A TWO-DIMENSIONAL PARTICLE SIMULATION OF THE MAGNETOPAUSE CURRENT LAYER

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Abstract

We have developed a $2^{1/2}$ -D (x, y, v_x, v_y, v_z) electromagnetic code to study the formation and the stability of the magnetopause current layer. This code computes the trajectories of ion and electron particles in their self-consistently generated electromagnetic field and an externally imposed 2-D vacuum dipolar magnetic field. The results presented here are obtained for the simulation of the solar wind-magnetosphere interaction in the subsolar region of the equatorial plane.

We observe the self-consistent establishment of a current layer resulting from both diamagnetic drift and $\mathbf{E} \times \mathbf{B}$ drift due to the charge separation. The simulation results show that during the establishment of the current layer, its thickness is of the order of the hybrid gyroradius $\rho_H = \sqrt{\rho_i \rho_e}$ predicted by the Ferraro-Rosenbluth model. However, diagnostics indicate that the current sheet is subject to an instability which broadens the width of the current layer. Ripples with amplitudes of the order of the ion gyroradius appear at the interface between the field and the particles. These perturbations are observed both on the electrostatic field and on the compressional component of the magnetic field. This instability has a frequency of the order of the local ion cyclotron frequency. However, the modulation propagates in the same direction as the electron diamagnetic drift which indicates that the instability is not a classical gradient-driven instability, such as the lower hybrid or ion drift cyclotron instabilities. The nonlinear phase of the instability is characterized by the filamentation of the current layer which causes anomalous diffusion inside the central current sheet.

Introduction

Recent controversies on the mechanism of the generation of flux transfer events (FTEs) has revived interest in the microstructure of the earth dayside magnetopause and shown the need for a new theoretical effort to study its stability. We present here the results of $2\frac{1}{2}$ -dimensional electromagnetic particle simulations which relate to the formation and stability of the current layer.

Historically, the theoretical investigation of the solar wind interaction with the earth's magnetic field started with Chapman and Ferraro (1931). Assuming a perfectly conducting and unmagnetized solar stream, they showed the confinement of the dipole field by currents flowing in the front surface of the stream. This is the current layer that now we call the magnetopause. Although Hoyle (1949) and Alfvén (1950) have already suggested the presence of an interplanetary magnetic field (IMF), early works only considered the problem of the equilibrium between an unmagnetized plasma and a vacuum region of constant magnetic field (Ferraro, 1952; Rosenbluth, 1957; Dungey, 1958, 1961; Grad, 1961; Longmire, 1963; Sestero, 1965; see also reviews by Willis, 1971, 1975).

With the start of spatial exploration, the first dayside magnetic field profiles (Cahill and Amazeen, 1963) confirmed the presence of a magnetized sheath between the shock and the magnetopause, constituted by a highly turbulent subsonic and super-Alfvénic plasma flow. Measurements also showed nonvanishing particle densities in a layer adjacent to the magnetospheric side of the discontinuity. These observations led to a second class of microscopic models of the magnetopause where the current sheet separates two magnetized plasmas. These magnetopause models are usually based on analytical solutions of Vlasov-Maxwell equations and differ mainly by the choice of initial velocity distribution functions, number of components, and boundary conditions (see review by Lee and Kan, 1979). There is no magnetic field component normal to the magnetopause surface. The boundary is impermeable to the solar wind flow. This type of magnetopause structure corresponds to the tangential discontinuity of the magnetohydrodynamics (MHD) description (Landau and Lifshitz, 1960).

A radically different type of magnetopause structure was introduced by Dungey's (1961) new developments on particle acceleration at magnetic neutral points (Giovannelli, 1946). Challenging Chapman and Ferraro's idea of a strict confinement of the geomagnetic field, his model suggests that a small fraction of the dayside earth's magnetic field lines can be connected to the interplanetary medium during periods of

southward IMF. As a consequence, instead of being a closed boundary, the magnetopause becomes open to the transfer of solar wind mass and momentum. Reconnection models (Petschek, 1964; Levy et al., 1964; Sonnerup, 1970; Yeh and Axford, 1970; Vasylunas, 1975) predict the existence of a magnetic component normal to the surface of the discontinuity and high-speed flows. The dayside magnetopause is then of the rotational discontinuity type (Su and Sonnerup, 1968; Hudson, 1971; Sonnerup and Ledley, 1974; Lee and Kan, 1982; Swift and Lee, 1983). This steady-state MHD flow model has provided the basis of the magnetopause paradigm that most of both experimenters and theoreticians have been pursuing during the past decade.

From the experimental point of view, the past twenty years of space exploration have permitted the accumulation of enough measurements to calculate typical plasma parameters at the dayside magnetopause (see reviews by Haerendel and Paschmann, 1982; Paschmann, 1984a; Elphic 1987). Magnetosheath plasma flow speeds (v_f) vary from one to a few hundred km/s. These velocities are small compared to the sound speed, $V_S \approx 200$ -400 km/s, but remain larger than the Alfvén velocity (V_A) which is of the order of 100 km/s for average density and magnetic field ($n_{MS} = 10 \text{ cm}^{-3}$, $B_{MS} = 15 \text{ nT}$). The temperature (T_i) of the bulk of the ion distribution varies between 5 and $10 \cdot 10^6$ K, so the magnetosheath ion gyroradius (ρ_i) is of the order of 100 km. The ratio of the ion to electron temperatures (T_i/T_e) is around 10 and the temperature anisotropy ($T_{\perp}/T_{\parallel}$) fluctuates between 1 and 2 in front of the magnetopause.

More recently, the ISEE multispacecraft missions allowed substantial progress in the investigation of magnetospheric boundaries, especially the magnetopause. Indeed, for the first time, most of the instruments were able to provide enough adequate time resolution measurements to study magnetopause phenomena, and simultaneous in-situ measurements at different locations were used to measure macroscopic features of the magnetopause boundary, such as the thickness (400 to 1600 km) and the velocity (10 to 80 km/s) of the current layer (Berchem and Russell, 1982). Of course, a great deal of effort has been devoted to analyzing correlative plasma and field measurements in order to identify dayside magnetic reconnection. The observations of some of the predicted signatures have to some extent supported the steady-state model (see reviews by Paschmann, 1978, 1984b; Russell, 1984 a,b; Sonnerup, 1984).

A new and important magnetopause feature emerged with the identification of temporally unsteady and spatially very localized reconnection processes, observed both in the polar cusps (Haerendel et al., 1978) and in the equatorial region (Russell and Elphic, 1978, 1979). These observations, which have been called flux transfer events (FTEs), have considerably altered the theoretical approach of the study of the microstructure and stability of the magnetopause. Indeed, although the equilibrium of the magnetopause was questioned in the late 1960's, those studies only focused on the

stability of the magnetospheric flank and tail boundaries. Eviatar and Wolf (1968) using quasilinear theory, suggested that the two-stream ion-cyclotron instability could scatter ions at a distance comparable to their gyroradius within a few gyroperiods. They calculated that for a 100-km-thick boundary, the transfer of momentum would be adequate to drive the magnetospheric convection postulated by Axford and Hines (1961). Attention was also drawn by the problematical confinement of the geomagnetic tail (Lerche, 1967; Parker, 1967a,b). Nevertheless, Alpers (1971), Su and Sonnerup (1971) demonstrated that an equilibrium does exist even in a region where the solar wind particles have a bulk velocity tangential to the geomagnetic field.

Before the FTEs became closely associated with the stability of the dayside magnetopause, Gary and Eastman (1979) suggested that the lower hybrid drift instability (Davidson and Krall, 1977) could account for the bursty low frequency ($<100\text{Hz}$) electric and magnetic fluctuations observed by Gurnett et al. (1979). However, they also noted that the mode is probably not excited by thick boundary layers, but should be driven only by smaller scale density gradients. The possible occurrence of another instability, the tearing mode (Furth et al., 1963), was also suggested from both laboratory terrella experiments (Dubinin et al., 1980) and space observations at the Earth and Jupiter's magnetopauses (Greenly and Sonnerup, 1981). In response to these observations, linear theory of the guide field tearing instability (Laval et al., 1966; Coppi et al., 1966; Galeev and Zeleny, 1978) was derived for high- β plasmas (Quest and Coroniti, 1981b) in order to study the tearing at the dayside magnetopause (Quest and Coroniti, 1981a). It was shown that the linear growth rate of the instability is a sensitive function of the IMF orientation, magnetopause thickness, and electron number density. Quest and Coroniti (1981 a,b) also emphasized that tearing should involve a two-dimensional wave spectrum ($k_y \neq 0$) tangential to the magnetopause and that tearing-reconnection should be spatially inhomogeneous and temporally unsteady. Thus, it is natural that the tearing instability has been proposed to explain FTE generation (Lee and Fu, 1985; Kan, 1988; La Belle-Hamer et al., 1988). Nevertheless, since the initial Russell and Elphic's (1978) FTE model of intermittent reconnection of individual flux tubes, many other mechanisms have also been advanced. They are based on stochastic percolation of the magnetic field through the magnetically destroyed surfaces (Galeev et al., 1986) or sporadic multiple reconnection processes (Ding et al., 1986; Sato et al., 1986; Shi et al., 1988; Sholer, 1988).

In this paper we examine the fundamental interaction between the subsonic magnetosheath flow and the magnetospheric field. We think that it is necessary to achieve a better understanding of the basic physical process before pursuing more sophisticated magnetopause models by including other fundamental ingredients such as

the magnetosheath field, the presence of energetic particles of magnetospheric origin, or the coupling with the ionosphere. We thus re-examine the classical Ferraro-Rosenbluth problem of the equilibrium between a plasma and a vacuum magnetic field. In fact, the kinetic treatment of this interaction has not been addressed in a self-consistent manner since the one-dimensional study by Hasegawa and Okuda (1967). Most of the past studies of the tangential discontinuity structure of the magnetopause assume an ad hoc initial Harris (1962) type of equilibrium. On the other hand, recent nonlinear numerical models, such as the ones proposed for the generation of FTEs, use MHD or hybrid codes and thus ignore electron Landau damping and electrostatic effects. However, as several studies (Quest and Coroniti, 1981b; Lembège and Pellat, 1982; Swift, 1986) have demonstrated, the electron dynamics is primordial in the physics of current sheets.

To study the formation and the stability of the subsolar magnetopause current layer, we have thus developed a two-dimensional simulation code which computes the trajectories of both ion and electron particles in their self-consistently generated electromagnetic fields and an externally imposed magnetic field. Our model consists of an unmagnetized subsonic super-Alfvénic plasma flow injected transversely to a two-dimensional vacuum dipolar magnetic field.

Simulation model and computational method

Figure 1 shows the geometry of our simulation system. The $L_x \times L_y$ domain represents the subsolar region of the equatorial plane in which we follow the trajectories of solar wind particles. The ion and electron particles are injected from the left, so the right boundary delineates the magnetospheric side of our system. Since for numerical convenience we fix the origin ($x=0$) of our system of coordinates at the injection boundary, our reference frame is oriented oppositely to conventional geocentric systems of coordinates. The x axis points towards the Earth, the y axis is directed from dusk to dawn, and the z axis points northward.

To model the Earth's magnetic field, we use a two-dimensional magnetic dipolar field generated by two infinite parallel wires aligned with the y direction. We can then assume that all the system is uniform and periodic along the y axis. The non-zero components of the field are given by:

$$B_x = 2B_0 \frac{-Xz}{(X^2+z^2)^2}$$

$$B_z = B_0 \frac{X^2-z^2}{(X^2+z^2)^2}$$

where $X = x_0 - x$, and x_0 is the location of the dipole axis. However, in order to keep a significant B_z variation within the limits of our system without imposing a too strong field on the left boundary, we have to position the axis of the dipole relatively close to the left boundary ($x_0 = 4L_x$). In return such a model leads to an excessive field line curvature in the immediate vicinity of the simulation plane. In reality, in the near equatorial region of the subsolar magnetopause, the radius of the field curvature is much larger than the thickness of the current layer and $B_z \gg B_x$. To study the limit of a small curvature of the field line, we assume that the field is independent of z ($B_x = 0$; $B_z = B_0/X^2$). The initial magnetic field strength at $z=0$ is such that before compression, the ratio of the average electron gyrofrequency to the electron plasma frequency (ω_{ce}/ω_{pe}) is equal to 0.5 at $x=0$ and to 3.125 at $x=L_x$. Furthermore, at the start of the simulation, B_0 is ramped linearly over the first 50 electron gyroperiods and then held constant. This "quiet start" technique allows us to apply smoothly and self-consistently the boundary conditions that we describe later on. We introduce a finite computational grid on the x - y plane, with a grid size Δ , such that $L_x = 200\Delta$ and $L_y = 128\Delta$, and define a dimensionless system of units by choosing Δ and ω_{pe} to be the spatial and time units of our system. We define c the speed of light such that the collisionless skin depth c/ω_{pe} is equal to 5Δ .

The calculation is then initialized by loading a solar wind particle reservoir adjacent to the left boundary of our simulation system. We use finite-size ion and electron particles to reduce the collision frequency and we distribute them uniformly to obtain a number density (n_0) equal to 4 per each Δ^2 grid cell of the reservoir. We also initialize the particle components of velocity v_x , v_y , and v_z from a Maxwellian velocity distribution function generator such that the electron thermal velocity v_{Te} is equal to $\Delta\omega_{pe}$. The grid size Δ then becomes identical to the Debye length (λ_D) of our system. Note that since we assume no spatial change along the z direction and define three velocity components, our model is $2^{1/2}$ -dimensional. Ion to electron mass (M_i/m_e) and temperature (T_i/T_e) ratios used are equal to 100 and 9, respectively. After initialization, a constant drift velocity $v_{dx} = 0.25\Delta\omega_{pe}$ is then added to the x -component of the solar wind particles to simulate the subsonic, super-Alfvénic magnetosheath flow. The sound speed is slightly greater than the ion thermal speed (v_{Ti}) which is equal to $0.3\Delta/\omega_{pe}$. The initial Alfvén speed before compression ($V_A = c(\omega_{ce}/\omega_{pe})(m_e/M_i)^{1/2}$) varies from $0.25\Delta\omega_{pe}$ at $x=0$ to $1.5\Delta\omega_{pe}$ at $x=L_x$.

We then start the iterative loop of the simulation by calculating the charge ρ and the current densities J produced by all particles entered in the system by using the linear interpolation scheme (Birdsall and Fuss, 1969; Morse and Nielson, 1969):

$$\begin{aligned}\rho(x,y) &= \sum q_j S_j(x,y) \\ J(x,y) &= \sum q_j v_j S_j(x,y)\end{aligned}$$

where $S_j(x,y) = \frac{1}{2\pi\Delta^2} e^{-\frac{[(x-x_j)^2 + (y-y_j)^2]}{2\Delta^2}}$ is a Gaussian-shaped factor for the finite-size particles used to reduce the collision frequency. Once J is calculated, we break down the transverse (J_t) and longitudinal (J_l) parts of the current density by assuming that the longitudinal current J_l is derived from a potential ($J_l = -\nabla\psi$), and proceed to the integration of the time-dependent Maxwell's equations:

$$\begin{aligned}\nabla \times E_t &= -\frac{1}{c} \frac{\partial B}{\partial t} \\ \nabla \times B &= \frac{1}{c} \frac{\partial E_t}{\partial t} + \frac{4\pi}{c} J_t\end{aligned}$$

To be consistent with our assumption of an unmagnetized plasma flow in the magnetosheath, we impose the condition that the B_y , B_z and E_y , E_z components vanish on the left and right boundaries, respectively, and proceed to the integration in time of these equations by using a second order accurate time-centered method (Langdon et al., 1976). We complete the set of electromagnetic fields by calculating the electrostatic field which is purely longitudinal:

$$\nabla E_z = 4\pi\rho.$$

We solve that equation by assuming no boundary perturbations, so the electrostatic potential ϕ vanishes ($E_z = -\nabla\phi$) on the left and right boundaries. After a Fourier transform of all the quantities along the y direction, the corresponding Poisson's second order differential equation is given by:

$$\frac{d^2\phi(x,k_y)}{dx^2} - k_y^2\phi(x,k_y) = 4\pi\rho(x,k_y),$$

which is easily solved by a triangularization method.

Finally, once the total $E(x,y)$ and $B(x,y)$ fields are completely determined, we use exact dynamics to advance the particles by solving the equations of motion:

$$m_j \frac{dv_j}{dt} = q_j \left(E_{tot} + \frac{1}{c} v_j \times B_{tot} \right).$$

using Buneman's (1966) scheme.

The code has been tested for numerical dissipation by loading the entire system of a magnetospheric plasma without injecting the solar wind. A run over $500/\omega_{pe}$ shows that the total energy is conserved within a few percent.

Results

Figure 2 shows a composite display of a few diagnostics of our simulation at time $t=140/\omega_{pe}$. On the left-hand column, we plotted as a function of the x position, the electron number density n_e (a), the E_x component of the sum of the ambipolar electrostatic field and the electric perturbation (b), the current density J_y (c), and the B_z component of the magnetic field (d). Here and after, all the one-dimensional plots that we show have been obtained by averaging each quantity both in time in order to eliminate high frequency noise, and spatially over the y -direction to show only the $k_y=0$ mode. On the right-hand column of the same figure, we display instantaneous two-dimensional scatter plots of the electrons and ions which show their locations [(e) and (f) respectively] in the x - y plane and their phase space distributions [(g) and (h) respectively] in the velocity phase space x - v_x . These diagnostics show the self-establishment of a current layer between the magnetosheath plasma and the magnetospheric field. The longitudinal electric field E_x (b), mainly due to the charge separation between ions and electrons, generates a negative ($y<0$) transverse electron current by a $cE \times B/B^2$ drift. An additional but smaller current is generated by the diamagnetic drift $J_p = -\nabla p \times B/B^2$ due to the strong density gradient seen in Fig. 2(a). Notice that the nonzero electric field and the density depletion observed at the injection boundary are just local effects due to the electron gyration at the edge of our simulation box and do not affect the rest of the system. Both $E \times B$ and diamagnetic currents flow in the same direction and contribute to the formation of the sharp current layer, J_y shown in Fig. 2(c), which compresses the magnetic field B_z (d) and produces the magnetic field gradient characteristic of the magnetopause. Nevertheless, as the ion velocity distribution in Fig. 2(h) shows, only a few ions have a negative v_x velocity. This means that the bulk of the particles have not reached the point where the particle pressure balances the magnetic pressure. Although a few particles have already penetrated deeply in the region of higher magnetic field, other diagnostics indicate that these ions also have negative v_y suprathermal velocities and thus have larger ion gyroradii than the bulk of the particles. Note that the electron v_x velocity distribution shown in Fig. 2(g) is already symmetrical because of the trapping by the local magnetic field.

As time progresses, the beam advances in the system, pushing and compressing the magnetic field, until ions get reflected. Naturally, the location of the reflection should be determined by the pressure balance condition. Since we suppose that the flow has already been shocked before the injection, when writing the balance condition, both terms for the dynamic and the thermal pressures have to be retained:

thus $2nM_i v_{dx}^2 + KT_i \approx B_z^2 / 8\pi$, where K is the usual Boltzmann's constant. The fact that the beam velocity $v_{dx} = 0.25\Delta\omega_{pe}$ is very close to the ion thermal speed $v_{Ti} = 0.3\Delta\omega_{pe}$ has further implications: we can use this to simplify the previous balance equation and rewrite it in terms of the ratio ω_{ce}/ω_{pe} . Firstly, only about half of the particles penetrate into the system. This is readily seen in Fig. 2(a) which shows that the beam has a number density equal to 2 instead of the initial 4 in the particle reservoir. Since we use the full Maxwellian velocity distribution functions to define the simulation plasma parameters, we can eliminate from the plasma pressure terms the factor 2 which takes into account the jump in density due to the specular reflection of the particles. Secondly, since both velocities are close, we can estimate that both dynamic and thermal pressure contributions are roughly equal, so in first approximation we have $2nM_i v_{dx}^2 \approx B_z^2 / 8\pi$. This allows us to rewrite the pressure balance condition as $\omega_{ce}/\omega_{pe} = 2\sqrt{M_i/m_e} v_{dx}/c$. With our set of initial values, this relation implies that the pressure balance is satisfied where $\omega_{ce}/\omega_{pe} \approx 1$. Taking into account our normalization factors, this ratio corresponds to $B_z/B_0 \approx 5$, so we expect the location of the reflexion point to be around half of the system length $x_p \approx 100\Delta$.

Figure 3 displays the same diagnostics as in Fig. 2, but at a later time, $t = 440/\omega_{pe}$. As expected, we now observe the start of the reflection of the bulk of the particles. This is seen on Fig. 3(a), which shows that the beam density increases significantly to almost 3 particles per cell, and in Fig. 3(h), where it is readily seen that the ion v_x velocity distribution is almost symmetrical. We also note that the current layer (3-d) advances to $x = 80\Delta$, sweeping the magnetic field. It turns out that the reflection point is closer to the left than we expected. This is due to the compression of the field. Nevertheless, the actual ratio B_z/B_0 agrees with the prediction of the pressure balance condition. Another interesting feature is seen in the current layer profile in Fig. 3(c). Although its peak amplitude is slightly smaller, the width of the current layer is substantially larger than the one observed at $t = 140/\omega_{pe}$. We will comment on the structure inside the layer later on. Examining the two-dimensional scatter plots of the particle x - y locations, Figs. 3(e) and 3(f), we notice that this broadening corresponds to the development of an instability at the interface of the plasma and the field. Naturally, the modulation appears clearer for the electrons (e) than for the ions (f) because of their smaller gyroradii. Before discussing the results of the spectral analysis of this instability, let us examine the time history of the current sheet formation and the development of the instability.

To follow the dynamics of the formation of the current layer we have compiled in Figs. 4(a) and 4(b) a series of diagnostics taken from time $t = 100/\omega_{pe}$ to time $t = 800/\omega_{pe}$. The time delay between each series is $\delta t = 100/\omega_{pe}$. We have organized the results in four columns. From left to right, we display two-dimensional plots of the

electron x-y locations, and one-dimensional plots of the electron density n_e , the current density J_y , and the magnetic field B_z , versus the x location. Inspection of the scatter plots of the electron x-y locations (left column) shows a phase shift of the instability and indicates that we can distinguish 3 stages in the evolution of the interaction between the plasma and the magnetic field: firstly, the establishment of the current layer, secondly, the linear growth of the instability, and finally, the nonlinear development.

From the start of the simulation to around time $t=200/\omega_{pe}$, particles are pushing the field at a speed close to the injection velocity. As we see in Fig. 2, only a few particles are already reflected back as the slight density increase indicates. This phase corresponds to the self-establishment of the current layer. During this first stage the current layer has an almost constant thickness of about 17Δ , which is much smaller than the approximately 40Δ ion gyroradius measured at the peak of the current density. We define the thickness of the current layer by measuring the width of the peak of the J_y component at one-half of its value. To compare the observed thickness to the classical results of the Ferraro-Rosenbluth model, we have to calculate the average hybrid gyroradius $\rho_H = \sqrt{\rho_i \rho_e}$, which becomes $\rho_H = \rho_i \sqrt{30}$, using thermal velocities to calculate both ion and electron gyroradii. Since at one-half of the peak the observed ρ_i varies between 140Δ to 30Δ , we obtain an average hybrid gyroradius of 15.5Δ , which is very close to the thickness observed.

At the time $t=300/\omega_{pe}$, the bump observed in the density profile indicates that more particles are reflected. However, as we know from Fig. 3, the bulk of the beam has not yet penetrated far enough for the particle pressure to balance the magnetic field pressure, so the beam is still progressing but has slowed down to almost half of the injection velocity. The current layer thickness is now at least of the order of one ion gyroradius. We also notice the presence of very small ripples on the scatter plots of the electrons (left column), which indicates the growth of an instability. This marks the beginning of the second stage which lasts until time $t=600/\omega_{pe}$. This period of time corresponds to the linear phase of the instability. Note that since the instability started growing before the bulk of incoming ions start reflecting, its frequency and growth rate should be then comparable to the ion gyrofrequency. This will be discussed in a more quantitative manner when reporting the spectral analysis of the instability. Nevertheless, we can note that the maximum amplitude from the crest to the trough of the modulation is of the order of one ion gyroradius, so the layer thickens substantially to almost two ion gyroradii.

Finally, from time $600/\omega_{pe}$ to the end of the simulation [Fig. 4(b)], we observe the development of the quasilinear and highly nonlinear phases of the instability. The plots of the electron locations show first the steepening of the modulation, and then,

starting from time $t=700/\omega_{pe}$, the formation of eddies. Looking at the current density profiles J_y/J_0 , it is also apparent that the layer has considerably spread out while the initial peak has been reduced to half of its value observed during the linear phase of the instability. This is due to the progressive pinch of the central current sheet which ultimately leads to the filamentation of the current layer at time $t=800/\omega_{pe}$. Examining Fig. 4(b) plots of the electron density profiles, we observe the formation of a small plateau which gradually evolves as the instability becomes highly nonlinear. This change in the density profile is very similar to the one associated to the nonlinear evolution of the collisionless drift-wave instability. Lee and Okuda (1976) showed that as the instability develops, the waves grow in amplitude, causing quasilinear diffusion in density. The change in the density profile then alters the local dispersions and becomes the dominant stabilization mechanism of the instability. Unfortunately for us, this anomalous transport coincides with the end of the simulation when we run out of particles in the reservoir, so it is difficult to verify that the diffusion saturates the instability. However, the resulting local density plateau turns off the diamagnetic currents which cause the local pinch observed in the current layer profiles.

To characterize the instability, we have carried out several Fourier analyses of the electromagnetic and electrostatic fluctuations. For various x -locations, we analyzed the first five modes along the y direction of the system. Naturally, the strongest fluctuation is observed on the B_z component for $k_y=0$, because it corresponds to the compression of the dipole field. Figure 5 shows the analysis of the magnetic perturbation δB_z measured into the reflection region at location $x=80\Delta$. The top panel displays time series of $\delta B_z/B_0$, where B_0 is a normalization factor. One can see that $\delta B_z/B_0$ first reaches a maximum around time $t=300/\omega_{pe}$ and then decreases to become negative around time $t=600/\omega_{pe}$. This variation is simply due to the advance of the current layer and appears as a peak at zero frequency in the power spectrum shown in the lowest panel of Fig. 5. Looking back at the time history, we also notice a superimposed modulation with a period around $64 \omega_{pe}t$. This is the signature of kinetic Alfvén waves associated with the compression of the magnetic field. As we calculated earlier, $\omega_{ce}/\omega_{pe} \approx 1$ at the reflection point, so the local Alfvén velocity ($V_A = c(\omega_{ce}/\omega_{pe})(m_e/m_i)^{1/2}$) is around $0.5\Delta/\omega_{pe}$. Using the magnetosonic dispersion relation $\omega = \pm k_y V_A$ and the kinetic approximation $k_y = 2\pi/p_i$ with $p_i = 30\Delta$, we find that $\omega = \pm 0.1$, which is consistent with the small twin peaks observed in the power spectrum.

For $k_y=0$, significant turbulence is observed on the electrostatic and compressional magnetic B_z components. Figures 6a and 6b display the time history and the power spectrum of the electrostatic potential ϕ and δB_z , respectively, measured

at $x=80\Delta$ for the dominant mode $k_y=4\pi/L_y$. For both we observe a strong peak at a very low frequency ($\omega \approx -0.018 \omega_{pe}$) which is close to the local ion cyclotron frequency ($\omega_{ci} \approx -0.01 \omega_{pe}$ at the reflection point). After normalization, calculation shows that both fluctuations are of the same order. To check this frequency against the phase shift that we noticed earlier in Fig. 4, we show in Fig. 7 a sequence of the scatter plots of the electron locations (left column) and the corresponding contours of the electrostatic potential (right column). We start at the bottom of the figure with plots measured at time $t=340/\omega_{pe}$ and display each series after a time delay $\delta t=40/\omega_{pe}$ until time $t=500/\omega_{pe}$. We observe clearly the motion of the structures in the $y>0$ direction, which means that the modulation propagates in the direction of the electron diamagnetic drift. The structures move about 10Δ between each snapshot, which corresponds to a phase velocity about $0.25\Delta\omega_{pe}$. Using the dominant mode $k_y=4\pi/L_y$, this translates into a real frequency $\omega \approx -0.025\omega_{pe}$. Taking into account that the Nyquist frequency $\omega \approx 0.008\omega_{pe}$, we find that this estimate is very close to the frequency obtained from the power spectrum analysis.

To explain the growth of an instability due to the presence of strong density and magnetic field gradients, it is natural to think of classical gradient-driven instabilities such as the lower hybrid drift or the drift cyclotron instabilities. However, for $k_x=0$ these instabilities propagate in the direction of the ion diamagnetic drift (Kadmonstev, 1965). Furthermore, the observed frequency is much lower than the lower hybrid. The occurrence of such instabilities may then be excluded, so we do not dispose of ready analytic predictions to compare to our simulation results. To this aim, we have initiated the theoretical analysis of the linear phase. Encouraging preliminary results (Hasegawa, 1988; personal communication) indicate that the instability can be explained by the combined effect of both magnetic and density gradients. The instability condition obtained with a local approximation can be expressed as $4v_0/\kappa\omega_{ci}>1$, where $\kappa=(1/B)\partial B/\partial x-(1/n)\partial n/\partial x>0$ and v_0 is the electron drift velocity. The characteristic frequency obtained is then $\omega=-\omega_{ci}k_y/\kappa$. Considering that the gradient scale length is of the order of 2 ion gyroradii or about 60Δ , the previous relation gives us a frequency $\omega \approx -1.5 \omega_{ci}$ for the dominant mode $k_y=4\pi/L_y$. This result shows that both directions of the phase velocity and frequency range observed in the simulation are reproduced by the analytical calculation. Details of the derivation will be reported elsewhere.

Discussion and Conclusions

In this paper we have examined the fundamental interaction between the magnetosheath flow and the magnetospheric field using a $2^{1/2}$ -dimensional simulation. We deliberately attacked this problem with a very simple model to be able to carry out a full kinetic electromagnetic simulation with a realistic computer time budget. We chose a 2-dimensional magnetic dipole with the field direction perpendicular to the simulation plane and imposed periodic conditions along the y direction. Of course these assumptions limit the extent of our study to the equatorial region of the subsolar magnetopause. In this model, the magnetopause is flat and particles cannot be convected out of the system as they do along the flanks of the real magnetopause. It is thus difficult to assess the relevance of our results to experimental observations. However, several comments can be drawn from the results presented above.

Firstly, one thing that stands very clearly is that because of the growth of the instability, the current layer appears to be thicker than the hybrid gyroradius predicted by the Ferraro-Rosenbluth model. In the past this discrepancy was resolved by arguing that in the presence of a magnetospheric plasma and a cold plasma trapped in the current layer, the ambipolar electric field should be short circuited by the ionospheric plasma being pulled along the field lines (Parker, 1967a). Once the polarization field neutralized, the ions should be able to penetrate fully in the magnetospheric field, and the thickness of the current layer should be just the average ion gyroradius. Utterior sophistications of our model that include magnetospheric plasma and the effects of the ionospheric conductivity should clarify the role of field-aligned currents at the dayside magnetopause. A further broadening of the current sheet due to the ionospheric coupling would certainly help us to use our simulation results to explain the magnetopause thickness measurements. Indeed, Berchem and Russell (1982) showed that the magnetopause current layer is thicker than about 5 ion gyroradii and that the thickening appears to be much better when ordered by the dipole magnetic latitude than any other parameters, so that the current layer is thinnest near the equator. Since the maximum thickness observed in our simulation is less than 3 ion gyroradii, it is difficult to use our present results to explain their observations. Since our simulation is limited to the equatorial plane, a further study of the instability dependence on the field line curvature is needed to resolve whether the instability can account for the latitudinal variation of the observations.

Secondly, another important point shown by our simulation results is the occurrence of a microinstability which has a frequency of the order of the local ion cyclotron frequency and propagates in the direction of the electron drift. Such fluctuations are indeed consistent with the ion cyclotron waves observed at the magnetopause (Fairfield, 1976; Anderson et al., 1982) and the bursty low frequency electrostatic and magnetic fluctuations reported by Gurnett et al. (1979).

Finally, the nonlinear evolution of the instability towards the filamentation of the current sheet by anomalous diffusion is somewhat surprising. However, we have to keep in mind that spacecraft experiments measure fields, not currents. Examination of high resolution magnetic field profiles across the dayside magnetopause very often reveals the presence of small structures inside the magnetic field gradient. Although it has always been difficult to resolve the ambiguity between spatial and temporal variations from the analysis of times series, ISEE dual magnetometer observations (Russell and Elphic, 1978) have shown several cases where spatial structures are clearly identified inside the magnetopause current layer. Examination of high resolution particle data, when they become available, should resolve whether these structures are caused by quasilinear diffusion.

Before concluding, we wish to remark that it is difficult to evaluate the role of the initial boundary conditions in the stabilization of the instability. However, earlier calculations (Berchem and Okuda, 1988) carried out with a much smaller system (128×64) led us to the same results as reported here. This seems to indicate that the walls are far enough away to affect significantly the saturation process. Definite resolution of this question will be possible by improving our code to load more particles and use an even bigger system than the (200×128) grid used here.

In summary, we have developed a $2^{1/2}$ -D (x, y, v_x, v_y, v_z) electromagnetic code to study the formation and the stability of the magnetopause current layer. Our model consists of an unmagnetized subsonic super-Alfvénic plasma flow injected transversely to a two-dimensional vacuum dipolar magnetic field. The results presented are obtained for the simulation of the solar wind-magnetosphere interaction in the subsolar region of the equatorial plane.

1. The simulation results show that during the establishment of the current layer, its thickness is of the order of the hybrid gyroradius as predicted by the Ferraro-Rosenbluth model. However, diagnostics indicate that the current sheet is subject to an instability which broadens the width of the current layer. Ripples with an amplitude of the order of the ion gyroradius appear at the interface between the field and the particles.

2. Perturbations are observed both on the electrostatic field and on the compressional component of the magnetic field. The instability has a frequency of the

order of the local ion cyclotron frequency. However, the modulation propagates in the same direction as the electron diamagnetic drift, which indicates that the instability is not a classical gradient-driven instability such as the lower hybrid drift or ion drift cyclotron instabilities.

3. The nonlinear phase of the instability is characterized by the filamentation of the current layer which causes anomalous diffusion inside the central current sheet.

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Figure Captions

Figure 1: Geometry of the simulation model. The x axis points towards earth, the y axis is directed from dusk to dawn, and the z axis points northward. The $L_x \times L_y$ domain represents the subsolar region of the equatorial plane in which we follow the trajectories of solar wind particles. The ion and electron particles are injected from the left boundary ($x=0$) transversely to a 2D-dipolar magnetic field. We assume that all the quantities are periodic along the y axis.

Figure 2: Composite display of several diagnostics of the simulation at time $t=180/\omega_{pe}$. On the left-hand column, we plotted as a function of the x position: (a) the electron number density n_e , (b) the component E_x of the total electric field, (c) the current density J_y , and (d) the magnetic fields B_z . On the right-hand column we display two-dimensional scatter plots of the electron (e) and ion (f) positions, and respectively (g) and (h) their phase space distribution in the velocity phase space $x-v_x$.

Figure 3: Composite display of the same diagnostics as in Fig. 2 but for time $t=440/\omega_{pe}$.

Figure 4 (a): Dynamics of the formation of the current layer. From left to right, one-dimensional plots of the magnetic field B_z , the current density J_y , the electron density n_e , and the electron locations in the x - y plane, taken from time $t=100/\omega_{pe}$ to time $t=400/\omega_{pe}$. The growth of the instability is clearly seen on the left column.

Figure 4(b): Dynamics of the formation of the current layer (continued). We use the same diagnostics as figure 4(a) from time $t=500/\omega_{pe}$ to time $t=800/\omega_{pe}$. This illustrates the quasilinear and the highly nonlinear phases of the instability.

Figure 5: Time history and power spectrum of the δB_z component of the magnetic field measured at $x=80\Delta$ for mode $k_y=0$. On the upper panel, $\delta B_z(x, k_y)$ is plotted as a function of time. The peak observed about the zero frequency range is due to the motion of the current layer during the compression of the magnetic field. The twin peaks observed at $\omega \approx -0.1\omega_{pe}$ are the signatures of compressional kinetic Alfvén waves.

Figure 6a: Time history and power spectrum of the electrostatic potential ϕ measured at $x=80\Delta$ for mode $k_y=4\pi/L_y$. On the upper panel, both real and imaginary parts of $\phi(x, k_y)$ are plotted as a function of time. The strong peak observed in the very low frequency range $\omega \approx -0.018 \omega_{pe}$ is due to the instability. The small twin peaks observed at $\omega \approx -0.1 \omega_{pe}$ are the signatures of compressional kinetic Alfvén waves.

Figure 6b: Time history and power spectrum of the δB_z component of the magnetic field measured at $x=80\Delta$ for mode $k_y=4\pi/L_y$. On the upper panel, both real and imaginary parts of $\delta B_z(x, k_y)$ are plotted as a function of time. The strong peak observed in the very low frequency range $\omega \approx -0.018 \omega_{pe}$ is due to the instability. The small twin peaks, almost unobservable at $\omega \approx -0.1 \omega_{pe}$, are the signatures of compressional kinetic Alfvén waves.

Figure 7: Scatter plots of the electron x-y location (left-hand column) and contour plots of the electrostatic potential (right-hand column) show the modulation propagates in the same direction as the electron diamagnetic drift at a phase velocity around $0.25\Delta\omega_{pe}$.

Figure 2: $t = 140 / \omega_{pe}$

Figure 3: $t = 440 / \omega_{pe}$

Figure 4a

Figure 4b

Figure 5: Time history and Power spectrum of δB_z ($x = 80\Delta$, $k_y = 0$)

Figure 6a: Time history and Power spectrum of $\Phi(x = 80\Delta, k_y = 4\pi/L_y)$

Figure 6b: Time history and Power spectrum of δB_z ($x = 80\Delta$, $k_y = 4\pi/L_y$)

Figure 7