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THE ROLE OF MICROPROCESSES IN MACROSCALE MAGNETOTAIL DYNAMICS

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The role of microprocesses in macroscale magnetotail dynamics

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Abstract. Microphysical processes influence the behavior of macroscopic variables and hence the dynamical processes on larger scales through the isotropization of the pressure tensor and through anomalous transport processes which affect the conservation and transport of momentum, energy, and magnetic flux. Using MHD simulations of the dynamic evolution of the magnetotail with deviations from ideal, isotropic, and isentropic MHD, it is demonstrated that each of these deviations may have drastic effects on the stability and dynamic behavior of the magnetotail.

1. Introduction

Large-scale models of the Earth's magnetotail typically are based on the MHD approximation, which incorporates microphysics assumptions in various places. The purpose of this paper is to identify the places where microphysics assumptions enter MHD and to discuss effects of modifications of these approximations within the framework of large-scale magnetotail dynamics. Obviously, the restriction to large scales does not permit a direct comparison of the MHD model and its modifications with a particle model. The results presented in this paper therefore mainly serve to illustrate possibilities of how microphysics might affect the macroscale dynamics rather than conclusive evidence.

2. Basics Equations and Approximations

Microscopic phenomena in space plasmas are usually investigated on the basis of the Vlasov description or by a full particle approach, which become essentially equivalent when the number of particles within a Debye sphere is large and particle-particle interactions are negligible. (The differences between these two approaches are not the subject of this paper.) The Vlasov description consists of Maxwell's equations

$$\frac{\partial \mathbf{B}}{\partial t} = -\nabla \times \mathbf{E} \quad (1)$$

$$\nabla \cdot \mathbf{B} = 0 \quad (2)$$

$$\nabla \times \mathbf{B} = \mu_0 \left(\mathbf{j} + \epsilon_0 \frac{\partial \mathbf{E}}{\partial t} \right) \quad (3)$$

$$\nabla \cdot \mathbf{E} = \frac{1}{\epsilon_0} \sigma \quad (4)$$

combined with the kinetic equations

$$\frac{\partial f_j}{\partial t} + \mathbf{u} \cdot \frac{\partial f_j}{\partial \mathbf{x}} + \frac{e_j}{m_j} (\mathbf{E} + \mathbf{u} \times \mathbf{B}) \cdot \frac{\partial f_j}{\partial \mathbf{u}} = 0 \quad (5)$$

where $f_j(\mathbf{x}, \mathbf{u}, t)$ denotes the phase space distribution of species j with mass m_j and charge e_j , σ the total charge density, and mks units are used. The macroscopic or fluid description of the plasma is usually derived by taking velocity moment integrals of Equation (5), using suitable assumptions. For more details we refer to appropriate text books. Here we point out some of the major assumptions that could affect the macroscopic behavior, using generalized forms of the moment equations. For simplicity we restrict the discussion to two (singly charged) particle species, ions and electrons. In that case each velocity moment integration of (5) yields two equations, one for each particle species, which are typically added together, using m_j or e_j as multipliers.

The lowest order (in terms of velocity powers) equations are the continuity equations, representing conservation of mass and charge

$$\frac{\partial \rho}{\partial t} + \nabla \cdot \rho \mathbf{v} = 0 \quad (6)$$

$$\frac{\partial \sigma}{\partial t} + \nabla \cdot \mathbf{j} = 0 \quad (7)$$

where ρ is the mass density. For large-scale phenomena with phase speeds considerably less than the speed of light it is justified to assume quasineutrality

$$n_i \approx n_e \quad \text{or} \quad \sigma \approx 0 \quad (8)$$

In that case the first term in (7) can be dropped, which is consistent with also dropping the displacement current term in (3), while (4) is replaced by (8).

In next higher order we obtain the momentum equation

$$\rho \frac{d\mathbf{v}}{dt} \equiv \rho \left(\frac{\partial \mathbf{v}}{\partial t} + \mathbf{v} \cdot \nabla \mathbf{v} \right) = -\nabla \cdot \underline{\mathbf{P}} + \mathbf{j} \times \mathbf{B} \quad (9)$$

and the generalized Ohm's law

$$\mathbf{E} + \mathbf{v} \times \mathbf{B} = \eta \mathbf{j} + \frac{1}{n_e} \mathbf{j} \times \mathbf{B} + \frac{1}{n_e} \nabla \cdot \left(\frac{m_e}{m_i} \underline{\mathbf{P}}_i - \underline{\mathbf{P}}_e \right) + \frac{m_e}{n_e^2} \left[\frac{\partial \mathbf{j}}{\partial t} + \nabla \cdot (\mathbf{j} \mathbf{v} + \mathbf{v} \mathbf{j}) \right] \quad (10)$$

where $\underline{P}$ is the total pressure tensor and $\underline{P}_i$ and $\underline{P}_e$ the ion and electron pressure tensor, respectively, all evaluated in the plasma rest frame. In (9) and (10) quasineutrality has already been used, and in some places the electron mass has been neglected against the ion mass. We note that the ohmic term $\eta \mathbf{j}$ follows directly from the kinetic equation only if a collision term is added on the right-hand side of (5) and momentum exchange is assumed to be proportional to the velocity difference of colliding particles. The reason for adding this term in (10) here is that it is the easiest way of approximating nonideal MHD effects that may result from one or more of the other terms on the right-hand side of (10), in order to incorporate it in a fluid model. It satisfies the heuristic argument that microscopic wave-particle scattering generates anomalous resistivity and the expectation that such processes are most likely concentrated in regions of strong gradients, that is, in thin current sheets. A further argument that is often used is that nonideal processes in space plasmas tend to be highly localized, so that it does not matter too much for the large-scale structure and behavior which process actually operates in such regions. This argument appears to be most suitable for cases of driven evolution.

Other major microphysics assumptions concern primarily the properties of the pressure tensor, which is usually assumed to be isotropic

$$P_{\perp} = P_{\parallel} \quad (11)$$

or gyrotropic

$$\underline{P} = P_{\perp} \underline{I} + (P_{\parallel} - P_{\perp}) \frac{\mathbf{B}\mathbf{B}}{B^2} \quad (12)$$

When finite Larmor radius effects are important, deviations from the gyrotropic form (12) may occur which are often put in the form of gyroviscous effects. Such models are used in the context of the magnetopause, but not yet for studies of tail dynamics.

A closure of the dynamic fluid equations requires equations for the pressure tensor, which are derived from the next higher moment integrations of (5) using further approximations. These

approximations are typically ad hoc and rarely based on microphysics of the plasma configuration under consideration. However, in view of the uncertainties of how microphysics affects the macrophysics, it may be justified to explore these effects by using a wide range of closure relations. Among such relations for isotropic pressure p are an isothermal assumption

$$p/\rho = \text{const} \quad (13)$$

which requires a fast energy exchange with some thermal reservoir. In the other extreme of negligible heat conduction one can derive the adiabatic law

$$\frac{dp}{dt} + \gamma p \nabla \cdot \mathbf{v} = (\gamma - 1) \eta j^2 \quad (14)$$

where $\gamma = 5/3$ and the ohmic term has been included to be consistent with its inclusion in (10).

Using different values of γ in (14), other approximations can be considered; for instance, the incompressible limit can be obtained for $\gamma \rightarrow \infty$ and negligible ohmic heating. Alternatively, an opening of field lines at the ionosphere may also allow for incompressible modes [Erickson and Heinemann, 1992] or some other process may lead to fast energy exchange, while maintaining approximate pressure balance in contact with some pressure reservoir. In such cases an isobaric law may also be possible [Birn et al., 1994]:

$$\frac{dp}{dt} = 0 \quad (15)$$

Anisotropic models have been considered in the gyrotropic approximation [e.g., Hesse and Birn, 1992b], where one obtains, again including effects of finite resistivity

$$\frac{dP_{\parallel}}{dt} + P_{\parallel} \nabla \cdot \mathbf{v} + 2P_{\parallel} (\nabla \mathbf{v})_{\parallel} - 2\eta j_{\parallel}^2 = 2R \quad (16a)$$

$$\frac{dP_{\perp}}{dt} + 2P_{\perp} \nabla \cdot \mathbf{v} - P_{\perp} (\nabla \mathbf{v})_{\parallel} - \eta j_{\perp}^2 = -R \quad (16b)$$

with

$$(\nabla \mathbf{v})_{\parallel} = \frac{\mathbf{B}}{B} \cdot \nabla \mathbf{v} \cdot \frac{\mathbf{B}}{B}$$

Neglecting the heat flux leads to $R = 0$ in (16), that is, the double adiabatic generalization of (14). Recently, *Gary et al.* [1994a] have proposed a new, limited closure relation for the temperature anisotropy of ions in plasmas of moderate β (ratio of plasma and magnetic pressure). This relation was discovered by *Anderson et al.* [1994] in the magnetosheath and has been interpreted by *Gary et al.* [1994a] as arising due to pitch-angle scattering of hot magnetospheric protons by the electromagnetic proton cyclotron instability. In the magnetosphere *Gary et al.* [1994b] express this closure as an upper bound on the hot proton temperature anisotropy given by

$$\frac{T_{\perp h}}{T_{\parallel h}} - 1 = \frac{0.5}{\beta_{\parallel h}^{0.44}} \left(\frac{n_h}{n_e} \right)^{0.3} \quad (0.05 \leq n_h/n_e \leq 1.0) \quad (17)$$

where $\beta_{\parallel h} \equiv 2\mu_o p_{\parallel h}/B_o^2$ and the subscript h denotes the hot proton component of the magnetosphere. Because a fluid approach does not discern the difference between the hot and cool magnetospheric ions, and because the energy density of the former is typically substantially greater than that of the latter, we here use this equation in the form [*Birn et al.*, 1995b]

$$\frac{P_{\perp}}{P_{\parallel}} - 1 = \frac{S'}{\beta_{\parallel}^{0.5}} \quad (18)$$

where S' is taken as a constant with values somewhat less than 0.5 and the exponent 0.44 in (17) has been approximated by 0.5. If we assume that anisotropies that exceed the threshold given by (18) are reduced to the marginal limit on a fast time scale τ_{micro} , we can represent the quantity R in (16) by

$$R = \text{Max} \left(\frac{1}{3} \frac{P_{\perp} - P_{\parallel} - SB\sqrt{P_{\parallel}}}{\tau_{micro}}, 0 \right) \quad (19)$$

where $S = S'/\sqrt{2\mu_o}$.

In the following sections we will primarily focus on the effects of different approximations of the energy or pressure equations in unstable models of the magnetotail where a tearing type instability is triggered by the occurrence of anomalous resistivity, using both isotropic and anisotropic models. The initial state and the numerical simulation model will be discussed in Section 3, the effect of various values of γ in isotropic models in section 4, and anisotropic models with various forms of anisotropy reduction in Section 5.

3. Initial State and Numerical Model

As an initial state we use a realistic isotropic three-dimensional equilibrium model of the magnetotail developed by *Birn* [1987], which was used previously for dynamic simulation studies [e.g., *Birn and Hesse*, 1991, *Hesse and Birn*, 1992a,b]. It is illustrated in Figure 1 in the form of a magnetic flux surface emanating from an assumed circular cross section at its near-Earth end. The configuration includes flaring of the tail in y and z and an increase in plasma sheet thickness from midnight toward the flanks by a factor of about 2 and a somewhat weaker increase downtail. The length units in Figure 1 are scaled by the current sheet half-thickness L_z at $y = 0$ and the near-Earth end $x = 0$ of the tail section considered. Here and in the following we use dimensionless units, normalized at the near-Earth boundary by the initial current sheet half-thickness L_z , a lobe magnetic field strength B_L , a characteristic Alfvén speed v_A , defined by the characteristic lobe magnetic field strength and the plasma sheet density, and suitable combinations of these quantities. For illustration we use dimensional units $L_z \approx 12000\text{km}$, $B_L = 40\text{nT}$, and $v_A = 1000\text{km/s}$. This yields, for instance, a time scale $t_A = L_z/v_A = 12\text{s}$.

The boundary conditions consist of solid ideally conducting walls at each of the boundaries $x=0$, $y=y_{max}=10$, and $z=z_{max}=10$ and an open boundary at $x = -60$. All velocity components are set to zero except at the distant boundary in x , where a free outflow condition is assumed. Neumann boundary conditions ($\partial/\partial n = 0$) are imposed on density, pressure, and the tangential magnetic field components, while the normal magnetic field is held fixed. Symmetry conditions are used at $y = 0$ and $z = 0$. The code consists of an explicit finite difference scheme which solves the time-dependent resistive MHD equations; it has been described earlier in more detail [e.g., *Birn and Hesse*, 1991]. Only constant uniform resistivity is considered here, primarily, to avoid prescribing where reconnection and neutral line formation occurs.

4. Unstable Evolution of Isotropic Models

For all but the isobaric model, which will be discussed later, the tail configuration is stable within our simulation period of a few hundred Alfvén times (comparable to about 1–2 hours) when the ideal MHD constraint $\mathbf{E} + \mathbf{v} \times \mathbf{B} = 0$ is used and no external electric field is imposed. The unstable evolution of the magnetotail hence is initiated by imposing finite resistivity. The uniform value chosen here is $\eta = 0.005$, corresponding to a Lundquist number (magnetic Reynolds number) $R_m = 1/\eta = 200$. The occurrence of finite resistivity typically leads, after some initial diffusion, to the evolution of a generalized tearing instability [e.g., *Birn and Hesse, 1991*] with neutral line formation and plasmoid ejection into the far tail and dipolarization and wedge type current diversion in the near tail.

Details of the magnetic field and flow characteristics have been presented earlier [e.g., *Birn and Hesse, 1991*; *Hesse and Birn, 1992a,b*] and are not shown here. We focus on a comparison between the different plasma models demonstrated by the temporal evolution of different local or global parameters which are characteristic of the unstable evolution. Figure 2 shows a comparison between isotropic models with various values of γ , ranging from $2/3$ to 4 . The figure shows (from top to bottom) (a) the reconnected magnetic flux, obtained by integrating $|B_z|$ in the equatorial plane over the region where it has become negative, as an indication of the start and rate of reconnection, (b) the maximum tailward speed v_x , (c) the total field-aligned current of region 1 signature, integrated at the near-Earth boundary $x = 0$, which is part of the substorm current wedge, and (d) the magnetic field component B_z at $x=-5.6$, $y=0$, $z=0$ earthward from the reconnection site, as a representation of the dipolarization of the field in the near tail. The figures are taken from the paper by *Hesse and Birn [1992a]*.

All quantities shown in Figure 2 behave quite similarly. There is, however, a tendency for near-Earth effects, dipolarization and substorm current wedge signatures, to be stronger for smaller values of γ . This effect becomes extreme when the limit $\gamma \rightarrow 0$ or the isobaric case is considered. As shown by *Birn et al. [1994]* through linear stability analysis and through two-dimensional

and three-dimensional numerical simulations [Birn *et al.*, 1995], this case becomes unstable even in the absence of resistivity. In that case the nonlinear evolution leads to the rapid formation of a thin current sheet in the near tail but also some dipolarization effects further earthward. This is demonstrated by Figure 3 taken from Birn *et al.* [1995]. The figure shows (from top to bottom) the magnetic field component B_z , the current density j_y , both along the x axis, and j_y as a function of z at $x=-5.6$, $y=0$ for various times indicated in the figure. Note that the changes in B_z and j_y occur within about 5 to 10 Alfvén times, corresponding to about 1 to 2 min. In the absence of resistivity the evolution does not lead to reconnection. However, when a small resistivity $\eta = 0.005$ is added, reconnection and plasmoid formation is initiated on the same fast time scale.

5. Effects of Anisotropy and Isotropization

The effects of anisotropy and its reduction through microscale processes are studied within four dynamic models: (A) an isotropic model with an adiabatic index $\gamma = 5/3$, as discussed in section 4, (B) an anisotropic model with isotropizing terms as in (19) but $S = 0$, that is, reducing any anisotropy, (C) an anisotropic model described by (19), which reduces anisotropies to the marginal proton cyclotron limit if the instability criterion is satisfied, and (D) an anisotropic model without any isotropization, $R = 0$ in (16), representing the double adiabatic approximation modified by ohmic heating. Again we use constant resistivity $\eta = 0.005$ and the initial state discussed in Section 3 for our three-dimensional simulations.

Figure 4 shows the evolution of characteristic parameters in the four models, (from top to bottom) the total reconnected flux, the maximum tailward speed, the maximum current density j_y , and the integrated region 1 type field aligned current at the near-earth end $x = 0$. All quantities in Figure 4 demonstrate clearly that the isotropic case is the most unstable while the double adiabatic case is the most stable; in fact it does not even lead to reconnection within the simulation period. (Note, however, that a similar two-dimensional run with a resistive double adiabatic model led to instability, although at a significant delay.) Consistent with our earlier

results [Hesse and Birn, 1992b] we find that the closer the model is to isotropy, the more unstable it is. Isotropization, even if it does not lead to full isotropy, has a significant destabilizing effect.

In addition, we find that the anisotropy model can also have a significant effect on the structure of the mode and, most significantly, on the location of neutral line formation. This is demonstrated by Figure 5, which shows the variation of B_z along the x axis for models B (which is close to the fully isotropic case) and C (which reduces anisotropies only to the marginal instability limit). Neutral line formation in model C not only occurs much later than in model B but also much further tailward.

Typical anisotropies are demonstrated for model C in Figure 6, which shows the variation of parallel and perpendicular temperatures with z at two locations in x at $y = 0$. For an indication of the thickness and extent of the plasma/current sheet, the variation of the cross-tail current density j_y is also shown. The location $x = -7.5$ represents a region of plasma sheet thickening and dipolarization, whereas the location $x = -30$ is close to the x-type neutral line where the plasma/current sheet has thinned considerably. At both locations (and others as well) the central part of the plasma sheet has been heated (the initial temperature in the chosen units was uniformly 0.5) but remained approximately isotropic. Large mirror-type anisotropies ($p_\perp > p_\parallel$) have developed in the boundary regions of the plasma/current sheet, extending into the lobes. A comparison with model D (not shown here) and with the marginal stability criterion (18) shows that these anisotropies are similar to those in the fully anisotropic model, but reduced in magnitude, primarily in the boundary regions, consistent with the marginal stability limit.

6. Summary and Discussion

Using isotropic and anisotropic resistive magnetofluid simulations, we have investigated the influence of microphysical processes on the large-scale dynamic evolution of the magnetotail. Obviously, the occurrence of anomalous resistivity is one of the most drastic possible consequences of microscopic processes, as it changes the stability of the tail and enables magnetic reconnection. Although observations have provided ample evidence that flux transfer from closed to open field

line regions and vice versa and plasmoid formation and release are an important part of magnetospheric activity, implying the occurrence of reconnection, the process which generates resistivity or an equivalent deviation from ideal MHD has not been clearly identified, despite more than 20 years of effort.

In contrast to the effects of anomalous resistivity or other possible deviations from ideal MHD, the role of microphysics elsewhere within a fluid plasma model of the magnetotail has received much less attention. In this paper we have therefore concentrated on several models of energy transport and pressure closure relations and their effects on the dynamic evolution of the magnetotail.

Isotropic models were varied by the value of the ratio of specific heats γ , which is equivalent to the polytropic index in the absence of nonadiabatic (here, ohmic) heating. Large values of γ represent nearly incompressible conditions, $\gamma = 5/3$ the adiabatic case, valid for an isotropic plasma in the absence of heat flux, $\gamma = 1$ the isothermal case, and $\gamma \rightarrow 0$ or $dp/dt = 0$ the extreme, isobaric case, which requires fast energy and/or mass exchange with some external reservoir while maintaining approximate pressure balance. All unstable resistive cases are characterized by the growth of a generalized tearing mode with neutral line formation and plasmoid ejection into the far tail, while the near-Earth region exhibits dipolarization of the magnetic field and the current reduction and diversion associated with the substorm current wedge [McPherron *et al.*, 1973].

A variation of the value of γ in the plausible range from about 1 to 4 has no drastic effect on the evolution of the tail. Reconnection proceeds similarly, and dipolarization and the generation of field-aligned currents associated with the substorm current wedge is qualitatively similar for different values of γ . There is, however, a tendency for the latter, near-tail effects to be more pronounced for smaller values of γ . In the extreme case of isobaric thermodynamic conditions, $dp/dt = 0$, the tail dynamics is drastically altered: the tail becomes unstable even in the absence of resistivity, leading to the rapid formation of a thin current sheet, and the dynamic time scale is much shorter [Birn *et al.*, 1994, 1995a].

Anisotropic simulations were based on three different plasma models, the double adiabatic model in which perpendicular and parallel pressure components are uncoupled, a model based on the marginal ion cyclotron instability limit, in which anisotropies that exceed this limit are reduced on a short time scale [Gary *et al.*, 1994a,b; Birn *et al.*, 1995b], and a model in which any anisotropy is reduced on such short time scale. Consistent with earlier results by Hesse and Birn [1992b], we find that the more isotropic models tend to be more unstable: reconnection starts sooner, the reconnection rate is larger, and region-1 type field-aligned currents associated with the substorm current wedge increase faster when the model is closer to isotropy. As in the driven nonresistive models [Birn *et al.*, 1995b] there is a strong influence of the anisotropy reduction model on the spatial variation of B_z in the neutral sheet $z = 0$. As a consequence the reconnection site and the location where a magnetic neutral line forms can vary significantly from model to model.

The effects of the anisotropy reduction on the stability of the tail in the resistive dynamic simulations are consistent with effects found in driven nonresistive simulations [Birn *et al.*, 1995b]. In these simulations a thin current sheet was formed within the near-tail plasma sheet combined with a reduction of B_z in the neutral sheet. Both the current density enhancement and the reduction of B_z were stronger the closer the model was to isotropy. This means that a more isotropic model will more likely, or sooner, lead to instability than a strongly anisotropic model (such as a double adiabatic model) when the instability threshold depends on the magnitude of the current density [e.g., Lui *et al.*, 1990] or on the magnitude of B_z in the neutral sheet [e.g., Schindler, 1974].

The anisotropic dynamic magnetotail models develop mirror-type anisotropies ($p_{\perp} > p_{\parallel}$) in the (low beta or low current density) boundary regions of the plasma sheet and the adjacent lobe regions, while the (high beta or high current density) center region of the plasma/current sheet remains close to isotropy. The type of anisotropy and its relation with the plasma beta are consistent with observations of Lui *et al.* [1992] after substorm onset.

Our results on the effects of anisotropy appear to differ from earlier studies of the collisionless tearing instability with anisotropic pressure which have been carried out using both analytic techniques [e.g., *Chen and Palmadesso*, 1984] and computer simulations [e.g., *Ambrosiano et al.*, 1986; *Pritchett and Coroniti*, 1990]. These investigations have shown that a mirror-type anisotropy present in the initial configuration or caused by quiet-time convection can significantly enhance the growth of this mode. However, these results cannot be compared directly to the conclusions drawn here and in *Hesse and Birn* [1992b] because the models are quite different. Our work has addressed the resistive tearing mode based on an isotropic equilibrium model; the reduction of instability growth which we demonstrate is a consequence of anisotropies which develop as the system evolves.

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Figure Captions

Fig. 1. Initial configuration of the simulations, showing a magnetic flux surface emanating from a circle of 10 units radius at $x = 0$ (which corresponds to a distance of $10\text{--}20R_E$ from the Earth). The dashed lines indicate the plasma sheet (current sheet) width. Length units are scaled with the current sheet half-width ($\approx 2R_E$) at $x = 0, y = 0$.

Fig. 2. Temporal evolution of characteristic parameters for the isotropic resistive simulations, showing from top to bottom: (1) the total reconnected flux, calculated from integrating B_z at $z = 0$ over the area where $B_z < 0$, (2) the maximum tailward speed, (3) the total field-aligned current of region 1 signature integrated at $x = 0$, and (4) the magnitude of B_z at $x = -5.6, y = 0, z = 0$ as an indicator of dipolarization in the near tail. The different values of γ , the ratio of specific heats, are indicated in the figures.

Fig. 3. Spatial and temporal variation of characteristic parameters for an isobaric nonresistive simulation, showing from top to bottom: (1) B_z as a function of x at $y = 0, z = 0$, (2) j_y as a function of x at $y = 0, z = 0$, and (3) j_y as a function of z at $x = -5.6, y = 0$. Different times are indicated in the figures.

Fig. 4. Temporal evolution of characteristic parameters for the nondriven resistive simulations, showing from top to bottom: (1) the total reconnected flux, calculated from integrating B_z at $z = 0$ over the area where $B_z < 0$, (2) the maximum tailward speed, (3) to maximum cross-tail current density, and (4) the total field-aligned current of region 1 signature integrated at $x = 0$. The four plasma models are indicated in the figures.

Fig. 5. Variation of B_z as a function of x at $y = 0, z = 0$, and different times as indicated for two of the four anisotropic resistive simulations: plasma model B (top), which is close to isotropy, and model C (bottom), based on the marginal proton-cyclotron instability criterion.

Fig. 6. Variation of the parallel and perpendicular temperatures with z at $x = -7.5$ (top) and $x = -30$ (bottom), $y = 0, t = 300$ for the anisotropic plasma model C, based on the marginal proton-cyclotron instability criterion. The current density J_y (dotted line) is shown also to indicate the extent of the plasma/current sheet.