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BALLOONING-TYPE INSTABILITIES AND WAVES IN THE EARTH'S MAGNETOSPHERE (REVIEW)

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Abstract. Ballooning-type mechanism is highlighted for the interpretation of three intriguing magnetospheric phenomena. These are: (1) hot filament extensions from the plasma sheet into the lobes and nightside polar cap arc formation; (2) destabilisation of the near-Earth tail current, accompanied by auroral arc activation prior to substorm onset; (3) flapping waves, as observed by Cluster spacecraft.

1. Ballooning instability at the plasma sheet-lobe interface and its implications for polar cap arc formation.

Huang et al. [1987, 1989] reported about hot filaments of plasma sheet origin filling the magnetospheric lobes during northward IMF. At the same time, cold tenuous transients are often observed in the plasma sheet. These features can be interpreted as a result of plasma exchange at the plasma sheet/lobe interface (PSLI) proceeding in a filamentary manner. We relate this process to destabilizing of the PSLI by the low-frequency ballooning. Although ballooning instability conditions are not typically met inside the plasma sheet [*Lee and Wolf,* 1992], they can be satisfied at the PSLI. Indeed, the near-Earth curved segment of the PSLI denoted as AA' in Figure 1, where it is shown for extremely quiet (left panel) and moderately quiet (right panel) conditions according to the T96 model, always separates the cold lobe population from the hot plasma sheet. Thus it has a pressure gradient in favour for the development of the instability and can become a generator region launching filamentation.

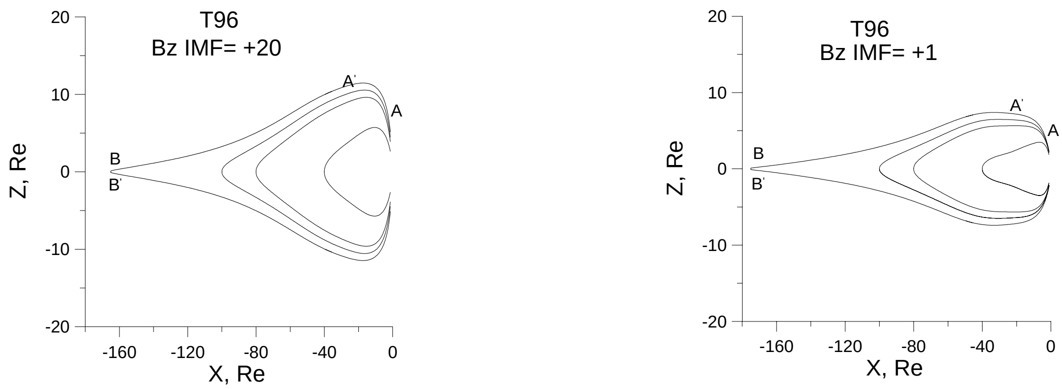

Figure 1. Statistical magnetic configurations for strong (left) and weak (right) northward IMF according to the T96 model, for $y_{GSM} = 0$ and zero dipole tilt. The near-Earth (AA′) curved site of the boundary magnetic field tube can become ballooning/interchange unstable. Note an increase in the near-Earth curvature of the boundary plasma sheet with increasing northward IMF Bz.

Previously, the dispersion relation for the ballooning mode considered in the local approximation was obtained by *Ohtani and Tamao* [1993] (in this paper we do not address the ballooning-type mode related to the field-aligned velocity perturbations [*Liu,* 1997]). It has the form

$$\omega^2 = V_A^2 k_{||}^2 - \omega_g^2 \qquad (1)$$

where V_A is the Alfvén velocity, $k_{||}$ is the wave vector component along the magnetic field and ω_g is given by

$$\omega_g^2 = \frac{\beta V_A^2}{1+\frac{\gamma\beta}{2}} k_c[k_p - \gamma(k_b + k_c)] \qquad (2)$$

where β is the ratio of the plasma pressure to the magnetic pressure, γ the adiabatic exponent, $k_c = 1/R_c$, R_c being the curvature radius, k_p and k_b are the reversed scale lengths of the pressure and magnetic field inhomogeneity in the background configuration. In deriving (1) it was supposed that the perpendicular component of the wave vector is directed along the y-axis and greatly exceeds any other scale size of the problem.

We note that the curvature of the generator region AA' is not very large. Further for the estimates we will take the curvature radius here to be $R_c \sim 20\ R_E$, R_E being the Earth's radius, that is $k_c \sim 0.05\ R_E^{-1}$. The curvature of the rest part of the boundary magnetic field tube, i.e. tailward from the point A', is neglected. (The near-equatorial curved site BB' is ignored either, for the magnetic field is singular around it, so that the total magnetic curvature is unclear. Besides, as differs from AA', there is no indication of strong pressure gradients necessary for a ballooning growth there, as the plasma sheet continues tailward of the last closed field line [*Song et al.,* 1999; *Kullen and Janhunen,* 2004]). At the same time, the scale length of the *e* fold pressure increase toward the curvature center in the near-Earth boundary plasma sheet is much smaller than R_c and hardly exceeds 2 R_E. This yields $k_p \sim 0.5\ R_E^{-1}$. Taking $\beta = \dfrac{2\mu_0 p}{B^2} \sim 0.5$ for the region considered (*p* being the plasma pressure, B the magnetic field, μ_0 the magnetic permeability of vacuum), we find from the force-balance condition $\beta/2\ k_p + k_b - k_c = 0$ that $k_b \sim -0.075\ R_E^{-1}$. The negative sign of k_b indicates that the magnetic field grows from the curvature center, i.e. towards the lobe. From eq. (2) it is easy to see that for the adopted values of the equilibrium parameters ω_g^2 is positive. Its magnitude can be estimated as $\omega_g \sim 0.015\ s^{-1}$.

Further we will roughly take that the equilibrium characteristics (V_A, k_c, etc.) do not change along the magnetic field line. At the ionospheric end of the magnetic tube (point A) the boundary condition of *Maltsev* [1977] and *Glassmeier* [1983] will be set, which states that the field-aligned currents in the ballooning perturbations must be closed by the ionospheric current

$$\frac{\partial\varphi}{\partial s}\Big|_A = i\omega\mu_0\Sigma_P\varphi \qquad (3)$$

where φ is the transverse displacement of the magnetic tube multiplied by the metric factor √B, Σ_P the height-integrated ionospheric Pedersen conductivity, the coordinate s is directed along the magnetic field.

At the point A' we will impose a condition of the emitting Alfvén wave

$$\frac{\partial\varphi}{\partial s}\Big|_{A'} = -i\omega\mu_0\Sigma_A\varphi \qquad (4)$$

where $\Sigma_A = 1/(\mu_0\ V_A)$ is the wave conductivity. Condition (4) implies that the perturbed field-aligned currents tailward of A' close via the Alfvén front going to infinity. As will be shown, the growth time of the instability we further reveal is by an order of magnitude shorter than the time an Alfvén wave travels along the last closed field line to the opposite ionosphere. (The configurations shown in Figure 1 suggest it to be about 1 hour). We will also argue that at the non-linear stage the instability results in hot tubes detached from the plasma sheet and progressing towards the dayside cusp. It will be shown that the foot of a tube travels a distance much larger than its diameter in the time required for an Alfvén wave to reach the opposite ionosphere and come back. All this makes condition (4), which presupposes that there is no essential Alfvén-wave communication with the opposite ionosphere, reasonable. Further we will adopt $\Sigma_A \sim 0.8$ S for the boundary plasma sheet and $\Sigma_P \sim 5$ S for its ionospheric footprint.

It can be shown that conditions (3), (4) yield one more relation between the frequency ω and field-aligned wave number $k_{||}$ of the perturbations

$$(k_{||}^2+\omega^2\ \mu_0^2\ \Sigma_P\Sigma_A)\ \mathrm{tg}\ k_{||}\ a = -\ i\omega\ \mu_0\ (\Sigma_P+\Sigma_A)\ k_{||} \qquad (5)$$

where the length of the curved site AA' is denoted as *a* (*a* ~ 20 R_E).

In the standing structure determined by (1) and (5) the most promising from the point of view of instability growth is the zeroth harmonics ($k_{||}\ a << 1$), for it is for this harmonics that the stabilizing term $V_A^2\ k_{||}^2$ in dispersion relation (1), proceeding from magnetic bending, is minimum.

For the zeroth harmonics eq. (5) takes the form

$$(k_{||}^2 + \omega^2\ \mu_0^2\ \Sigma_P\ \Sigma_A)\ a = -\ i\ \omega\ \mu_0\ (\Sigma_P + \Sigma_A) \qquad (6)$$

Substituting $k_{||}$ from (1) into (6), we have

$$[\omega_g^2+\omega^2\ (1+\mu_0^2 V_A^2\Sigma_P\ \Sigma_A)]\ a = -\ i\ \omega\ \mu_0\ V_A^2(\Sigma_P+\Sigma_A) \qquad (7)$$

Note that in eq. (7) both bending and inertial terms for the zeroth harmonics are preserved. That is, eq. (7) is a rigorous dispersion relation for the considered mode. Only starting from this step we will restrict ourselves with the low frequency limit, such that

$$\omega^2 << \frac{\omega_g^2}{(1+\dfrac{\Sigma_P}{\Sigma_A})} \qquad (8)$$

Then the second term in the LHS of (7) can be neglected and we obtain the solution

$$\omega = \frac{i\omega_g^2 a}{V_A(1+\dfrac{\Sigma_P}{\Sigma_A})} \qquad (9)$$

For the chosen characteristic values of the parameters, $\omega = \dfrac{i}{250s} = 0.004\ s^{-1}$, which satisfies inequality (8). Generally, the larger is the ratio of the ionospheric conductivity Σ_P to the wave conductivity Σ_A, the better is the accuracy with which inequality (8) holds. We can also see that the growth time of the instability is by an order of magnitude shorter than the travel time of the Alfvén wave to the opposite ionosphere and, consequently, condition (4) is justifiable.

As was shown above, ω_g^2 is positive for the near-Earth boundary plasma sheet and thus (9) describes an unstable mode. This mode has both ballooning and interchange sense. Formally, $k_{||}$, though small, is not zero for it, which is a feature of ballooning perturbations. On the other hand, solution (9) describes aperiodic regime, for which field-aligned Alfvén propagation and associated inertial currents are not essential, and this is a feature of interchange motions.

We have found that the growth time of the instability under study is as short as 4 minutes. This suggests that most of the time it evolves in the non-linear regime, which will further be considered.

At the non-linear stage, plasma sheet extensions into the lobes can either be filaments, i.e. hot tubes detached from the plasma sheet and making up roads along convection streamlines, or solid sheets in the XZ plane. If we carefully look at the snapshots or images of extension footprints in the ionosphere (further Figure 2), we will see that they are never solid or smooth but always made up of smaller-scale features of increased luminosity. This observation infers filamentary rather than solid manner of extension. The tubes detached from the plasma sheet are still curved in the near-Earth region and filled with hot plasma. Then they tend to continue their motion away from the plasma sheet, as the final cause of this motion is a charge separation due to a divergence of energetic ion drift current in a curved/inhomogeneous magnetic field. A similar evolution of the interchange instability was observed by *Alfvén* [1981] in a laboratory plasma.

Let us find the footprint velocity v_{hi} of a hot filament with a circular cross-section in the reference frame of the background convection. Previously, similar treatments of transient motions were performed by *Lyatsky and Maltsev* [1984] and *Pontius and Wolf* [1990] for the Earth's plasma sheet and by *Pontius and Hill* [1989] for Iogenic plasma.

The divergence of the gradient-curvature current $\mathbf{j}_\perp$ is given by

$$\mathrm{div}\,\mathbf{j}_\perp = \frac{1}{B^2}\,\mathrm{div}\,[\mathbf{B}\times\nabla p] + \nabla\left(\frac{1}{B^2}\right)\cdot[\mathbf{B}\times\nabla p] \qquad (10)$$

Noticing that $\mathrm{div}\,[\mathbf{B}\times\nabla p] = \nabla p\cdot\mathrm{rot}\,\mathbf{B} - \mathbf{B}\cdot\mathrm{rot}\,\nabla p$, in the RHS of which both terms are zero, we have

$$\mathrm{div}\,\mathbf{j}_\perp = -\mathbf{B}\,\left[\nabla\left(\frac{1}{B^2}\right)\times\nabla p\right] = \mathbf{B}\,\left[\frac{\nabla B^2}{B^4}\times\nabla p\right] \qquad (11)$$

Using the force balance equation in the form $\nabla B^2 = -2\,\mu_0\,\nabla p + 2\,\mathbf{k}_c\,B^2$, we can rewrite (11) as

$$\mathrm{div}\,\mathbf{j}_\perp = \frac{2}{B}\,\mathbf{e}_b\,[\mathbf{k}_c\times\nabla p] \qquad (12)$$

where $\mathbf{e}_b$ is the unit vector along the magnetic field. Then the field-aligned current $j_{||}$ in the northern ionosphere is

$$j_{||} = -B_i \int \frac{\mathrm{div}\,j_\perp}{B} ds = -B_i \int \frac{2}{B^2}\,\mathbf{e}_b\,[\mathbf{k}_c\times\nabla p]\,ds \qquad (13)$$

where B_i is the magnetic field at the ionospheric level. It is seen that integration in (13) extends only over the near-Earth curvilinear segment of the hot tube, for which $k_c \neq 0$. It is also seen that the field-aligned current is positive (into the ionosphere) at the dusk side of the hot filament and negative (out of the ionosphere) at its dawn side.

Now we use the current continuity equation

$$-B_i \int \frac{2}{B^2}\,\mathbf{e}_b\,[\mathbf{k}_c\times\nabla p]\,ds = (\Sigma_P + \Sigma_A)\,\mathrm{div}\,\mathbf{E} \qquad (14)$$

where $\mathbf{E}$ is the polarization electric field to be found. We will consider the difference in pressure between the hot tube and the background to be a step function, the pressure of the lobe plasma taken zero. Then the solution of (14) is the electric field $\mathbf{E}$ which is uniform and directed from dusk to dawn inside the filament and subsides as a 2D dipole field from its boundary outward (e.g. [*Pontius and Wolf*, 1990]). Taking into account that the tangential component of the electric field is continuous at the filament boundary and the radial component suffers a jump equal to the double value of the electric field inside the filament, we can find the value of this latter field as

$$E = -\frac{B_i^{1/2}\,p}{(\Sigma_P + \Sigma_A)} \int \frac{k_c}{B^{3/2}}\,ds \qquad (15)$$

where it has been suggested that $|\nabla_\perp\, p| = (B/B_i)^{1/2}\,|\nabla_{\perp i} p|$. Correspondingly, the footprint velocity of a hot filament $v_{hi} = \dfrac{E}{B_i}$ can be written as

$$v_{hi} = -\frac{p}{(\Sigma_P + \Sigma_A)\,B_i^{1/2}} \int \frac{k_c}{B^{3/2}}\,ds \qquad (16)$$

in which the minus sign indicates that the motion is outward from the curvature centre. It can be shown that the velocity v_{hi} ranges from 100 m/s to 1 km/s, depending on which parameters are chosen. Taking the same values of the parameters as for the estimation of ω from (9), we get $v_{hi} \approx 200$ m/s. These values are consistent with the rates of sun-aligned arc protrusion into the polar cap (e.g., [*Hoffman et al.*, 1985]).

The background convection is imposed on the interchange motions. The large-scale convection associated with southward Bz IMF suppresses hot filament progression into the lobes, while promoting penetration of tenuous transients into the plasma sheet. In opposite, during northward IMF the convection may favour the extension of plasma sheet filaments into the lobes and their subsequent protrusion towards noon. This process is signified in the ionosphere by the occurrence of nightside originating polar arcs. Time evolution of a multiple arc event seen on the Polar UVI imager is illustrated in Figure 2. Thus, the low-frequency ballooning instability can lead to destabilization of the boundary plasma sheet and at the non-linear stage be responsible for nightside polar cap arc formation.

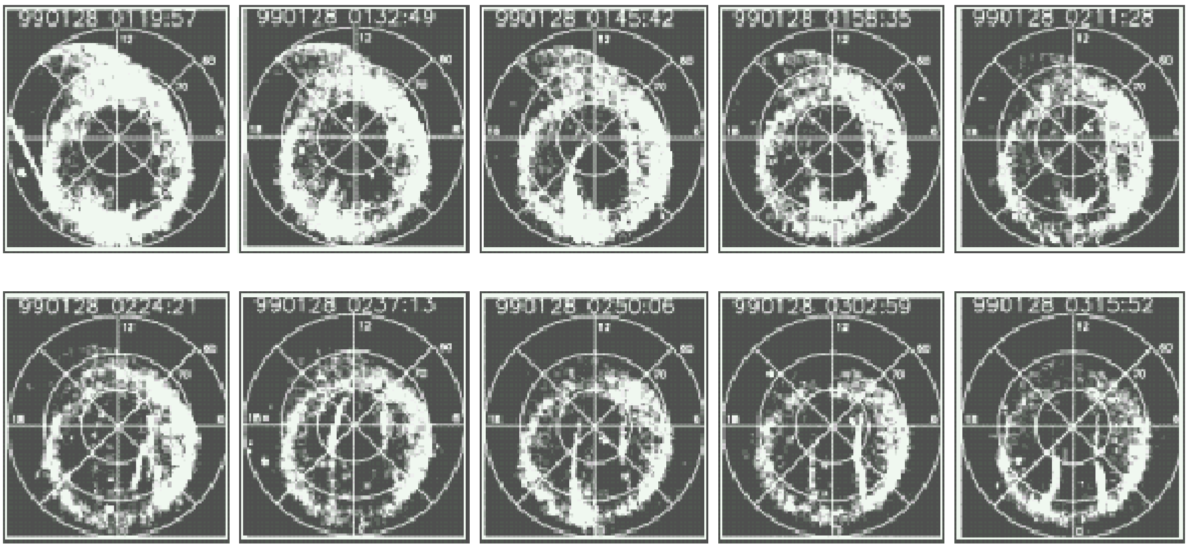

Figure 2. Time evolution of a multiple arc event seen on the Polar UVI imager.

2. Pressure azimuth asymmetry driven (or baroclinic) interchange instability as a trigger of substorm activation.

Efforts of numerous studies are directed toward finding an internal magnetospheric instability that could spontaneously lead to cross-tail current disruption at substorm onset within the inner plasma sheet. It is generally accepted that enhanced plasma transport to the Earth during substorm growth phase leads to a significant pressure buildup in the near- and middle tail at premidnight-to-postmidnight MLTs (sketched by the thin isocontours in Figure 3). Then a possibility of near-Earth tail current destabilisation by the ballooning/interchange mechanism, tending to destroy large pressure gradients forming in this region, may be quite plausible [e.g., *Swift*, 1967; *Roux et al*., 1991; *Ivanov et al.,* 1992; *Ohtani and Tamao*, 1993; *Liu*, 1997; *Golovchanskaya et al.*, 2004]. The stability of the radial pressure profile near geosynchronous orbit was first examined statistically by *Ohtani and Tamao* [1993]. The authors demonstrated that at substorm growth phase this profile does not develop earthward pressure gradients large enough to satisfy $k_p - \gamma (k_b + k_c) > 0$ criterion for the ballooning instability and thus refuted the basic ballooning as a possible substorm trigger. The radial pressure profile is also stable to interchange perturbations, which is signified by pV^{γ} (V is the unit magnetic flux tube volume) increasing tailward for both quiet and stretched magnetospheric configurations [*Lee and Wolf*, 1992]. Consequently, the manner of disruption of the near-Earth pressure buildup shown in Figure 3 (left panel) and originally suggested by *Roux et al*. [1991] can hardly be expected.

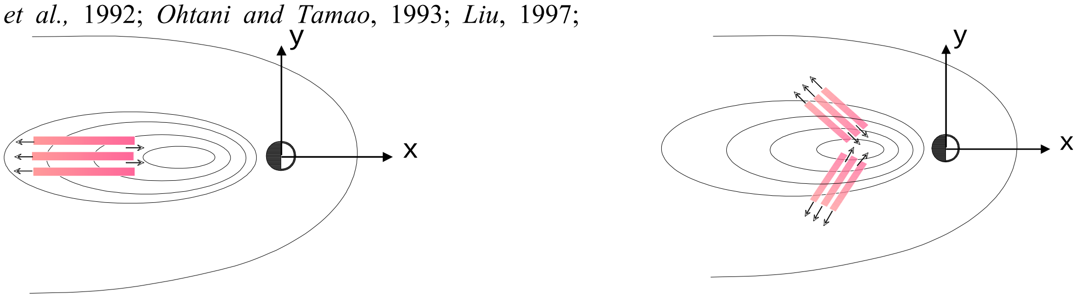

Figure 3. Disruption of the near tail pressure buildup at the late growth phase by the basic ballooning/interchange instability (left) and baroclinic interchange instability (right). Hot and dense plasma filaments are shown dark grey. Colder and more tenuous plasma features are blank. The arrows indicate directions of velocity perturbations inside the interchange structures. The solid lines are pressure buildup contours.

Later, the ballooning mechanism with regard to substorm triggering was revised by *Liu* [1997], who performed a more rigorous analysis of the MHD equations in the inhomogeneous case, the field-aligned plasma transport included. One more ballooning-type mode was revealed, which, unlike the basic ballooning, is essentially related to the field-aligned rather than transverse velocity perturbations and thus has more in common with a slow magnetosonic than a shear Alfvén wave. This mode has a looser destabilisation criterion and suggests rapid magnetic reconfigurations (on timescale δ 1 min) that may be relevant to the explosive growth phase of the substorm [*Liu*, 1997].

Alternatively, in [*Volkov and Maltsev*, 1986; *Golovchanskaya and Maltsev*, 2003] an interchange modification related to the longitudinal asymmetry in a plasma distribution (Figure 3, right panel) has been proposed (see also [*Liu*, 1996]). This mode is unstable whenever the magnetic tension force, acting on hot plasma in the curvilinear magnetic field, and plasma pressure gradient are not strictly anti-parallel. We term this mode as a 'baroclinic interchange' one, meaning a close analogy with the atmospheric baroclinic instability. The latter develops if atmospheric pressure gradient gets inclined to the gravity direction, owing, for example, to horizontally non-uniform heating. The measure of the directional displacement between ∇V and ∇p vectors in the magnetosphere is a background field-aligned current (FAC) generated by the *Vasyliunas* [1970] mechanism. For substorm growth phase the occurrence of the large-scale FACs in the near-to-middle tail at night MLTs that have a sense of the Region 2 FACs is well known [*Iijima et al.*, 1993; *Watanabe and Iijima*, 1993]. The average intensities of the FACs are increased up to ~ 10^{-6} A/m^2 when reduced to the ionosphere. In [*Golovchanskaya et al.*, 2004] we performed a numerical simulation of the decay of these FACs into smaller-scale structures.

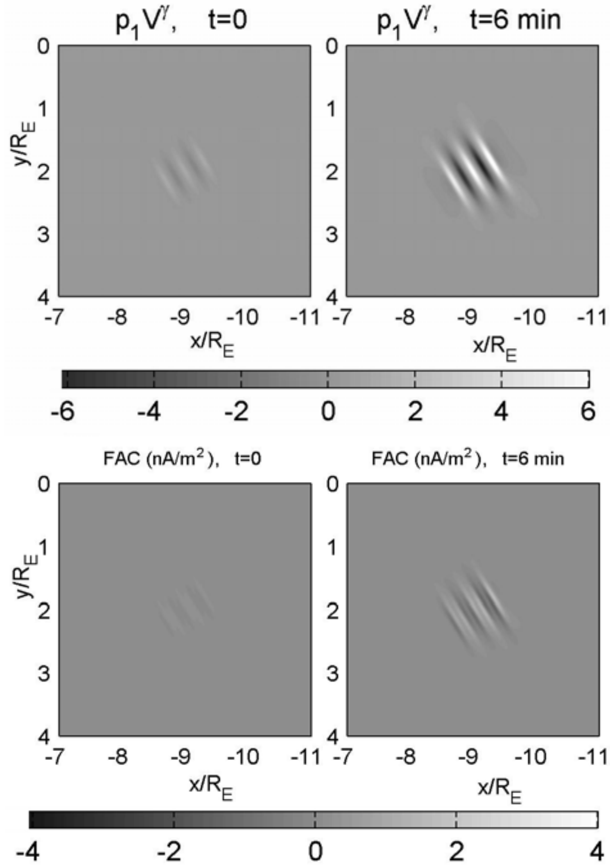

Figure 4. Perturbations in pV^{γ} (upper panels) and in field-aligned currents (lower panels) at t = 0 (left column) and t = 6 min (right column).

The run presented in Figure 4 reproduces the growth in time of small perturbations in pV^{γ} set in the premidnight sector as a ripple inclined to azimuth (upper panels) and of associated small-scale FACs (lower panels). One can see that by the 6th minute of evolution the perturbed FACs reach great values sufficient to cause brightening of accompanied discrete auroral arcs.

We conclude that the baroclinic interchange instability that develops in the near-Earth tail during substorm growth phase can result in cross-tail current structuring, the sites with more dipolarized magnetic field alternating radially and azimuthally with the regions of more tail-like field configuration. In the ionosphere this process is signified by auroral arcs growing in intensity.

3. Identification and theory of plasma sheet flapping waves observed by Cluster

An intriguing feature of the observed plasma sheet flapping waves, that is large-amplitude oscillations in the GSM B_x component traveling in the $\pm y_{\mathrm{GSM}}$ directions, is their propagation from the magnetospheric center toward both dawn and dusk flanks [*Runov et al.*, 2003; *Sergeev et al.*, 2003, 2004]. The group velocity ranges from a few tens of km/s for quiet conditions to a few hundreds of km/s for disturbed conditions. This velocity is much lower than the magnetosonic speed (~1000 km/s), making a number of MHD modes inappropriate for interpretation of the phenomenon (see [*Sergeev et al.,* 2004] and references therein). Nor can the flapping waves be explained by the modes related to ion drift motion, for such an explanation would be consistent only with duskward, but not dawnward, propagation. The kink mode in the Harris-like current sheet [*Volwerk et al.*, 2003] could be relevant to the standing oscillations of the neutral sheet but it fails to explain propagation toward the flanks.

Golovchanskaya and Maltsev [2005] argue that the flapping waves are ballooning waves propagating in hot inhomogeneous plasma confined by a curvilinear magnetic field. The theory of such waves was previously considered in [*Safargaleev and Mal*tsev, 1986]. However, a suggestion of small magnetic curvature (low β), used in that treatment, does not permit to directly apply the results for the interpretation of Cluster observations related to strongly stretched magnetic configurations in the magnetospheric tail. *Golovchanskaya and Maltsev* [2005] generalized the ballooning wave theory and obtained the following dispersion relation

$$\omega^2 = V_A^2 k_\parallel^2 + \frac{V_A^4 k_\parallel^2}{2\gamma c_s^2} + V_A^2 k_c \frac{k_y^2}{k_\perp^2}(k_b + k_c \pm \pm\sqrt{(k_b+k_c)^2 + 4k_\parallel^2 \frac{k_\perp^2}{k_y^2} + \left(\frac{V_A^2 k_\parallel^2}{2\gamma k_c c_s^2}\frac{k_\perp^2}{k_y^2}\right)^2}\,) \qquad (17)$$

where $c_s^2 = \frac{p}{\rho}$, ρ being the mass density, $k_\perp^2 = k_n^2 + k_y^2$, k_n being the component of the wave vector of the perturbation along the direction of magnetic curvature.

In case β >> 1, eq. (17) simplifies and takes the form

$$\omega^2 = V_A^2 k_\parallel^2 + V_A^2 k_c (k_b + k_c \pm \pm\sqrt{(k_b+k_c)^2 + 4k_\parallel^2 \frac{k_\perp^2}{k_y^2}}\,)\frac{k_y^2}{k_\perp^2} \qquad (18)$$

At this step in [*Golovchanskaya and Maltsev*, 2005] the second term under the square root was neglected (this can be done for sufficiently small $k_\parallel$), and the dispersion relation reduced to

$$\omega^2 = V_A^2 k_\parallel^2 + \omega_g^2 \frac{k_y^2}{k_n^2 + k_y^2} \qquad (19)$$

which those authors further analyzed. In (19) $\omega_g^2 = 2V_A^2 k_c(k_b + k_c)$ is the frequency determined solely by the equilibrium characteristics. The plus sign before the square root in (18) has been chosen as it refers to the case of propagation (real ω) that we are analyzing. The minus sign could lead to destabilizing of the ballooning mode [*Liu*, 1997], which is beyond the scope of the present study. The case with the term $4k_\parallel^2 \frac{k_\perp^2}{k_y^2}$ preserved in (18) should be further considered in more rigorous treatments.

From eq. (19) the group velocity in the $\pm y$ directions can be found as

$$v_y^{group} = \frac{\sqrt{2}\, v_T\, k_n^2}{h\left(k_n^2 + k_y^2\right)^{3/2}} \qquad (20)$$

where v_T is the average thermal velocity of ions in the neutral sheet, h the characteristic width of the current sheet. To estimate the group velocity, we take $k_n \sim k_y$, and thus get $v_y^{group} \approx (\sqrt{2})\, v_T/(3hk_y)$. The typical proton energy in the plasma sheet is 6 keV, which corresponds to the thermal velocity of ~1000 km/s. We suggest $h \sim 2\ R_E$ for the quiet current sheet and $h \sim 0.5\ R_E$ for the active sheet, and take from the observations $\lambda_y = 1\text{–}3R_E$, ($k_y = 2\pi/\lambda_y \approx 2 - 6\ R_E^{-1}$). Then v_y^{group} can be estimated as 40 km/s to 100 km/s for quiet conditions and 160 km/s to 400 km/s for active conditions. These values of v_y^{group} are consistent with the observed ones [*Sergeev et al.,* 2004].

Consequently, plasma sheet flapping waves are most likely to be ballooning waves propagating in hot plasma confined in curvilinear magnetic field. These waves are similar in nature to the internal gravitational waves in the atmosphere.

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