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Kinetic Alfvén waves as a source of plasma transport at the dayside magnetopause

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Abstract. As the shocked solar wind with variable plasma density and magnetic field impinges on the dayside magnetopause, it is likely to generate large-scale Alfvén waves at the solar wind magnetosphere interface. However, large gradients in the density and magnetic field at the magnetopause boundary effectively couple large-scale Alfvén waves with kinetic Alfvén waves. In this paper we propose that the wave power converted into kinetic Alfvén waves may play an important role in plasma transport at the dayside magnetopause and in electron acceleration along field lines. The transport can occur because, unlike the magnetohydrodynamic (MHD) shear Alfvén wave, the kinetic Alfvén wave has an associated parallel electric field which breaks down the "frozen-in" condition and decouples the plasma from field lines. We calculate the average deviation of the plasma from the field line from which we estimate the diffusion coefficient associated with these "bundles" of decoupled plasma to be approximately $10^9 m^2/s$. The parallel electric field also may lead to acceleration of electrons along field lines in the magnetopause boundary and may possibly provide an explanation for observed counterstreaming electron beams characterized by energies of 50-200 eV.

1. Introduction

The interface between the shocked solar wind and the magnetosphere, the magnetopause, is a transition layer characterized by strong gradients in the characteristic plasma quantities. These strong gradients contain much free energy which can be released in the form of various plasma processes such as magnetic reconnection, Kelvin-Helmholtz instabilities, lower-hybrid drift instability, and kinetic Alfvén waves (KAW). Indeed, the present description of the magnetopause boundary layer is essentially a montage of these various processes which has been constructed over the last 30 years in an attempt to explain the entry of solar wind plasma into the magnetosphere.

The various mechanisms for plasma entry into the magnetosphere can be grouped into the three following classes: viscous interaction, reconnection, and impulsive penetration. Historically, the viscous model proposed by *Axford and Hines* [1961] and the magnetic reconnection proposed by *Dungey* [1961] have been the most widely accepted models, although an alternative mechanism of impulsive plasma penetration has also been recently suggested. Viscous models are based primarily upon the existence of some form of anomalous viscosity in the magnetopause boundary. The diffusion occurs primarily via wave-particle collisions which cause the particles to random walk across the boundary. The flux of particles across the boundary is primarily dependent upon the spectral power of the waves which give rise to the anomalous viscosity. Viscous penetration of particles may arise from lower-hybrid turbulence [*Labelle and Treumann*, 1988; *Gary and Sgro*, 1990; *Treumann et al.*, 1992] or whistler mode waves driven by current gradients [*Drake et al.*, 1994], while momentum can be transported by means of Kelvin-Helmholtz turbulence [*Miura*, 1984]. Enhanced fluctuations in the magnetic and electric fields have been observed at the magnetopause [*Gurnett et al.*, 1979; *Tsurutani et al.*, 1981; *Anderson et al.*, 1982].

Magnetic reconnection provides an alternative model in which the plasma penetrates along reconnected field lines which provide direct access from the solar wind to the magnetosphere. Observational evidence of flux transfer events [*Russell and Elphic*, 1978] and high-speed magnetopause and boundary layer plasma flow [*Sonnerup et al.*, 1981; *Paschmann et al.*, 1986; *Gosling et al.*, 1990] suggest that at least to some degree, bundles of flux do enter the magnetosphere via this mechanism.

The model of impulsive penetration consists primarily of a filament of enhanced momentum which impinges on the magnetopause boundary [*Lemaire et al.*, 1979; *Lundin and Dubinin*, 1984]. It would appear that this process is efficient only for nearly parallel or

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nearly antiparallel orientation of the solar wind magnetic field with respect to the Earth's magnetic field. Efficient transport occurs for antiparallel orientation because conditions are conducive to magnetic reconnection, while for parallel orientation solar wind plasma can penetrate as the result of field line interchange which may occur when the restoring magnetic tension is small. However, the orientations which are favorable for impulsive plasma penetration are rather restricted, and, in general, the filament, instead, bounces off the magnetopause, giving rise to a large-scale surface magnetohydrodynamic (MHD) wave along the boundary [*Ma et al.*, 1991].

In this paper it is our purpose to add transport due to KAWs to the present collection of mechanisms. As the solar wind impinges on the magnetosphere, large-scale MHD waves are generated along the boundary by solar wind fluctuations and the Kelvin-Helmholtz instabilities which can couple to KAWs. The magnetopause is characterized by large gradients in the Alfvén velocity, and these gradients provide an effective mechanism for coupling large-scale structure into small-scale structure when the gradients are the order of an ion gyroradius. As a result, power may be converted from the surface MHD wave to the KAW [*Hasegawa and Chen*, 1975; *Hasegawa*, 1976]. However, as we shall discuss, the KAW is characterized by a parallel electric field so that the plasma may be decoupled from the magnetic field. The field-aligned particle velocities mix the plasma and provide a mechanism for cross-field transport and diffusion. There are two important findings of this model. First, we present an estimate of the diffusion into the magnetosphere, and secondly, we discuss electron acceleration by the parallel electric field of the wave which may be comparable with observations of counterstreaming electrons [*Ogilvie et al.*, 1984; *Takahashi et al.*, 1991].

2. Kinetic Alfvén Waves

Coupling between macroscopic driving forces such as the fluctuating solar wind and microscopic processes such as localized diffusion and heating is of fundamental concern in describing the processes which occur in the magnetosphere. In the most simplistic model the solar wind, with variable plasma density, velocity, or magnetic field, impinges on the magnetospheric cavity (sharp boundary) and sets up discrete plasma surface waves along the magnetopause [*Chen and Hasegawa*, 1974]. If, on the other hand, the magnetopause boundary varies smoothly, then the discreet spectrum breaks up into a continuous spectrum and all modes whose frequency ω satisfies the resonance condition $\omega = k_{\|} v_A(x)$, where v_A is the Alfvén velocity and $k_{\|}$ is the parallel wave number for some point x in the system, are resonant and are characterized by a logarithmic singularity. *Hasegawa* [1976], who first studied the KAW in the context of auroral electron acceleration, showed that this resonance may be resolved by expanding the kinetic equations to first order in the ion gyroradius. The singularity is removed because, unlike the electrons, the ions do not move with the field lines when the perpendicular wavelength of the field is the order of the gyroradius. As a result, the ions on different field lines interact electrostatically so that an isolated resonance (infinite current sheet) on a particular field line will not occur. The charge separation not only allows the wave to travel across field lines, but also gives rise to a parallel electric field.

The kinetic Alfvén waves can be generated at the magnetopause due to the presence of strong gradients in the density and magnetic field [*Eastman and Hones*, 1979]. In addition, many other wave modes can also be excited in the presence of strong gradients. These phenomena can be described in terms of an approximation to the coupled Vlasov-Maxwell set of equations. It is important to consider what approximations might be relevant for considering KAW generation.

The dominant process that gives rise to KAW generation is the existence of field line resonances. These resonances are associated with a singularity in the cold plasma description of the shear Alfvén mode. As such, it can be expected that as a first approximation compressional effects will not affect the resolution of the singularity. In fact, if one considers a full MHD description of a thermal plasma, it is easy to show that coupling between the shear and compressional modes takes place primarily via curvature in the magnetic field lines [*Cheng*, 1991]. Curvature couples compressional perturbations to shear perturbations through the so-called "ballooning" term, which can be important for a high-β plasma. In the presence of "good curvature," as is found at the magnetopause, this term does not give rise to a ballooning instability but, rather, primarily serves to modify the location of the field line resonance by increasing the effective frequency of the wave. In effect, even in a high-β plasma the fundamental physics is dominated by the shear Alfvén wave and the resolution of the field line resonance so that as a first approximation, one can neglect the compression of the magnetic field.

We should point out that the observed plasma β in the magnetosheath ranges from 0.2 to 4 [*Paschmann et al.*, 1986; *Gosling et al.*, 1990], and there may be coupling between the shear Alfvén mode and compressional mode in the high-β case. However, we expect that coupling will not affect the fundamental physics involved in the mode conversion process. In the magnetospheric side of the magnetopause the plasma β is usually small ($\beta < 0.2$). In the depletion layer just outside the magnetopause, where the plasma is squeezed into the flanks by the pileup of magnetic field, the plasma β can also be low, and an incompressible approximation is valid.

Another important wave which is associated with strong density gradients is the drift wave. For the parameters of interest, electron drift waves and waves generated near the lower-hybrid frequency are unimportant for a linear calculation because their frequency is much larger than the Alfvén frequency. Any coupling that takes place will be nonlinear in nature. The ion drift frequency is, however, comparable with the Alfvén frequency so that ion drift waves may be important. The primary effect of ion drift waves on the shear Alfvén mode is to split the mode, shifting it to higher and lower frequency. The field line resonances still exist but are shifted to a different location. The resolution of the resonance should remain the same, although the range of the continuous spectrum can be increased.

In order to describe the physics most effectively and simply, near the field line resonance we assume magnetic incompressibility and ignore modifications due to drift waves. We have discussed the nature of these approximation above. This model should be sufficient to describe the physics that occurs at the magnetopause. To resolve the MHD resonance, the ion Vlasov equation is expanded in terms of the ion gyroradius. Assuming variations in density and magnetic field along the x direction only, the Vlasov-Maxwell equations reduce to [*Hasegawa*, 1976]

$$\left(\frac{\omega^2}{k_\parallel^2 v_A^2}\frac{3}{4}\rho_i^2\frac{d^3}{dx^3}+\frac{d^2}{dx^2}\frac{1}{g}\frac{T_e}{T_i}\rho_i^2\frac{d}{dx}\right)\left(g\frac{d\phi}{dx}\right)$$
$$+\left[\frac{d}{dx}\left(\frac{\omega^2}{k_\parallel^2 v_A^2}g-1\right)\frac{d}{dx}-k_y^2\left(\frac{\omega^2}{k_\parallel^2 v_A^2}g-1\right)\right]\phi=0 \qquad (1)$$

Here ρ_i is the ion gyroradius; g is the density profile, $g = n(x)/n_{\max}$, normalized to the maximum density, $n_{\max}$; T_e and T_i are, respectively, the electron and ion temperatures. In (1) the fields have been expressed in terms of a perpendicular scalar potential ϕ such that $\delta\mathbf{E}_\perp = -\nabla_\perp\phi$, where $\delta\mathbf{E}_\perp$ is the perturbed perpendicular electric field. For a homogeneous plasma this equation decouples into two MHD surface waves (represented by the dominant balance of the third and fourth terms) and two KAWs (represented by the dominant balance of the first, second, and third terms) with the respective dispersion relations $\omega = \pm k_\parallel v_A$ and $\omega^2 = k_\parallel^2 v_A^2(1 + k_\perp^2\rho^2)$, where $\rho^2 = (3/4 + T_e/T_i)\rho_i^2$. Without the first and second terms, (1) exhibits the logarithmic singularity associated with the MHD surface wave resonance, however, the first and second terms resolve the singularity. Indeed, using a suitable linear approximation for the density and magnetic field gradients near the singular point ($x \sim 0$), it is possible to reduce these equations to a driven, second-order ordinary differential equation for the x component of the perturbed electric field δE_x

$$\rho^2(d^2\delta E_x/dx^2) + \kappa x\delta E_x = E_0 \qquad (2)$$

where κ is the inverse scale length of the density gradient, $\kappa \equiv |dln(n)/dx|$, and E_0 is the integration constant obtained by integrating (1). This equation can be interpreted as describing KAWs driven by an MHD surface wave with amplitude E_0. The extent of the linear coupling of the KAWs to the MHD driver can be determined by solving (2) and matching to the appropriate boundary conditions (no incoming wave power in the KAW mode). The solution of (2) is readily found in terms of the Airy functions Ai, Bi, and Gi as functions of x/δ [*Hasegawa*, 1976], where $\delta \equiv (\rho^2/\kappa)^{1/3}$. The KAW has the asymptotic behavior

$$\delta E_{x,\mathrm{KAW}} \sim \exp[i\frac{2}{3}(-\frac{x}{\delta})^{3/2}] \qquad (3)$$

which is oscillatory for $x < 0$, where the density is higher and exponentially decaying for $x > 0$, where the density is lower. The surface MHD wave has the asymptotic behavior $E_0/\kappa x$, which is the balance of the zeroth-rder term with the inhomogeneous term in (2). In Figure 1 we show a plot of the solution to (2). In magnetospheric context the coordinate x is across the magnetopause (from solar wind to magnetosphere) and $x = 0$ defines the location of the resonant point where $\omega = k_\parallel v_A(x)$ (the surface wave is driven at the frequency ω). For $x/\delta < 0$ the solution is dominated by the KAW, while for large $x/\delta > 0$ the solution is dominated by the MHD surface wave in that the KAW decays exponentially. The asymptotic forms of the MHD surface wave and KAW have been plotted for comparison with the solution of (2). Physically, the driven MHD wave couples through the density gradient to the KAW, which propagates to the left into the solar wind and also penetrates into the magnetosphere with decaying exponential behavior. For $\rho \sim 50$ km and $\kappa^{-1} \sim 500$ km [*Berchem and Russell*, 1982] the penetration at the magnetopause boundary can be estimated from the wave scale length δ to be ~ 100 km, which is comparable to the width of the magnetopause.

In the following two sections we will exploit some of the properties of the KAW to estimate the amount of

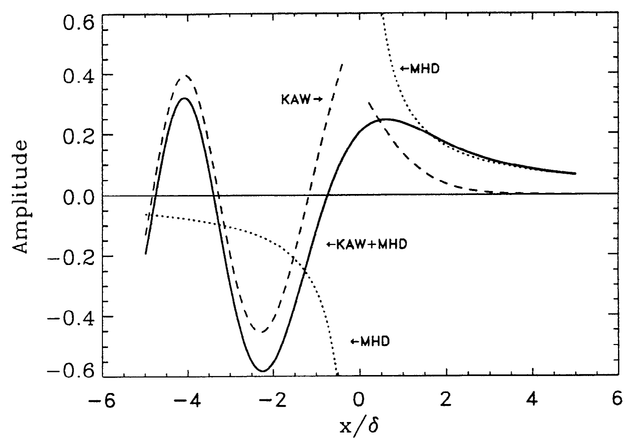

Figure 1. The wave solution of equation (2) near the resonant point $x=0$ is plotted as the solid line. The coordinate x is across the magnetopause from solar wind to magnetosphere and is normalized to the scale length δ of the kinetic Alfvén waves (KAW). In the asymptotic regions the solution can be decomposed into the asymptotic forms of the magnetohydrodynamic (MHD) and KAW, as indicated by the dotted and dashed lines, respectively. On the magnetosphere side the KAW decays after several scale lengths and the solution is asymptotic to the MHD surface wave. The resonant region is dominated by the KAW which reflects back into the solar wind.

diffusion associated with a driven MHD surface wave and secondly, we will estimate the magnitude of the parallel electric field associated with the wave.

2.1. Diffusion and Kinetic Alfvén Waves

In this section we estimate the diffusion associated with a driven KAW. In this approach we assume that there is a single KAW with a fixed frequency. This approach is to be differentiated from the situation in which there is a broadband spectrum of wave power at low frequencies. As such, this process does not lie in the same class as lower-hybrid and Kelvin-Helmholtz turbulence.

Unlike the shear Alfvén wave the KAW exhibits certain properties conducive to plasma transport. The KAW is characterized by charge separation of ions and electrons which leads to the generation of a field-aligned electric field ($\delta E_{\parallel}$) and cross-field wave propagation. The parallel electric field $\delta E_{\parallel}$ is related to the perpendicular field $\delta E_{\perp}$ by

$$\delta E_{\parallel} = \rho_i^2 \frac{T_e}{T_i} \frac{\partial}{\partial z} (\nabla \cdot \delta \mathbf{E}_{\perp}) \tag{4}$$

The existence of $\delta E_{\parallel}$ is associated with the breakdown of ideal MHD and decoupling of the plasma from the magnetic field. This decoupling can be estimated from the characteristics of the KAW. Consider a configuration with x across the magnetopause and with z along and y across the unperturbed magnetic field. In the following sections we will emphasize transport in the x direction which is across the magnetopause. For this configuration, $k_x \sim \delta^{-1} \gg k_y \gg k_z$ and the dominant components of the Alfvén wave are E_x and B_y. The dispersion relation of the KAW is

$$\omega^2 = k_{\parallel}^2 v_A^2 (1 + k_x^2 \rho^2) \tag{5}$$

The magnetic field components are modified accordingly

$$\delta \mathbf{B} = \frac{k_{\parallel}}{\omega} \hat{b} \times \delta \mathbf{E} (1 + \frac{T_e}{T_i} k_x^2 {\rho_i}^2). \tag{6}$$

The position of the field line deviation perpendicular to the unperturbed magnetic field direction $\hat{b} = \hat{z}$ is given by the vector $\mathbf{r_B}$. This deviation is best determined in the wave frame. If z is the direction along the field line, then $z^* = z - \omega t / k_{\parallel}$ is the z coordinate in the wave frame so that $d\mathbf{r_B}/dz^* = \delta \mathbf{B}/B_0$ and hence

$$\frac{d\mathbf{r_B}}{dt} = \frac{\delta \mathbf{B}}{B_0} \frac{dz^*}{dt} = -\frac{\delta \mathbf{B}}{B_0} \frac{\omega}{k_{\parallel}}. \tag{7}$$

Thus

$$\frac{d\mathbf{r_B}}{dt} = \frac{\delta \mathbf{E} \times \hat{b}}{B_0} (1 + \frac{T_e}{T_i} k_x^2 {\rho_i}^2), \tag{8}$$

whereas the plasma displacement $\mathbf{r_P}$ is given by the $\mathbf{E} \times \mathbf{B}$ drift

$$\frac{d\mathbf{r_P}}{dt} = \frac{\delta \mathbf{E} \times \hat{b}}{B_0}. \tag{9}$$

The relative difference in the plasma and magnetic field velocities in the x direction v_D is

$$v_D = \frac{T_e}{T_i} k_x^2 {\rho_i}^2 (\frac{\delta B_x}{B_0}) v_A. \tag{10}$$

The deviation between the plasma and the field lines in the x direction (across the magnetopause) can be obtained by integrating the velocities over half a wave period and comparing the difference in the displacement.

$$\Delta x \equiv \Delta x_B - \Delta x_P = 2 \frac{T_e}{T_i} \frac{k_x^2 \rho_i^2}{k_{\parallel}} (\frac{\delta B_x}{B_0}). \tag{11}$$

In Figure 2 we illustrate the displacement of the plasma and magnetic field. The magnetic field is more deformed than the plasma so that, in effect, plasma is left behind on the magnetospheric side of the magnetopause. Using typical parameters, $T_e/T_i \sim 0.2$, $k_x^2 \rho_i^2 \sim \rho_i^2/\delta^2 \sim 0.25$, $\lambda_{\parallel} = 2\pi k_{\parallel}^{-1} = 10 R_E$, $\delta B_x / B_0 = 0.2$, $\delta B_y / B_0 \sim 1$, and $v_A = 400$ km/s [*Berchem and Russell*, 1982; *Russell and Elphic*, 1978; *Labelle and Treumann*, 1988; *Song et al.*, 1993a, b], the deviation can be estimated as $\Delta x \sim 200$ km which is comparable to the width of the magnetopause, and the velocity in (10) can be estimated as $v_D \sim 4$ km/s. Note that the frequency of KAWs at the dayside magnetopause is typically $f \lesssim 0.1$ Hz. The observed wave amplitude δB at this frequency is $\delta B \sim 30$ nT$\sim B_0$.

In the context of fluid theory the plasma oscillates about the zeroth-order magnetic field, and although the perturbed field and fluid are decoupled, no net transport of fluid particles will occur. There are two mechanisms, however, which could lead to transport of particles. First, parallel particle velocities cause individual particles to shift from their initial field lines which may lead to a diffusive process which should be reasonably efficient for $v_{\parallel} \sim v_A$. This effect is illustrated schemati-

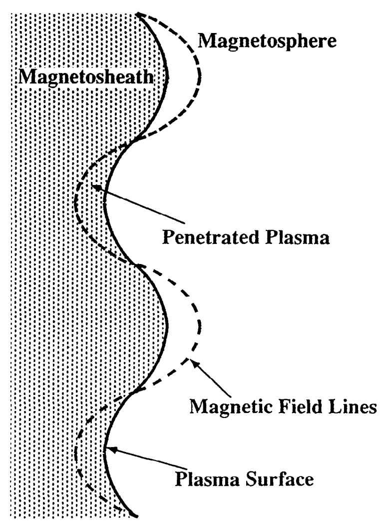

Figure 2. Plasma penetration into the magnetosphere may occur when the magnetic field decouples from the plasma at the magnetopause boundary. As shown by the dashed line, the magnetic field displacement exceeds the plasma displacement and leaves plasma behind in the magnetosphere. Sudden damping of the wave or particle motion along the magnetic field can lead to plasma transport across field lines.

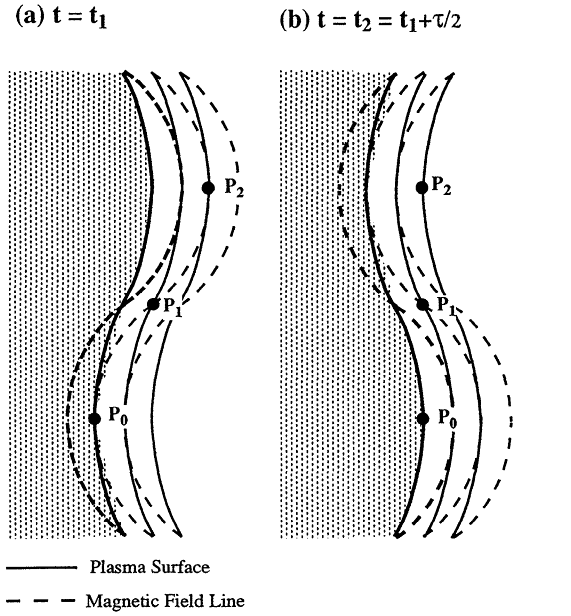

Figure 3. Parallel particle velocities can lead to transport of particles across field lines in the presence of a KAW. (a) An ion found at point P_0 streams along the magnetic field toward point P_1 or P_2 in time $\tau/2$, depending on its parallel velocity. (b) Due to the finite Lamor radius effect, the ion which moves to P_2 is somewhat unmagnetized and decoupled from its initial field line. Other nonresonant particles such as those that reach P_1 may later follow the field lines back into the magnetosheath.

cally in Figure 3. A particle initially at the point P_0 has a parallel velocity along the magnetic field and moves freely along the magnetic field to a point P_1 or if the velocity is larger, to the point P_2 during the half-period of the wave $\tau/2$. Because $k_\perp \rho_i \sim 1$, the ions are not strongly magnetized so that they also move across field lines. Therefore, a particle which travels to P_2 will have a net deviation from its initial field line. The deviation can be estimated in terms of the plasma velocity across the field line which is the order of Δx obtained in (11). This process is most efficient for resonant particles with $v_\| \sim v_A$.

An alternative way that plasma can enter the magnetosphere might be from dissipation of the KAW. We have not heretofore discussed the nonlinear behavior of these modes, obviously an important effect if $\delta B \sim B_0$. As has been discussed by *Hasegawa and Chen* [1976], the KAW tends to dissipate rapidly through nonlinear ion Landau damping and phase mixing of several modes. The relevant decay process in this instance is coupling between a strongly damped ion acoustic wave and two KAWs and involves nonlinear ion Landau damping. We should stress that this mechanism works efficiently even when the ion acoustic wave is heavily damped, as is the case for $T_e/T_i \sim 0.2$. Hence when the KAW is excited, the plasma decouples from the field line. If the KAW dissipates suddenly on the Alfvén timescale, then the plasma would be left decoupled from the original field line and particle transport could occur. Clearly, the nonlinear evolution of the KAW is an important aspect of this problem and is a clear direction for further consideration.

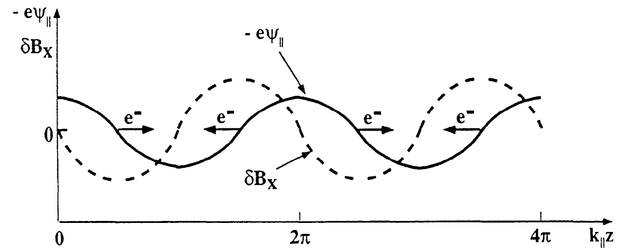

Figure 4. The parallel electric potential $\Psi_\|$ associated with the KAW provides an explanation for measurements of low energy counterstreaming electron beams which are observed in the low-latitude boundary layer. Electrons may be accelerated up to 100 eV in the potential wells which have macroscopic scale along the field lines.

As a result, filaments of plasma can enter into the magnetosphere with a typical cross-field velocity given by the relative displacement between the field lines and plasma in bundles v_D which has characteristic scale size Δx. An estimate for the diffusion coefficient D resulting from this process is

$$D \sim v_D \Delta x \sim 2(\frac{T_e}{T_i} k_x^2 \rho_i^2)^2 \frac{v_A}{k_\|} (\frac{\delta B_x}{B_0})^2. \tag{12}$$

For the typical parameters discussed above the diffusion is estimated to be the correct order of magnitude (10^9 m^2/s). As will be discussed later, the estimated diffusion coefficient in (12) is virtually identical to that associated with each wave component in the quasi-linear theory.

2.2. Parallel Electric Field and Counterstreaming Beams

As we mentioned, one of the characteristics of KAWs is the presence of a field-aligned electric field which leads to the decoupling of plasma from magnetic field. Another important consequence of the parallel electric field is the possibility that it may accelerate electrons of ionospheric origin. As illustrated in Figure 4, a strong parallel electric field can accelerate low-energy ionospheric electrons in the potential well in opposite directions so that they may appear to be counterstreaming. This property of the KAW suggests an explanation for measurements of low-energy counterstreaming electron beams observed in the low-latitude boundary layer [*Ogilvie et al.*, 1984; *Takahashi et al.*, 1991]. Indeed, electron beams are observed simultaneously with strong wave activity at low frequencies [*Song et al.*, 1993c].

The parallel potential $\Psi_\|$ can be estimated for the KAW [*Hasegawa*, 1976]

$$\Psi_\| \approx \frac{T_e}{T_i} k_x^2 \rho_i^2 \frac{v_A \delta B_x}{k_y}, \tag{13}$$

which is ~100 V for typical parameters $T_e/T_i \sim 0.2$, $k_x^2 \rho^2 \sim 0.25$, $v_A \sim 400$ km/s, $\lambda_y = 2\pi k_y^{-1} \sim 4000$ km, $\delta B_x/B_0 \sim 0.2$, and $B_0 \sim 40$ nT . The observed electron distribution function in the low-latitude boundary layer indicates that counterstreaming electrons are present with typical energies ranging from 50 to 200 eV. These electrons could either be observed in the presence of the wave or could be accelerated by the wave and continue to persist after the wave has dissipated.

3. Comparison of Filamentary and Turbulent Diffusion Models

In this paper we have presented a model in which plasma decouples from the magnetic field lines and, as such, can be compared with observations of mixed filamentary structures which characterize the magnetopause [*Eastman and Hones*, 1979; *Sckopke et al.*, 1981]. In this model, we consider only a single driven surface Alfvén wave and estimate the filamentary diffusion associated with it. If, instead, the MHD wave gives rise to a broadband spectrum of turbulence (which is not unlikely because the KAW spreads in frequency rapidly, either by decay to ion acoustic waves or through electron Landau damping), then the diffusion coefficient can be calculated using the standard quasi-linear theory [*Hasegawa and Mima*, 1978]. In fact, wave observations near the subsolar magnetopause clearly show the existence of a low-frequency broadband spectrum of waves in the sheath transition layer [*Song et al.*, 1993a, b]. Because these waves are found in a region with a strong density gradient, it is reasonable, based upon our previous discussion of KAWs, to assume that a substantial fraction of this wave power is in the KAW mode and has a parallel electric field.

In general, it can be shown that for arbitrary low-frequency turbulence with a Gaussian distribution function the ion diffusion coefficient is

$$D_{\text{turb}} = \left(\frac{\pi}{8}\right)^{1/2} \sum_k \frac{k_y^2 |E_{\|k}|^2}{k_\|^2 B_0^2} \frac{1}{|k_\||v_{ti}} \exp(-v_A^2/2v_{ti}^2) \tag{14}$$

where v_{ti} is the ion thermal velocity. From (14) it is clear that a parallel electric field is associated with the diffusion process. Inserting $E_\|$ for the KAW into this expression, the diffusion coefficient can be estimated as

$$D_{\text{turb}} = \sum_k \left(\frac{T_e}{T_i} k_x^2 \rho_i^2\right)^2 \left(\frac{v_A}{k_\|}\right) \left(\frac{\delta B_{xk}}{B_0}\right)^2 \left[\left(\frac{\pi}{8}\right)^{1/2} \left(\frac{v_A}{v_{ti}}\right) \times \exp(-v_A^2/2v_{ti}^2)\right] \tag{15}$$

where we have assumed that $T_e/T_i < 1$. For each wave in the k summation the contribution is the same as (12), modified by the term in brackets which can be interpreted as a geometrical factor multiplied by the fraction of particles resonant with the wave. If the observed low-frequency waves at the magnetopause are KAWs, then the waves may lead to $D_{\text{turb}} \sim 10^9$ m^2/s.

4. Conclusions

In summary, we have discussed the KAW as a source for particle entry across the magnetopause. Because of the large gradients in density and magnetic field that are found at the magnetopause, it is likely that solar wind fluctuations generate surface MHD waves which mode convert into KAWs. Because the scale length of KAWs is the order of ρ_i, the waves can easily couple with particles and cause transport across the ambient magnetic field.

In this paper we have estimated the decoupling of plasma from the field and have found the deviation to be ~200 km which is the order of the width of the magnetopause. On the slow timescale of the wave the diffusion process may be considered to be the decoupling of bundles of plasma from magnetic field at the magnetopause boundary. This mechanism may explain the observed filamentary structures observed in the low-latitude boundary layer [*Eastman and Hones*, 1979; *Sckopke et al.*, 1981]. An order of magnitude estimate of the diffusion coefficient yields the correct order of magnitude (10^9 m^2/s). The effects of parallel velocity on the diffusive process are critical and allow particles to decouple from their initial field lines and be transported perpendicular to the field. An obvious extension of this work will be to compute the orbits of test particles in a KAW, using different initial velocities.

A further implication of our analysis is that a strong parallel electric field is associated with the KAW which may energize particles to energies of the order of 100 eV. This characteristic of the KAW may be applicable to the observations of counterstreaming electron beams which are frequently observed in the magnetopause boundary layer.

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