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BALLOONING INSTABILITY AS A MECHANISM OF THE NEAR-EARTH ONSET OF SUBSTORMS

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Abstract. After introducing a mathematical definition of the tail-like equilibrium and the dipole-like equilibrium in the magnetosphere, it is shown by using physical intuition based on the Energy Principle that the incompressible assumption for the ballooning instability is more valid for the tail-like configuration when the unstable ballooning mode is strongly localized near the equator. Therefore, before the substorm onset, the near-Earth plasma sheet becomes more tail-like and more likely to be subject to the ballooning instability without the stabilizing influence of the compressibility, when the critical plasma β due to the stabilizing tension force is exceeded. The onset of the ballooning instability in the near-Earth plasma sheet seems promisingly relevant to the substorm onset phenomena. Also, the effect of the stochastic plasma dynamics on the ballooning and interchange instabilities is clearly shown.

1. Introduction

The substorm is a global self-organization process in the magnetosphere, which reconfigures magnetospheric magnetic fields having a highly stretched nightside tail to a potential field (dipole field). Although the exact and quantitative details of how the field lines are dragged tailward by the solar-wind have not been completely agreed upon, the existence of the growth phase of the substorm, in which the solar-wind energy is stored in the tail of the magnetosphere in the form of the magnetic energy of the highly stretched tail field, has been tested and confirmed by intensive observations using in situ and ground-based observations. However, the specific physical mechanism of the onset of the expansion phase of the substorm remains unidentified even on a qualitative basis. Since the onset of the expansive phase is the initiation of the release of enormous energy in the magnetosphere in a short time interval, the understanding of the specific mechanism leading to the onset of the expansion phase is a key to understanding the whole substorm process in the magnetosphere. Several different mechanisms have been invoked to explain the substorm onset. Among other processes Roux et al. (1991a,b) have suggested that observed phenomena associated with the substorm onset are consistent with the ballooning instability in the plasma sheet and thus the ballooning instability may be an onset mechanism.

The ballooning instability is a pressure-driven ideal magnetohydrodynamic (MHD) instability growing in a fast MHD time scale in a high-β plasma and it occurs where the pressure gradient vector and the field line curvature vector are in the same direction. In this sense the driving mechanism of the ballooning instability is similar to the interchange instability, which has also been studied intensively in the magnetospheric context (e.g. Gold 1959). Although the ballooning instability has been intensively studied in the context of fusion plasma confinement, the application of the ballooning instability to the magnetospheric context is only recent. Viñas and Madden

Space Science Reviews **95**: 387-398, 2001

(1986) applied the ballooning instability to the plasmapause and investigated the effects of the azimuthal shear flow on the ballooning instability. Miura et al. (1989) showed by the numerical eigenmode analysis that the plasma sheet is subject to the ballooning instability and Ohtani et al. (1989b) showed by the numerical eigenmode analysis that the outer edge of the ring current is subject to the ballooning instability and/or interchange instability. Roux et al. (1991a,b) investigated an isolated dispersionless substorm with ground based observations and in situ observations by a geostationary satellite and have suggested, based on the studies of Miura et al. (1989) and Ohtani et al. (1989a,b), that the near-Earth plasma sheet is subject to the ballooning instability and they attributed the partial cancellation of the tail current, the resulting particle injection, and the development of a westward traveling surge to the emergence of the ballooning instability.

In the high-β plasma the Alfvén mode and the compressible slow mode are coupled where there is a field line curvature (Southwood and Saunders 1985). Miura et al. (1989) and Ohtani et al. (1989a) derived coupled eigenmode equations of the ballooning instability, which show the coupling of the Alfvén mode and the slow mode in the one-fluid MHD and in the two-fluid equations, respectively. In the numerical eigenmode analysis of the ballooning instability in a tail-like equilibrium Miura et al. (1989) assumed intuitively that the parallel component of the velocity perturbation vanishes and the ballooning mode in the magnetosphere is incompressible, because the ballooning instability arises in the shear Alfvén branch and the slow mode is only stabilizing. They found that the plasma sheet is subject to the ballooning instability and the fundamental symmetric mode is destabilized by the instability. However, by employing the Energy Principle (Bernstein et al. 1958, Freidberg 1987), Lee and Wolf (1992) have argued that the ballooning instability is stabilized by the compressibility in the highly stretched tail-like field configuration. The purpose of the present paper is to show intuitively, based on the Energy Principle, that the incompressible assumption for the ballooning instability as adopted by Miura et al. (1989) is more valid for the tail-like configuration, when the ballooning mode is strongly localized near the equator, and to suggest on a more solid basis the relevance of the onset of the ballooning instability in the tail-like equilibrium to the onset of the substorm expansion phase. More detailed analysis of the validity of the incompressible assumption based on mathematical formulation is deferred to another paper (Miura 1999).

2. Requirement for the Theory of the Near-Earth Substorm Onset

The importance of the near-Earth region or even the region as close as the geosynchronous region in the substorm onset has been emphasized by intensive observations. Those intensive and crucial observations of the substorm onset have shown that the location of the onset is near-Earth between 6 R_e and 11 R_e in the premidnight, where R_e is the earth's radius, it is localized azimuthally, and the time scale of the onset is very fast, less than a few minutes. The importance of the near-Earth region in the substorm onset has also been supported by evidence that the auroral arc that brightens first during a substorm maps to the inner edge of the plasma sheet. Therefore, any potential mechanism of the substorm onset must be able to explain (1) the near-Earth location of the onset, its localized nature in the local time, its fast time

Space Science Reviews **95**: 387-398, 2001

scale, and (2) the relation of the specific onset mechanism to the global reconfiguration of the magnetosphere, to the energy release of the magnetic energy stored in the highly stretched tail field lines, and also to the auroral brightening in the magnetosphere-ionosphere coupling system associated with the onset.

Whereas the ion tearing instability in the 2-D model of the thick quasineutral sheet is stabilized by the electron compressibility effect due to a normal magnetic field component (Lembege and Pellat 1982, Pellat et al. 1991), which is stronger nearer to the earth, and a small azimuthal wavenumber is favorable for the tearing instability, a large azimuthal wavenumber is favorable for the ballooning instability and the growth rate of the ballooning instability in the plasma sheet is larger nearer to the earth (see Figure 6 of Miura et al. (1989)). Therefore, the localized nature of the ballooning instability in the azimuthal direction and its preference for the near-Earth region in the plasma sheet are favorable for the substorm onset.

Since the field-lines in the near-Earth plasma sheet are stretched and become tail-like during the growth phase (Kaufmann 1987), the field line curvature radius near the equator may become comparable to the Larmor radius of bulk of ions. In such a case the ion stochastic dynamics become important in the near–Earth plasma sheet (e.g. Büchner and Zelenyi 1989, Chen 1992). Existence of such stochastic plasma has been suggested by observations in the near-Earth plasma sheet by Lui et al. (1992). Therefore, any potential mechanism of the substorm onset must be robust even in the presence of the stochastic plasma dynamics.

3. Tail-like Equilibrium and Dipole-like Equilibrium

Although there is no quantitative definition of the tail-like configuration and the dipole-like configuration, which is important in understanding the dynamical change of the magnetospheric configuration in substorm, we propose that a good criterion of the equilibrium configuration determining whether it is tail-like or dipole-like may be the smallness parameter ε, which is equal to the ratio of the perpendicular gradient-B scale length to the field-line curvature radius. When $\varepsilon \ll 1$ is satisfied on average except near the equator, the configuration is considered to be tail-like and when $\varepsilon \gg 1$ the configuration is considered to be dipole-like. Here, we should note that since the dipole field in vacuum is a zero-β field (zero pressure), the dipole field in vacuum satisfies $\varepsilon = 1$. This means that in the dipole field in vacuum the tension force by the field line curvature is balanced with the magnetic pressure gradient force.

When $\varepsilon \ll 1$ on average except near the equator, the pressure gradient force is balanced with the magnetic pressure gradient force. Such an equilibrium is possible only when the equilibrium is tail-like and the field lines are stretched substantially, so that the field line curvature radius is large except for a tiny region around the equator. Figure 1(b) shows schematically this case, where the configuration is tail-like. The extreme limit ($\varepsilon = 0$) of this case is the one-dimensional neutral sheet (Harris sheet), in which the pressure gradient force is outward away from the neutral plane and this force is balanced with the magnetic pressure gradient force directing toward the neutral plane. In the highly stretched 2-D plasma sheet, such small ε limit is satisfied except near the equator (see Figure 4 of Miura et al. (1989)). In the opposite limit $\varepsilon \gg 1$ the pressure gradient force is balanced with the magnetic tension force due to the field-line curvature.

Space Science Reviews **95**: 387-398, 2001

Therefore, this case corresponds to that, where the field line is substantially curved or dipole-like. Figure 1(a) shows schematically this case, where the field line is substantially curved and the tension force is nearly balanced with the pressure gradient force. We expect that the outer edge of the ring current is an example of this large ε equilibrium, because the field line is not strongly stretched from the dipole as in the plasma sheet. Ohtani et al. (1989b) obtained by numerical iteration a model equilibrium representing the outer edge of the ring current. It is shown in Figure 2 of Ohtani et al. (1989b) that $\varepsilon > 1$ is indeed satisfied in their model.

4. Ballooning Instability versus Interchange Instability

It is well known that the ballooning instability and the interchange instability are analogous to the gravitational Rayleigh-Taylor instability, in which the effective gravity is given by equating the gravitational drift with the combined gradient-B and curvature drifts. Whereas the interchange or flute instability extends uniformly along the entire length of the field line, the ballooning instability is localized to a finite region of unfavorable curvature. Therefore, the growth rate of the ballooning instability should be smaller than the growth rate of the interchange instability. In the interchange instability, the electrostatic potential perturbation, which describes the azimuthal electric field perturbation (see Miura et al. (1989)), is uniform along the field line and the field line tension does not prevent the instability. Therefore, only compressibility is the stabilizing factor in the interchange instability. Although the interchange instability has traditionally been investigated for a dipole-like magnetospheric field configuration (e.g. Gold 1959), the interchange instability is also possible in the tail-like equilibrium (Horton et al. 1999). On the other hand, the ballooning instability or the unstable ballooning mode in the tail-like equilibrium is strongly localized near the equator (see Figure 7 of Miura et al. (1989)). Since the ballooning mode is strongly peaked at the equator and is decaying exponentially toward the ionosphere, the electrostatic potential perturbation is a function of the distance along the field line and the field line tension adds to the compressibility as stabilizing factors.

5. Ballooning Instability in the Tail-like Equilibrium

5.1. Ideal MHD Plasma

When there is a field line curvature in the high-β plasma, the Alfvénic perturbation becomes a source of the compressional perturbation and the compressional effect propagates along the field line with the slow mode speed. Since the stability analysis studies the steady state, when the perturbation becomes steady long after the initial transient phase, the compressional effect given at the equator, for example, is averaged along the field line and the compressional factor, which is the divergence of the perturbed velocity vector, is homogeneously distributed and becomes constant along the field line. This is what one of the minimization conditions of the Energy Principle with respect to the parallel displacement (Bernstein et al. 1958, Freidberg 1987, Lee and Wolf 1992) requires. Notice that the minimization condition is derived from the minimization of the square of the absolute value of the divergence of the displacement

Space Science Reviews **95**: 387-398, 2001

vector and this guarantees that the compressional factor becomes constant along the field line. Although the ballooning driving term in the ballooning eigenmode equation (Miura 1999), which is proportional to the perturbed electrostatic potential, is a local quantity changing with the distance along the field line and is strongly peaked at the equator, the compressional stabilizing term is proportional to the flux average of the perturbed electrostatic potential, which is a global quantity and arises from the spreading of the compressional effect over the entire field line as we have seen above. When the magnetic field lines are more stretched tailward and the small ε approximation for the equilibrium along with the long-thin assumption for the perturbations are more satisfied on average except near the equator, the local instability driving term is more strongly peaked at the equator, but the compressional stabilizing term is proportional to the flux average of the perturbed electrostatic potential, which is still averaged along the entire field line, however strongly the field lines become stretched and tail-like. Here, the long-thin perturbation means that the perpendicular wavelength of the perturbation is much smaller than the parallel wavelength. Therefore, the local ballooning driving term dominates the global compressional stabilizing term near the equator. Thus, the incompressible assumption is more likely to be satisfied in the tail-like equilibrium near the equator. Away from the equator the contribution to the energy integral due to the field line tension may be more important than the compressibility contribution to the energy integral. Therefore, the incompressible assumption seems to be more valid in the tail-like equilibrium.

The Energy Principle adopted by Lee and Wolf (1992) is consistent with the eigenmode analysis adopted by Miura et al. (1989) when the energy mimimization condition is taken into account and the compressibility term is retained (Miura 1999). In order to find more quantitatively the conditions under which the compressibility can be neglected for the ballooning instability, Miura (1999) calculated each energy term in the Energy Principle for the tail-like equilibrium and has shown that the compressible stabilizing term can be neglected for the tail-like equilibrium, when the ballooning eigenmode is strongly localized near the equator. More specifically, it has been found that when the plasma β at the equator is much larger than $6/\Gamma$, where Γ is the ratio of specific heat and is equal to 5/3 for the adiabatic 3-D plasma, the compressible stabilizing term can be neglected for the strongly localized mode, which is strongly peaked at the equator and is decaying exponentially toward the ionosphere. It has also been found that the interchange mode is completely stabilized by the compressibility when the plasma β at the equator is larger than $6/\Gamma$. This is consistent with findings of Horton et al. (1999) that the interchange mode is stabilized by the compressibility when the plasma β at the equator is larger than 1.5-3.0. The difference of the critical β at the equator for the interchange mode between the present investigation and Horton et al. (1999) is due to the difference of the specific field models. It has also been found that for the plasma β at the equator much smaller than $6/\Gamma$, the interchange (flute) mode becomes essentially incompressible. We note that for the ballooning instability the ionospheric boundary condition is such that the electrostatic potential perturbation is zero at the ionosphere or the field line displacement is zero at the ionosphere. For the interchange mode the ionospheric boundary condition is such that the field-aligned derivative of the electrostatic potential perturbation is zero at the ionosphere, which is satisfied for zero ionospheric conductivities. When the compressibility can be negleted

Space Science Reviews **95**: 387-398, 2001

for the ballooning instability, the only stabilizing factor is the field line tension. Therefore, the condition for the ballooning instability in the incompressible limit becomes a familiar form that the plasma β at the equator must exceed the product of the pressure gradient scale length and the field line curvature radius at the equator divided by the square of the parallel scale length of the ballooning mode (Miura et al. 1989, Roux et al. 1991a,b).

Since the unstable ballooning eigenmode found in Miura et al. (1989) (see their Figure 7) using the incompressible assumption satisfies the condition that the plasma β at the equator is much larger than $6/\Gamma$ and the unstable ballooning mode is strongly localized near the equator, the incompressible assumption in their calculation seems to be justified a posteriori. Therefore, the instability is expected to occur when the critical β calculated using the incompressible assumption is exceeded by the increase of the near-Earth plasma β.

5.2. Stochastic Plasma

When the field line becomes very tail-like, so that the field line curvature radius near the equator becomes comparable to the ion Larmor radius, the ion motion becomes stochastic and the stochastic ion dynamics become important and must be taken into account in the ballooning stability analysis (Hurricane et al. 1995). The stochastic ion dynamics may be important even as close as in the near-Earth plasma sheet (Lui et al. 1992).

According to the formulation of Hurricane et al. (1995), the ballooning eigenmode equation including the stochastic dynamics is similar to the ideal MHD eigenmode equation, but the compressibility term, which is expressed with a weighted flux average of the electrostatic potential perturbation, must be replaced by a weighted flux average of the electrostatic potential perturbation with the weighting function including the field-line curvature. This indicates that the effective compression term in the stochastic plasma can also be represented by the weighted flux average of the electrostatic potential perturbation with the weighting function including the field-line curvature. Since the stochastic dynamics are only important where the curvature radius is very small near the equator, we reasonably assume that the curvature radius is much smaller than the perpendicular gradient-B scale length near the equator and the integral only in the interval near the equator contributes most to the weighted flux average of the electrostatic potential perturbation. Then, it is shown (Miura 1999) that there is no unstable interchange mode with a constant electrostatic potential perturbation and the interchange mode is only marginal in the stochastic plasma. This is because the interchange mode is completely stabilized by the effective compressibility in the stochastic plasma. It is also shown that, the ballooning mode in the stochastic plasma cannot be completely stabilized by the effective compressibility alone. For a strongly localized mode, which is strongly peaked at the equator and is exponentially decaying toward the ionosphere, the effective compressional term in the stochastic plasma can be neglected and the ballooning eigenmode equation under such a condition in the stochastic plasma becomes the same as the ideal incompressible MHD eigenmode equation. This is very reasonable, because the stochastic dynamics are to change the equation of state from the adiabatic equation of state used in the ideal MHD analysis.

Space Science Reviews **95**: 387-398, 2001

Since the incompressible treatment does not use the equation of state and instead uses a mechanical equation that the divergence of the perturbed velocity vector is zero for closure of the fluid equations, the incompressible equation should be valid in the appropriate limits considered above irrespective of whether the plasma is adiabatic or stochastic. It has also been found that there is no condition for plasma β for the validity of the incompressible assumption in the stochastic plasma, although there is a critical β, which is set by the stabilizing tension force. By comparing the eigenmode equations for the ideal MHD and stochastic plasmas it has been found that for the plasma β much larger than $6/\Gamma$ the stochastic plasma is less stable than the ideal adiabatic MHD plasma. This is consistent with the findings of Hurricane et al. (1995).

6. Relation of the Ballooning Instability to the Auroral Break Up

By including in the ballooning eigenmode equation an ion diamagnetic drift term, which is the most dominant non-ideal MHD term, Miura et al. (1989) showed that the ballooning instability destabilizes the drift Alfvén mode (Tamao, 1984) in the magnetosphere, which is propagating westward. Ohtani et al. (1989a,b) further extended the eigenmode analysis of Miura et al. (1989) by employing full two-fluid equations. Pu et al. (1997) also used two-fluid equations in their study of the ballooning instability based on the local approximation. The real frequency of the unstable drift ballooning mode due to the presence of the ion diamagnetic drift term in the eigenmode equation is important and should be taken into account in understanding the substorm onset phenomena. The ion diamagnetic drift effect, which is a dominant correction to the ideal MHD equations, is responsible for the westward drift of the unstable drift Alfvén wave and this westward drift may be relevant to the inherent westward drift of the travelling auroral surge.

A consideration of the non-ideal MHD and kinetic effects (e.g. Chen and Hasegawa 1991, Cheng and Lui 1998) on the instability may further suggest the relevance of the ballooning instability to the observed onset phenomena. The strongly localized electrostatic potential perturbation, which is peaked at the equator and is decaying exponentially toward the ionosphere, may cause a parallel electric field for auroral particle accelaration when the full kinetic effects are taken into account, although in the ideal MHD approximation the parallel electrostatic electric field is cancelled by the parallel inductive field, with a consequence that there is no net parallel electric field. The shear in the ion diamagnetic drift velocity is known to cause a Kelvin-Helmholtz instability (Yoon et al., 1996), but coupled with the ballooning instability it might cause complexity in auroral deformation following the onset. Thus at such a kinetic level the ballooning instability may naturally be consistent with the auroral brightening and deformation in the magnetosphere-ionosphere coupling system following the onset.

7. Nonlinear Growth of the Ballooning Instability

Recently, a comprehensive study of the nonlinear development of the ballooning instability in the near- Earth plasma sheet has been carried out by Pritchett and Coroniti (1999) by using a 3-D full particle simulation and their simulation shows that the

Space Science Reviews **95**: 387-398, 2001

near-Earth plasma sheet does indeed become subject to the ballooning instability, when the plasma β exceeds a critical β calculated by using the incompressible assumption (Miura et al. 1989, Roux et al. 1991a,b). Their simulation demonstrates clearly that the westward propagating drift Alfvén wave, which is diamagnetic in character (the pressure perturbation and the magnetic pressure perturbation are out of phase), can be destabilized by the ballooning instability in the parameter range predicted by Miura et al. (1989). Thus, their simulation supports the validity of the incompressible assumption and the linear analysis of Miura et al. (1989). Although the nonlinear evolution of the ballooning instability clearly shown in their simulation is beyond the scope of the present linear investigation, it is conjectured that during the development of the ballooning instability, the curvature of the field line and the pressure gradient, which drive the instability, are reduced, because any instability evolves to reduce the inhomogeneity in the unperturbed state, which drives the instability. Therefore, it is suggested that the ballooning instability may be a dipolarization process.

8. Discussions and Conclusion

The growth rate of the incompressible ballooning instability in the tail-like equilibrium is proportional to the square root of the sum of the ion and electron temperatures divided by the product of the perpendicular pressure gradient scale length and the curvature radius at the equator. If we assume, according to observational results of Korth et al. (1991) and Pu et al. (1992), that the ion temperature, which is much larger than the electron temperature, is 1keV and the pressure gradient scale length and the field line curvature radius at the equator in the near-Earth plasma sheet are both 10000km, the e-folding time of the ballooning mode growth becomes 32sec. This e-folding time is fast enough to account for the rapid onset of the substorm expansion phase. Notice that although the present incompressible ballooning instability arises in the shear Alfvén branch, this e-folding time is much smaller than the bounce time of the Alfvén wave between the ionospheres of both hemispheres. This is because the Alfvén wave is trapped in a region around the equator (effective potential well) and is not bouncing back and forth between the ionospheres. According to the particle simulation of Pritchett and Coroniti (1999), the growth rate and the real frequency of the unstable drift ballooning mode due to the ion diamagnetic drift effect are comparable. Therefore, we expect that the wave period of the unstable drift ballooning mode is a few tens of seconds. Holter et al. (1995) found in situ diamagnetic hydrodynamic oscillations with periods of ~ 45-65sec, which are observed during the most active phase of the substorm breakup. The wave period and the diamagnetic character of these observed hydromagnetic oscillations are consistent with diamagnetic drift ballooning modes, which are strongly localized near the equator (Miura et al. 1989).

According to Lui et al. (1992) the plasma β as large as ~70 has been observed in the near-Earth plasma sheet before the substorm onset. Therefore, the plasma β much larger than one may not be a rare case in the near-Earth plasma sheet before the substorm onset. Therefore, the present result in section 5 showing that the ideal MHD ballooning mode, which is strongly localized near the equator, is essentially incompressible for the plasma β at the equator much larger than $6/\Gamma$ seems to be applicable to the near-Earth plasma sheet before the substorm onset. As we have seen in

Space Science Reviews **95**: 387-398, 2001

section 5.2, if the plasma is stochastic in the near-Earth plasma sheet, the same conclusion holds irrespective of the plasma β.

Although the initial perturbation or seed for the ballooning instability in the particle simulation of Pritchett and Coroniti (1999) is provided by the noise inherent in the particle simulation, in the real magnetospheric plasma the finite amplitude initial perturbation is necessary for the growth of the ballooning instability. This fact may explain the somewhat sporadic occurrence of the substorm onset, which has been reported to be influenced by internal triggers and external triggers outside the magnetosphere. It might also be conjectured that an earthward flow induced by the reconnection in the downstream region plays a role of the initial perturbation for the ballooning instability in the near-Earth plasma sheet.

We should mention here that the stability analysis of the ballooning instability in the tail-like configuration has also been performed by Bhattacharjee et al. (1998), but with a different conclusion. They showed by an eigenmode analysis that a 2-D magnetotail, obtained by 2-D time dependent simulations of the magnetotail in the high-Lundquist-number regime, is subject to an ideal MHD compressible ballooning instability (symmetric mode) with a high azimuthal wave number. They also showed that the same magnetotail configuration is not subject to the ideal MHD incompressible ballooning instability. Lee (1998) also supports their conclusion. Since the compressibility has always a stabilizing influence on the ballooning instability, it is not certain why the calculation of Bhattacharjee et al. (1998) showed that the same equilibrium is subject to the compressible ballooning instability but is not subject to the incompressible ballooning instability, because if an unstable compressible mode is found by the eigenmode analysis, that mode should become an unstable trial function in the Energy Principle, in which the compressible stabilizing term is neglected. The validity of using the dynamic equilibrium configuration as an unperturbed state in the ballooning stability analysis, whose eigenmode equation is obtained for the static equilibrium, is not guaranteed. Hence, the reason why Bhattacharjee et al. (1998) could not find an incompressible unstable mode may be due to the fact that they used a dynamic equilibrium solution instead of the static equilibrium as an unperturbed state.

The present results justify the incompressible assumption made in the ballooning stability analysis of Miura et al. (1989) for a tail-like equilibrium. Furthermore, they suggest that when the critical β due to the stabilizing tension force is exceeded, the ballooning instability is a viable instability in the near-Earth plasma sheet, which is strongly localized near the equator and may become essentially incompressible before the substorm onset. We have also found that the growth rate of the ballooning instability is larger nearer to the earth, the ballooning mode has a large azimuthal wave-number, and the growth time (e-folding time) of the ballooning instability becomes a few tens of seconds. All of these characteristics of the unstable drift ballooning mode seem to satisfy the requirements in (1) described in section 2 for the onset mechanism, i.e., the near-Earth location, the localized nature of the onset phenomena in the local time, and the time scale of the onset phenomena. Furthermore, the robustness of the ballooning mode against the stochastic dynamics, in contrast to the susceptibility of the interchange mode, suggests that the ballooning mode may be relevant to the near-Earth onset of substorms.

Although a complete scenario of the substorm onset and the subsequent

Space Science Reviews **95**: 387-398, 2001

physical processes has not been presented in the present investigation owing to lack of understanding of the specific relation of the ballooning instability to the global features in (2) described in section 2, i.e., to the global reconfiguration of the magnetosphere and to the energy release of the magnetic energy stored in the highly stretched tail field lines, the present investigation suggests at least the viability of the ballooning instability in the near-Earth plasma sheet prior to the substorm onset and its possible role in the near-Earth onset of the expansion phase of substorms.

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References

Bernstein, I.B., Frieman, E.A., Kruskal, M.D., and Kulsrud, R.M. 1958, *Proc. Roy. Soc. (London), A277*, 17.

Bhattacharjee, A., Ma, Z.W., and Wang, X. 1998, *Geophys. Res. Lett.*, **25**, 861.

Büchner, J., and Zelenyi, L.M. 1989, *J. Geophys. Res.*, **94**, 11821.

Chen, J. 1992, *J. Geophys. Res.*, **97**, 15011.

Chen, L., and Hasegawa, A. 1991, *J. Geophys. Res.*, **96**, 1503.

Cheng, C.Z., and Lui, A.T.Y. 1998, *Geophys. Res. Lett.*, **25**, 4091.

Freidberg, J.P. 1987, *Ideal Magnetohydrodynamics*, Plenum Press (New York), p.231.

Gold, T. 1959, *J. Geophys. Res.*, **64**, 1219.

Holter, Ø., Altman, C., Roux, A., Perraut, S., Pedersen, A., Pecseli, H., Lybekk, B., Trulsen, J., Korth, A., and Kremser, G. 1995, *J. Geophys. Res.*, **100**, 19109.

Horton, W., Wong, H.V., and Van Dam, J.W. 1999, *J. Geophys. Res.*, **104**, 22745.

Hurricane, O.A., Pellat, R., and Coroniti, F.V. 1995, *J. Geophys. Res.*, **100**, 19421.

Kaufmann, R.L. 1987, *J. Geophys. Res.*, **92**, 7471.

Korth, A., Pu, Z.Y., Kremser, G., and Roux, A. 1991, in *Magnetospheric Substorms*, Geophys. Monogr. Ser., vol. 64, T.A. Potemra, T. Iijima, S. Kokubun, and J.R. Kan, eds. (Washington D.C.: American Geophysical Union), p. 343.

Lee, D.-Y. 1998, *Geophys. Res. Lett.*, **25**, 4095.

Lee, D.-Y., and Wolf, R.A. 1992, *J. Geophys. Res.*, **97**, 19251.

Lembege, B., and Pellat, R. 1982, *Phys. Fluids*, **25**, 1495.

Lui, A.T.Y., Lopez, R.E. Anderson, B.J. Takahashi, K. Zanetti, L.J. McEntire, R.W. Potemra, T.A. Klumpar, D.M. Greene, E.M., and Strangeway, R. 1992, *J. Geophys. Res.*, **97**, 1461.

Miura, A. 1999, submitted to *J. Geophys. Res.*

Miura, A., Ohtani, S., Tamao, T. 1989, *J. Geophys. Res.*, **94**, 15231.

Ohtani, S., Miura, A., and Tamao, T. 1989a, *Planet. Space Sci.*, **37**, 567.

Ohtani, S., Miura, A., and Tamao, T. 1989b, *Planet. Space Sci.*, **37**, 579.

Pellat, R., Coroniti, F.V., and Pritchett, P.L. 1991, *Geophys. Res. Lett.*, **18**, 143.

Pritchett, P.L., and Coroniti, F.V. 1999, *J. Geophys. Res.*, **104**, 12289.

Pu, Z.Y., Korth, A., and Kremser, G. 1992, *J. Geophys. Res.*, **97**, 19341.

Pu, Z.Y., Korth, A., Chen, Z.X., Friedel, R.H.W., Zong, Q.G., Wang, X.M., Hong, M.H., Fu, S.Y., Liu, Z.X., and Pulkkinen, T.I. 1997, *J. Geophys. Res.*, **102**, 14397.

Space Science Reviews **95**: 387-398, 2001

Roux, A., Perraut, S., Morane, A., Robert, P., Korth, A., Kremser, G., Pedersen, A., Pellinen, R., and Pu, Z.Y. 1991a., in *Magnetospheric Substorms, Geophys. Monogr. Ser.*, vol. 64, T.A. Potemra, T. Iijima, S. Kokubun, and J.R. Kan, eds. (Washington, D.C.: American Geophysical Union), p. 201.

Roux, A., Perraut, S., Robert, P., Morane, A., Pedersen, A., Korth, A., Kremser, G., Aparicio, B., Rodgers, D., and Pellinen, P. 1991b, *J. Geophys. Res.*, **96,** 17697.

Southwood, D.J., and Saunders, M.A. 1985, *Planet. Space Sci.*, **33,** 127.

Tamao, T. 1984, *Planet. Space Sci.*, **32,** 1371.

Viñas, A.F., and Madden T.R. 1986, *J. Geophys. Res.*, **91,** 1519.

Yoon, P.H., Drake, J.F., and Lui, A.T.Y. 1996, *J. Geophys. Res.*, **101,** 27327.

Figure captions

Fig. 1. Schematic view of the dipole-like field and the tail-like field. (a) In the dipole-like field, $\varepsilon \gg 1$ holds and the pressure gradient force (black arrow) is nearly balanced with the tension force (white arrow) due to the field line curvature. (b) In the tail-like field, $\varepsilon \ll 1$ holds on average except near the equator and the pressure gradient force (black arrow) is nearly balanced with the magnetic pressure force (white arrow) away from the equator.

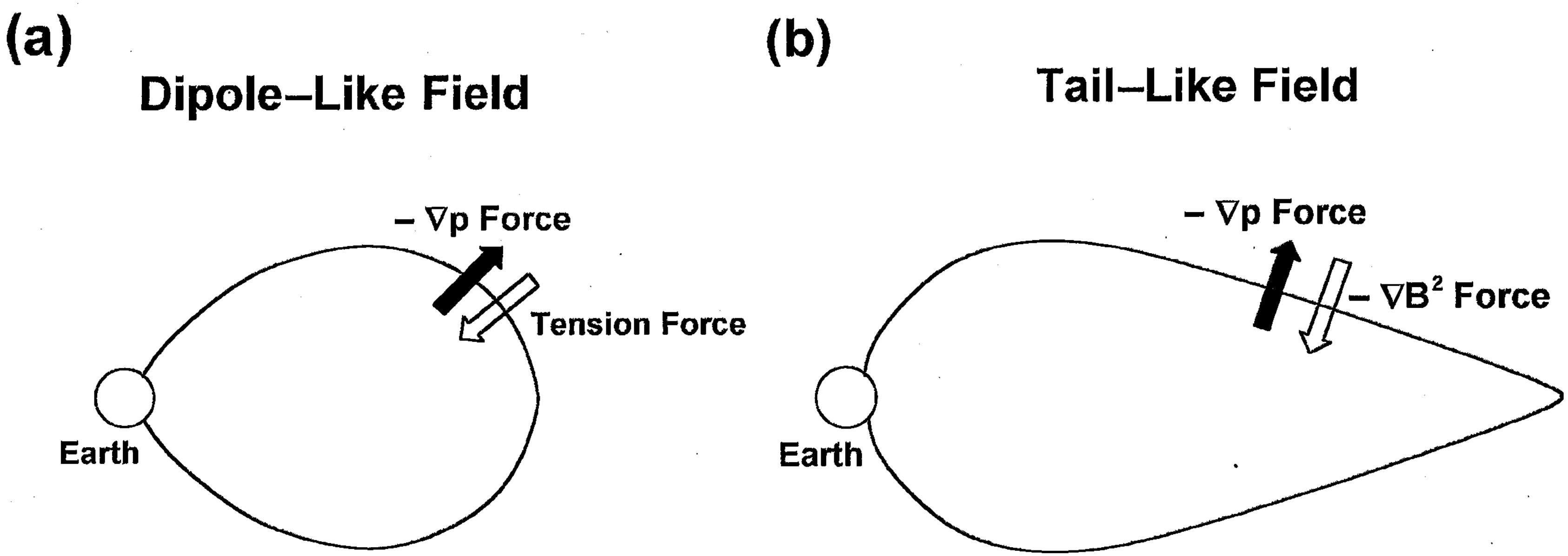
(a)
Dipole–Like Field
– ∇p Force
Tension Force
Earth
(b)
Tail–Like Field
– ∇p Force
– ∇B² Force
Earth

Fig. 1