial form of bounce oscillations); (b) particle displacement near chaos/adiabaticity boundary as it looks on the scale of the parameter κ at different phases of an additionally imposed electrostatic wave perturbation, along with the model diffusion coefficient (right panel).

chaos and linear approximation in perturbation amplitude, this energy change does not yet provide an energy exchange between the particle and the wave, because the energy change becomes zero after averaging either over the equilibrium pitch angle distribution or over the wavelength of the perturbed field. The situation is drastically different for dynamical chaos which acts in the neighborhood of a relatively clear-cut lower-energy edge of the chaotic domain. By noting that particle energy change can occur only in the chaotic regime, accelerated particles going from the adiabatic to the chaotic orbit regime can be accelerated further, whereas decelerated particles going from the chaotic to the adiabatic orbit regime will not be decelerated (Figure 1b). The above description applies to particles in which the average diffusive displacement in their energy exceeds the energy width of chaos/adiabatic boundary. This asymmetry in particle energy change is the essense of PSA. It is an irreversible process and can thus provide dissipation to tearing mode perturbation.

The transition of orbit characteristics is sharp. In the case of a parabolic field and not very small pitch angles, numerical modeling yields the boundary between chaos and adiabaticity in terms of the parameter $\kappa \sim \left(mv_0^2\right)^{-1/4}$: $\kappa = \kappa_1 \approx 1.6$, with the boundary thickness $\delta\kappa < 0.1$ [Sitnov, 1991; Malova and Sitnov, 1993] (since at small pitch angles the boundary value κ_1 as a function of the pitch angle grows logarithmically from $\approx$1.6 to $\sim$3 [e.g., Büchner and Zelenyi, 1989], the width $\delta\kappa$ has a sense only as some effective parameter). At high energies with $\kappa \ll 1$ the restoration of adiabatic motion proceeds smoothly, with the effective coefficient of diffusion over the adiabatic invariant reduced according to power law in energy, while at $\kappa \simeq \kappa_1$ the corresponding dependence is exponential [Büchner and Zelenyi, 1989; Malova and Sitnov, 1989].

The results of numerical modeling of the PSA effect are shown in Figure 2. The Hamiltonian of a particle in a parabolic magnetic field and electrostatic perturbation field is

$$H=\frac{1}{2}\left(\frac{dx}{dt}\right)^2+\frac{1}{2}\left(\frac{dz}{dt}\right)^2+\frac{1}{2}\left(\frac{z^2}{2}-\kappa x\right)^2-\alpha x\phi(z) \tag{1}$$

with $\alpha = eE_{1x}(\rho_0 L)^{1/2}/\left(mv_0^2\right)$ and dimensionless variables $\mathbf{r} \rightarrow \mathbf{r}/(\rho_0 L)^{1/2}$ and $t \rightarrow tv_0/(\rho_0 L)^{1/2}$, where v_0 is the modulus of the initial velocity. The electric field is modulated along the z direction due to the factor $\phi(z) = \left(1+z^2\right)^{-1/2}$ to provide the scaling law $\phi \propto z^{-1}$ at $z \gg 1$. This modulation allows one to exclude the additional amplification of the effect under study, which is due to the electric field induced loss cone and works at not very small α. The induced loss cone may break down the self-consistency of the theory, because the electric field E_{1x} itself is generated by the $E_{1y} \times B_n$ drift of trapped particles. In case of more realistic tearing-like perturbation this z modulation is provided in the simulation due to the finite wavelength of the perturbation and sufficiently small values of the parameter α.

The energy gain of a typical single particle at different phases (signs) of the electric field is shown in Figure 2a, with the parameter $x\phi$ being averaged over the particle gyroperiod. The parameter α is chosen in the form $\alpha = \alpha_0\eta$, with $\eta = 1, 0, -1$. But the most vivid demonstration of the effect is obtained after statistical averaging over an ensemble of particles with different initial values of the parameters κ, $\mathbf{r}$, and $\mathbf{v}$ (Figure 2b). The particles are launched from the plane $z = 0$. In this case their initial positions may be uniquely determined by the following dimensionless parameters: κ, pitch angle θ, and azimuthal angle φ; $dx/dt = \sin\theta\cos\varphi$, $dz/dt = \cos\theta$. The corresponding array of initial conditions was as follows: $\kappa = 1.3 + 0.025\,(NK-1)$ with $NK = 1, ..., 16$; $\theta = (\pi/18)\,(NT-0.5)$ with $NT = 1, ..., 9$; $\varphi = (\pi/8)\,(NF-1)$ with $NF = 1, ..., 16$. The dependence of all dimensionless parameters on the particle initial energy is taken into account so that the particular value of the parameter $\alpha = 0.075$ in Figure 2 corresponds to the midpoint of the energy interval $\kappa = 1.4875$. The initial pitch angle distribution is $f(\theta) = \sin^2\theta$, in accordance with the observation of pancake electron distributions just prior to current disruptions [Lui et al., 1992]. Figure 2b clearly demonstrates the asymmetry between the absolute values of the kinetic energy variations of accelerated and decelerated particles.

To understand better the physical reason for the PSA effect, it is useful to compare it with the acceleration of particles on Speiser orbits [Speiser, 1970]. In the latter case the increase of particle kinetic energy is only due to the earthward motion of the field lines, like the effective energy gain after an elastic collision against an approaching wall. The PSA effect can be envisioned as another "wall" in particle energy, which stops the deceleration of particles during the tailward motion of field lines arising from the perturbation field. Therefore the PSA effect should not depend critically on the direction (earthward/tailward) of the irreversible particle displacement due to the dynamical chaos in the strongly curved equilibrium magnetic field, and hence upon the plasma anisotropy. Nevertheless, the tailward scattering seems to contribute less to the effect, because it is provided by relatively weak single-particle diffusion. These features are confirmed by the results of simulation of the PSA for butterfly pitch angle distribution $f(\theta) = \cos\theta\sin\theta$, with the mirrors being placed at $z = \pm 30$ to provide zero net fluxes of particles outside the sheet (Figure 2c).

To elucidate how the electromagnetic component of tearing perturbation may influence the effect decribed above, let us consider a more realistic Hamiltonian with perturbed potentials $\tilde{\phi}$ and $\tilde{A}$ being the solutions of the appropriate eigenmode problem for quasi-neutral sheet plasma [Pritchett et al., 1991],

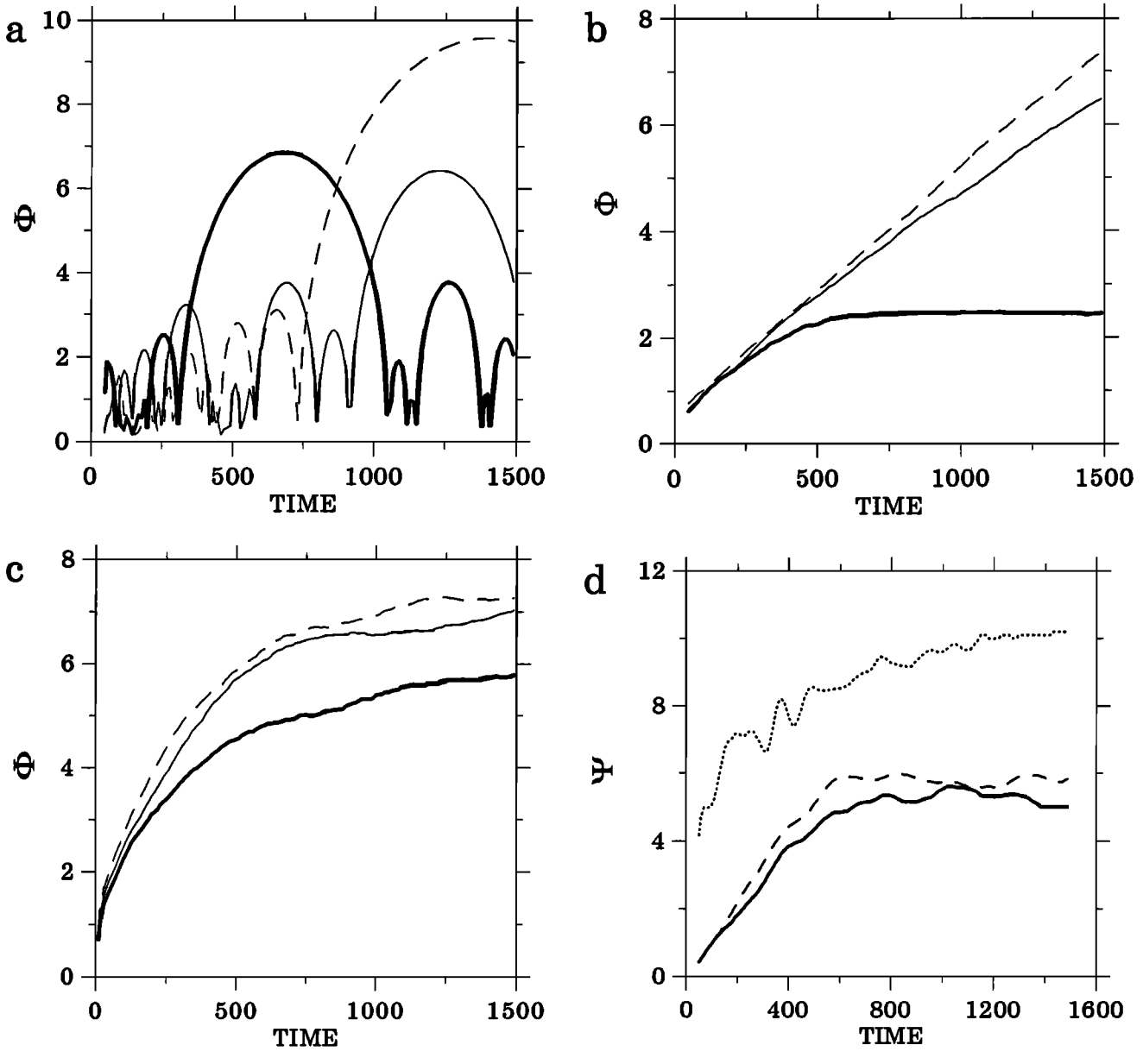

Figure 2. Simulation of the effect PSA: temporal variation of the particle kinetic energy divided by the parameter α (to compare the perturbed dynamics with the unperturbed one) near the chaos/adiabaticity boundary for different values of electrostatic field: (a) typical single-particle motion in terms of the parameter $\Phi = x\phi(z) = x/(1+z^2)^{1/2}$ with $\alpha_0 = 0.075$, $\eta = -1$, 0, 1 denoted respectively by bold solid, light solid, and dashed curves; (b) the same dependence averaged over an ensemble of 2304 particles with pancake pitch angle distribution; (c) averaged motion for butterfly pitch angle distribution with mirrors in PSBL region; and (d) net energy gain in the case of tearinglike perturbation in terms of the parameter $\psi = \tilde{\phi}(z)\sin\left(\tilde{k}x\right)/\tilde{k}$ averaged over the wave phase φ_w ($\tilde{k}x \to \tilde{k}x + \varphi_w$) (total number of particles is 18432) with $\tilde{k} = 0.02$, $\tilde{L} = 10$, and $\alpha = 0.075$, $\beta = 1$ (dotted curve); $\alpha = 0.0075$, $\beta = 1$ (dashed curve); and $\alpha = 0.0075$, $\beta = 3$ (solid curve).

$$H = \frac{1}{2}\left(\frac{dx}{dt}\right)^2 + \frac{1}{2}\left(\frac{dz}{dt}\right)^2 - \alpha x\Psi(x,z) + \frac{1}{2}\left[\frac{z^2}{2} - \kappa x - \alpha\beta\kappa\left(A(x,z) - A(x(0),z(0))\right)\right]^2 \quad (2)$$

where $\beta = mv_0^2/\left(T_i\tilde{k}\right)$, $A = \tilde{\phi}(z)\cos\left(\tilde{k}x\right)/\tilde{k}$, $\Psi = \tilde{\phi}(z)\sin\left(\tilde{k}x\right)/\tilde{k}$ with the z profile

$$\tilde{\phi}(z) = \left[1 + \tanh\left(|z|/\tilde{L}\right)/kL\right]\exp\left(-\tilde{k}|z|\right) \quad (3)$$

for $\rho_{0i} \ll L$ and

$$\widetilde{\phi}(z) = \left[\cosh\left(z/\widetilde{L}\right)\right]^{-kL} \tag{4}$$

for $\rho_{0i} \sim L$. Here $\widetilde{k} = k(\rho_0 L)^{1/2}$, $\widetilde{L} = (L/\rho_0)^{1/2}$, ρ_{0i} is the ion gyroradius in the field B_0. The parameter β is chosen to comply with the quasi-neutrality condition and the relation

$$\langle \widetilde{n}_1 \rangle = ikn_0 \langle A_{1y} \rangle / B_n \tag{5}$$

which connects, according to Coroniti [1980], Galeev [1984], and Pritchett at al. [1991], the bounce-averaged (terms in angle brackets) vector potential with the bounce-averaged part of the electron density perturbation, $\langle \widetilde{n}_1 \rangle = \int \left\langle \widetilde{f}_{1e} \right\rangle d^3v$. Here $\widetilde{f}_{1e} = f_{1e} - A_{1y}(\partial f_{0e}/\partial A_{0y})$, where f_{0e} and f_{1e} are the equilibrium particle distribution and its perturbation, respectively, for electrons.

As follows from (2)—(4), a neutral line appears when $\alpha\beta = 1$ or $B_{1z} = B_n$. The influence of the electromagnetic component of the tearing mode on the effect of preferential acceleration is determined by the corresponding electromagnetic perturbation of the parameter κ: $\delta\kappa_{em} \sim \alpha\beta\kappa_1$. Hence to neglect the electromagnetic influence, it is sufficient that $\delta\kappa_{em}$ be no more than the natural nonlinear width of the chaos/adiabaticity boundary $\delta\kappa$. In particular, for electrons at $B_n/B_0 \sim 0.2$, $mv_0^2 \sim 0.2T_i$, and $kL \sim 0.3$ it is sufficient that the parameter α be no more than $0.3 \times \delta\kappa/\kappa_1^2$. In reality, the influence of the electromagnetic component appears to be even weaker, because the phases of the fields E_{1x} and B_{1z} are shifted by $\pi/2$ in relation to each other. Therefore within a half period of fixed sign of the field E_{1x} the correction term $\delta\kappa_{em}$ alters its sign, so that its effect is reduced further by the additional averaging over the parameter κ.

The simulation result of the PSA effect for the Hamiltonian (2) with the z profile (4) is shown in Figure 2d. An additional averaging over the wave phase φ_w ($\widetilde{k}x \to \widetilde{k}x + \varphi_w$) is performed; $\varphi_w = (\pi/4)(NP - 0.5)$ with $NP = 1, ..., 8$. No additional z modulation is assumed. The real tearing mode provides the corresponding effective modulation due to its finite wavelength. Thus we have demonstrated that preferential stochastic acceleration is present for different pitch angle distributions and perturbation wave polarizations and that it can occur at tearing amplitudes less than the amplitude needed to form a neutral line.

An important question is whether the real parameters of the Earth's magnetotail are appropriate to reach the lower energy edge of the electron chaos region where the PSA works. Pulkkinen et al. [1994] for the CDAW 6 substorm (March 1979) found that κ_e could be about unity (their Figure 4). Lui et al. [1992] on the basis of in situ measurements from 15 CDs detected by CCE found $\kappa_e \geq 4$. However, with κ_e calculated from the actual ratio between the field line curvature and gyroradius rather than the crude estimate of $(B_n/B_0)(L/\rho_0)^{1/2}$, Lui [1993] showed that $\kappa = 1.7$ in one of the CD events considered by Lui et al. [1992], which is surprisingly close to κ near the chaos/adiabaticity boundary for particles with not too small pitch angles.

3. Nonlinear Tearing Instability

The inclusion of the dynamical PSA effect in the kinetic treatment of the tearing instability is based on the analogy between the PSA and the inverse Cerenkov effect providing ordinary linear Landau damping of electrostatic waves in isotropic plasma. The only difference is that in the latter case the resulting energy gain is determined by a larger number of lower-energy particles in equilibrium distribution near the Landau resonance, i.e., by the derivative of the energy distribution, while near the chaos/adiabaticity boundary it is determined by energy asymmetry between the processes of acceleration and deceleration, i.e., by the energy distribution itself.

The effective width in κ of the layer with preferential acceleration for the tearing-like perturbation (14)—(17) at $\alpha\beta \ll 1$ is given by the expression [Sitnov, 1991; Malova and Sitnov, 1993]

$$\Delta\kappa \simeq -\kappa_1 \frac{D_0^2 \omega_B}{mv_0} \int_0^t eE_{1x}(\tau)\, d\tau \tag{6}$$

where D_0 is a dimensionless factor in the following approximation of the diffusion coefficient over the magnetic moment,

$$D_{\mu\mu} \simeq 4D_0 \left(\mu B_n / \left(mv_0^2\right)\right)^3 [1 - \mathrm{sgn}(\kappa - \kappa_1)] \tag{7}$$

and $\omega_B = v_0/(\rho_0 L)^{1/2}$. Under the condition $\Delta\kappa \geq \delta\kappa$ the effect becomes formally linear, and in particular, the resulting energy gain within the layer $\Delta\kappa$ becomes proportional to E_{1x}^2, as occurs in the case of linear Landau damping. This result means that this essentially nonlinear effect may be included in a formal scheme of linear stability analysis of the tearing mode.

The corresponding equation for perturbed electron density $\langle \widetilde{n}_1 \rangle$ averaged over the bounce period may be derived from the linearized quasi-hydrodynamic equations for the two-component electron fluid. The first component describes the main body of relatively cold ($T_e \ll T_i$) electrons (both adiabatic and chaotic) complying with the drift approximation [Galeev, 1984; Wang and Bhattacharjee, 1993],

$$\begin{aligned} -e\langle E_{1y}\rangle + e\langle v_{1x}\rangle B_n/c &= m\langle dv_y/dt\rangle \\ +e\langle v_{1x}B_x(z)\rangle/c &\simeq 0 \end{aligned} \tag{8}$$

$$-ikT_i\langle \widetilde{n}_1\rangle + \left\langle \widetilde{j}_{1y} \right\rangle B_n/c \simeq 0 \tag{9}$$

(equation (9) also involves ion parameters due to quasi-neutrality), while the second one describes a relatively small amount of fast accelerating electrons from the neighborhood of the chaos/adiabatic boundary,

$$\left\langle n_1^{(a)} \right\rangle \simeq n_0 f_0(\kappa_1) \langle \Delta\kappa \rangle \tag{10}$$

$$v_x^{(a)} \simeq 2 v_0 D_0^2 \omega_B |t| \tag{11}$$

These parameters are similar to resonant particles in the case of Landau damping.

The connection between these two dynamically different components of electrons is provided by the electron density conservation equation and the quasi-neutrality condition,

$$\frac{\partial \langle \tilde{n}_1 \rangle}{\partial t} + \frac{\partial}{\partial x}\left(n_0 \langle v_{1x} \rangle + \left\langle n_1^{(a)} \right\rangle v_x^{(a)} \right) = 0 \tag{12}$$

$$e \langle E_{1x} \rangle n_0 = ikT_i \langle \tilde{n}_1 \rangle \tag{13}$$

The resulting equation takes the form [Malova and Sitnov, 1993]

$$\frac{\partial \langle \tilde{n}_1 \rangle}{\partial t} + \frac{\kappa_1 f_0(\kappa_1) D_0^4 \omega_B^2 k^2 T_i}{m} \int_0^t |t| \tau \langle \tilde{n}_1 \rangle (\tau)\, d\tau = i \frac{k n_0}{B_n} \frac{\partial \langle A_{1y} \rangle}{\partial t} \tag{14}$$

Here $f_0(\kappa)$ is the initial electron distribution over the parameter κ. The right-hand side of (14) describes the above mentioned electron compressibility effect (compare equations (14) and (5)) due to the drift motion of the main group of electrons in the fields E_{1y} and B_n. The dissipative properties of the second term on the left-hand side of (14) result from the dependence of the resonant particle density upon the sign of the energy variation (6).

According to (14), under the condition on the tearing growth rate

$$\gamma \le \gamma_a = D_0 \sqrt{\omega_B k} \left[\kappa_1 f_0(\kappa_1) T_i / m \right]^{1/4} \tag{15}$$

the perturbation of the plasma density is determined mainly by "resonant" near-chaos/adiabatic boundary particles and is considerably reduced in comparison with the case of pure adiabatic response (5):

$$\langle \tilde{n}_1 \rangle \simeq i \left(\gamma |\gamma|^3 / \gamma_a^4 \right) (k n_0 / B_n) \langle A_{1y} \rangle \tag{16}$$

A similar reduction was suggested previously by Galeev as a result of collisions and loss cone precipitation.

Now consider the energy balance equation of the tearing mode,

$$\frac{1}{8\pi} \frac{\partial}{\partial t} \int d\mathbf{r} \left(|\mathbf{B}_1|^2 + |\mathbf{E}_1|^2 \right) = - \int d\mathbf{r} \mathbf{j}_1 \mathbf{E}_1^* \tag{17}$$

Its right-hand side can be represented as a sum of the following three terms [e.g., Galeev, 1984]:

Current filamentation power

$$\frac{\partial}{\partial t} \int d\mathbf{r} \int d\mathbf{v} |A_{1y}|^2 \frac{q_\alpha v_y}{2c} \frac{\partial f_{0\alpha}}{\partial A_{0y}} = \frac{1}{8\pi} \frac{\partial}{\partial t} \int d\mathbf{r} \frac{2 |A_{1y}|^2}{L^2 \cosh^2(z/L)} \tag{18}$$

The corresponding perturbation of the distribution function and the current density is connected with the magnetic part of the perturbed field in the linearized part of the Vlasov equation.

Drift motion power

$$- \int d\mathbf{r} \tilde{j}_{1y} E_{1y}^* = ik \frac{T_e + T_i}{B_n} \int d\mathbf{r} \langle \langle \tilde{n}_1 \rangle \rangle \frac{\partial \langle A_{1y}^* \rangle}{\partial t} \tag{19}$$

Here the current density is generated by gradient drift in the field B_n (see (9)). The additional angle brackets are introduced to take into account that the coherent density perturbation in response to the tearing electric field E_{1y} of the type of (5) or (16) will be valid unless the particle displacement along the Earth-Sun direction exceeds half the tearing wavelength [Galeev, 1984].

Ion Landau dissipation power

$$- \int d\mathbf{r} j_{1y}^{(L)} E_{1y}^* = - \frac{\gamma^2 \omega_{pi}^2}{8\sqrt{\pi} c^2 |k| v_{Ti}} \int_{-d_i}^{d_i} dz \int dx dy |A_{1y}|^2 \tag{20}$$

Here $d_i = (\rho_{0i} L)^{1/2}$, and ω_{pi} is the plasma frequency for ions. The right-hand side of (20) has been estimated in the thick sheet approximation ($\rho_{0i} \ll L$). Note that both the second term in (14) and expression (20) describe the special groups of electrons and ions, respectively, and thus they reflect the kinetic nature of the process in the form of some effective four-fluid hydrodynamics.

When the stochastic electron acceleration is negligible, the term (19) with $\langle \tilde{n}_1 \rangle$ given by the right-hand side of (14) will overbalance the negative energy contribution from (18) and make the energy of the tearing mode positive. Therefore the dissipation (20) results in the damping of the mode; i.e., (17) has only solutions with $\gamma < 0$.

When the electron stochastic acceleration dominates the motion, the term (19) due to (16) takes the form

$$- \frac{(T_e + T_i) k^2 n_0}{2 B_n^2} \frac{\gamma^2 |\gamma|^3}{\gamma_a^4} \int d\mathbf{r} |\langle A_{1y} \rangle|^2 \tag{21}$$

i.e., term (19) not only is decreased in comparison with the case of adiabatic electrons due to small-parameter γ/γ_a, but becomes dissipative like term (20). Therefore the energy of the tearing mode becomes negative, and the dissipative processes (20) and (21) drive the instability. Its growth rate may be inferred from the energy principle by using expressions (3) and (4) as trial func-

tions for the perturbed potential A_{1y}. In particular, in the case of a thick sheet ($\rho_{0i} \ll L$) and relatively strong electron dissipation, the growth rate can be expressed due to (21) in the form

$$\gamma = \gamma_e = \gamma_a \left(4\frac{1-k^2L^2}{kL+\Delta}\right)^{1/4} \left(\frac{B_n}{B_0}\right)^{3/4} \tag{22}$$

where $\Delta = (2B_n/(kLB_0))^{1/2}$. If this growth rate is sufficiently strong (γ due to dissipative term (21) is more than γ due to (20)), it will be saturated at the level of the so-called ion tearing instability proposed by Schindler [1974],

$$\gamma \to \gamma_i^{(1)} = \Omega_{0i} \left(\frac{\rho_{0i}}{L}\right)^{5/2} \frac{1-k^2L^2}{kL+\Delta} \tag{23}$$

where Ω_{0i} is the ion gyrofrequency in the field B_0.

Pritchett et al. [1991] have extended this estimate of maximum permissible growth rate in the case of extremely thin isotropic sheets with $\rho_{0i} \sim L$,

$$\gamma \le \gamma_i^{(2)} = \frac{\Omega_{0i}}{\sqrt{\pi}} \frac{kL\,(2+kL)\,(1-kL)}{1+2\,(\rho_{0i}/L)^2} \tag{24}$$

Even though at $kL \le 1$ and $\rho_{0i} \sim L$ the ratio $\gamma_i^{(2)}/\gamma_i^{(1)} \sim kL \le 1$, both these estimates seem to be quite compatible with the corresponding substorm timescales. Hence it is more important to clear up other limitations of the instability.

The most obvious limitation is determined by (15). In particular, for tail parameters $kL \sim 0.3$, $L \sim 6000$ km, $mv_0^2/2 \sim T_e \sim 2$ keV, $T_i \sim 10$ keV, $D_0 \sim 0.01$ [see Malova and Sitnov, 1993], and $f_0(\kappa_1) \sim 0.1$ one finds that $\gamma \le \gamma_a \sim 1\ \mathrm{s}^{-1}$. As a result the growth rate of the instability (22) can actually reach the level of ion tearing instability described by (23) and (24).

Two other limitations are connected with the ranges of applicability of the linear approximation (6). They determine the threshold of the instability in wave amplitude. First, the minimum length of acceleration corresponding to the minimum shift in κ ($\Delta\kappa \sim \delta\kappa$) may not exceed half the wavelength,

$$(\delta\kappa/\kappa_1)\,(mv_0^2/(eE_{1x})) \le 2\pi/k \tag{25}$$

From comparison of (25) with the condition of neutral line formation $\alpha\beta = 1$ one can infer the inequality $T_i \ge mv_0^2\,(\delta\kappa/\kappa_1)$, which implies a considerable range of permissible pre-neutral line amplitudes. As a result the instability, having been restored by the mainly earthward stochastic electron acceleration, also provides the earthward flows of heated ions with $\kappa_i \sim 0.1 \times \delta\kappa$, which are forced to move along Speiser-like transient orbits near considerably disturbed but not completely topologically altered magnetic field lines. Hence, having been heated in the central sheet region due to Landau resonance, those ions are ejected outside in the earthward direction.

Second, the minimum acceleration time (to gain $\Delta\kappa \sim \delta\kappa$) should be less than the instability growth time, $\gamma_t^{-1} \le \gamma^{-1}$ with

$$\gamma_t = D_0\sqrt{\kappa_1 eE_{1x}\omega_B}/\sqrt{2\delta\kappa mv_0} \tag{26}$$

At $B_0 \sim 30$ nT (other parameters have been given above) and $\gamma \sim \gamma_t \sim 1$ s the smallest possible amplitude inferred from (26) is $E_{1x} \sim 5 \times 10^{-8}L$ [km]mV/m, well within the range of typical fluctuations [e.g., Pedersen et al., 1984].

4. Discussion

One of the key features of the tearing instability model described above is its nonlinearity. It is the nonlinearity that provides the destabilization of the tearing mode, shown by Pellat et al. [1991] to be linearly stable because of the conservation of canonical momentum P_y. The latter condition holds throughout our model. In particular, it is explicitly taken into account in the Hamiltonian (2) and in the linear part of equation (14), which is equivalent to equation (14) of Pellat et al. [1991] inferred directly from this condition. In terms of linear theory the nonlinear part of (14) associated with PSA resembles the loss of particles from the system. Under the condition of conservation of the total number of electrons it becomes possible because of dynamical separation of electrons in the phase space into two populations: one drifting in the crossed fields E_{1y} and B_n in accordance with linear theory and the other preferentially accelerated because of PSA.

The essential features of our model of tearing instability, its possible relation to other key elements of the substorm, and unavoidable nonlinearity of the instability as a constituent of the substorm trigger may be considered in more detail by the theory of catastrophes [e.g., Haken, 1978]. This theory is widely accepted as an appropriate mathematical tool to describe abrupt changes in a low-dimensional system driven by quasi-stationary evolution of a set of control parameters. This theory can be applied here if we treat the tearing instability as a process for the growth of a large-scale and one-mode perturbation. The state of the system may be described by a single state-parameter such as the dawn-dusk inductive electric field E_{1y}. The proper control parameter should vary continiously in the neighborhood of the catastrophe and can be inferred from the linear theory of stability. For instance, in the case of the isotropic sheet, according to Coroniti [1980], Lembege and Pellat [1982], and Galeev [1984], the damping rate of the tearing mode is controlled by the parameter $b_{eff} = (\pi kLB_0/(4B_n))$ as $\gamma \sim \gamma_0\,(b_{eff})\,(1-b_{eff})$ with $\gamma_0(1) \ne 0$. On the other hand, according to Schindler and Birn [1982], the evolution of this parameter obeys the law $b_{eff} \propto \left(B_o^2/(8\pi)\right)^{1/2-1/\nu}$, where ν is a polytropic index equal to 1 for isothermal and 5/3 for isen-

tropic processes. Therefore b_{eff} tends to be diminishing during the growth phase. It is an example of coherent evolution of the tail toward the catastrophe (driven by the single parameter B_0). Then the transition to instability appears as a minimum dimension catastrophe or the so-called fold catastrophe with one state and one control parameter.

On the plot of the potential energy of the system as a function of the state parameter (Figure 3) the fold catastrophe resembles a process of quasi-stationary transformation of the potential minimum into the point of inflection. Being located near the bottom of the potential well before the catastrophe, the system is linearly stable because of positive energy of small deviations from the minimum. The transition to instability is possible only at the moment of the catastrophe or before the catastrophe under the influence of a finite amplitude perturbation necessary to surmount the potential barrier. In both cases the destabilization of the system proves to be nonlinear. In the first case it is because the linear eigenmodes become zero-energy modes. This explicitly nonlinear excitation can be described by the theory of soft modes [Haken, 1978]. The second case assumes a pseudo-linear treatment of the system excited by a finite amplitude perturbation before the catastrophe, with the energy of the corresponding eigenmodes being negative because of negative curvature of the potential energy plot near the potential barrier (Figure 3). Note that the inference of Lembege and Pellat [1982] and Pellat et al. [1991] about the overall linear tearing stability of the geomagnetotail current sheet argues in favor of the second case regime of the development of the instability, i.e., the regime explored above.

Let us try to extend the description of the nonlinear tearing instability in terms of the theory of catastrophes on a more complex set of substorm events. The observational reasons for such an extension are the small effective dimension of the magnetosphere, inferred from AE and AL series by the use of correlation dimension technique, and the nonlinearity of the reaction of the magnetosphere in response to solar wind parameter variations based on the results of linear and nonlinear prediction filter studies [Klimas et al., 1996, and references therein].

The extension of the theory of catastrophes on a more complex set of substorm events requires a corresponding sophistication of the catastrophe implied, and in particular, the enlargement of its dimension. The only scheme of substorm dynamics based on a more dimensional pleat catastrophe has been proposed by Lewis [1991]. The main features of this scheme are reflected by the system potential energy of the form $P(x) = x^4 + 2ax^2 + 4bx + \text{const}$, with the control parameters a and b and the state parameter x. In comparison with the case of the fold catastrophe this potential may have more (two) minima corresponding to different equilibrium states (Figure 4). The substorm is modeled by the disappearance of the upper minimum due to a local fold catastrophe. In spite of some reasonable speculations concerning the real magnetospheric analogues of the parameters a, b, and x, Lewis' scheme suffers from too abstract a description of the catastrophe. In particular, neither the local fold catastrophe nor the relaxation of the magnetosphere near its lower-energy

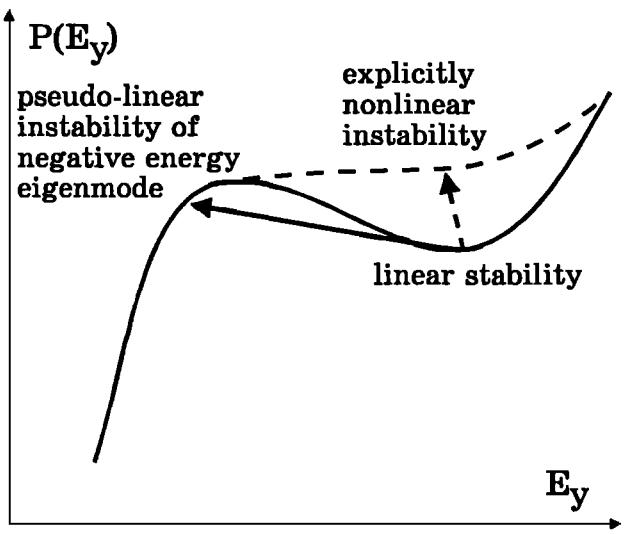

Figure 3. Effective potential energy of the current sheet near the marginal stability. The dashed arrow marks the transformation of the potential minimum into the point of inflection, that is, the fold catastrophe itself. The solid arrow denotes the transition to instability before the catastrophe due to a finite amplitude "germ" perturbation.

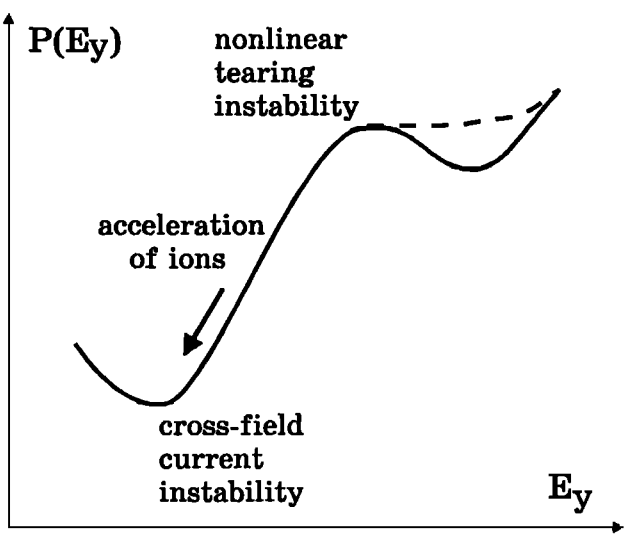

Figure 4. Effective potential of the current sheet in the vicinity of a substorm reconstructed within the framework of the pleat catastrophe model on the basis of the properties of nonlinear tearing and cross-field current instabilities.

state is coordinated with specific substorm phenomena. We suggest in the capacity of such phenomena the nonlinear tearing instability described above and the cross-field current instability [Lui et al., 1991].

Actually, as is shown above, the nonlinear tearing instability based on PSA can be presented as one of the regimes of the fold catastrophe. Therefore the corresponding energy balance equation (see (17) and (18)—(20)) can be used to approximate the whole effective potential $P(x)$ near its upper minimum and the potential barrier alongside (Figure 4). On the other hand, typical CCI whistlers are known to be positive energy excitations. Their free energy source is a linearly unstable current formed mainly by ions prepared (accelerated or unmagnetized) by some other process. The result of the growth and nonlinear saturation of CCI should be an increase of electromagnetic fluctuations and a reduction of the initial current due to ion deceleration with a relatively small increase of the plasma internal energy. Taking into account these properties of CCI, one can coordinate it with positive energy excitations of the system (tail current sheet) near the lower minimum of the potential energy (Figure 4). Then CCI will develop in the process of the acceleration of ions by the dawn-dusk inductive electric field E_{1y} [Speiser, 1970] under a suitable phase of the nonlinear tearing mode, while the relaxation of the system near its lower-energy minimum can be described by the saturation of the development of CCI due to quasi-linear effects [Lui et al., 1993].

5. Conclusion

In this work we tried to resolve a number of problems, in particular, the linear stability of the tearing mode as a mechanism of the explosive reconnection in the magnetospheric tail and a substorm trigger. It has been revealed that the quasi-stationary evolution of the parameters of the tail current sheet toward its marginal stability leads to transition from linear stability to a nonlinear development.

To overcome the linear stability, we adopt the suggestion by Malova and Sitnov [1993] that some electrons near the equilibrium chaos/adiabatic boundary undergo preferential stochastic acceleration from tearing perturbation electric fields and provide a means of dissipation to promote nonlinear tearing. A set of single-particle orbits with realistic pitch angle distributions in a given field model (equilibrium magnetic field plus tearing field) is computed and averaged over the particle ensemble and wave phases to verify the preferential acceleration of electrons. The resulting nonlinear tearing instability is shown not to modify the field topology (i.e., no neutral line formation) within a large range of tearing amplitudes. It has sufficiently small timescales, about the growth time of the hypothetic linear ion tearing instability [Schindler, 1974], well within the smallest timescales of the substorm.

Thus the nonlinear approach gives a chance for tearing instability to remain a key element of the substorm onset mechanism. Moreover, the tearing instability appears to be consistently connected with another important substorm element, namely, the cross-field current instability of the current disruption phenomenon, within the whole picture of the substorm explosion in the tail as a low-dimensional catastrophe. With the neutral lines not yet created, the nonlinear tearing instability serves as a process of current filamentation to form ion bulk flows (intense current density) for linear cross-field current instabilities rather the real tearing of the initial field lines. The quasi-linear relaxation of CCI promotes further reduction of the current in the sheet. The tearing itself with the formation of neutral lines and magnetic islands could be developed afterward, probably as a result of CCI-CD phenomena, as suggested in the synthesis substorm model of Lui [1991]. Naturally, we regard this possibility as one of many scenarios for substorms.

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