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Kinetic Alfvén waves and plasma transport at the magnetopause

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Abstract. Large amplitude compressional type ULF waves can propagate from the magnetosheath to the magnetopause where there are large gradients in density, pressure and magnetic field. These gradients efficiently couple compressional waves with shear/kinetic Alfvén waves near the Alfvén field-line resonance location ($\omega = k_{\|}v_A$). We present a solution of the kinetic-MHD wave equations for this process using a realistic equilibrium profile including full ion Larmor radius effects and wave-particle resonance interactions for electrons and ions to model the dissipation. For northward IMF a KAW propagates backward to the magnetosheath. For southward IMF the wave remains in the magnetopause but can propagate through the $k_{\|} = 0$ location. The quasilinear theory predicts that transport due to KAWs at the magnetopause primarily results from the perpendicular electric field coupling with magnetic drift effects with diffusion coefficient $D_{\perp} \sim 10^9$ m^2/s. For southward IMF additional transport can occur because magnetic islands form at the $k_{\|} = 0$ location. Due to the broadband nature of the observed waves these islands can overlap leading to stochastic transport which is larger than that due to quasilinear effects.

Introduction

Observations in the magnetosheath and magnetopause indicate that ultra low-frequency (ULF) waves with frequency less than 1 Hz are the dominant fluctuations. However, the origin of these ULF waves at the magnetopause remains an ongoing issue. In the magnetosheath substantial ULF compressional wave activity exists nearly all of the time, and under most conditions the power level of wave activity is usually 10–100 times larger than the wave activity in the solar wind or the magnetosphere[*Anderson*, 1995, and references therein]. Specific magnetopause crossings have been studied for both northward [*Song et al.*, 1993b] and southward [*Rezeau et al.*, 1989; *Song et al.*, 1993a] IMF conditions. Interestingly enough, ULF wave activity in the magnetopause exhibits strikingly different characteristics than waves observed at the magnetosheath: magnetosheath waves are primarily compressional while waves at the same frequency range at the magnetopause are transverse and often have larger amplitude. For southward IMF the fluctuation level of the transverse waves at the magnetopause are typically enhanced by a factor of 10 over the level of the compressional waves in the magnetosheath[*Rezeau et al.*, 1989]. These observations led *Belmont et al.* [1995] to suggest that coupling between the compressional waves and the Alfvén resonance based on ideal MHD theory was responsible for the observed wave signatures. However, until now there has been no satisfactory theory for these wave observations. In this paper we present the theory of mode conversion of compressional waves into kinetic Alfvén waves (KAWs) at the magnetopause to explain these observed wave features at the magnetopause.

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Presently, it seems to be generally accepted that the various high frequency (larger than ion cyclotron frequency) waves cannot provide the necessary transport [*LaBelle and Treumann*, 1988; *Treumann et al.*, 1994] to explain the amount of solar wind plasma that enters the magnetosphere because the high frequency power spectrum near the magnetopause is weak and obeys a power law which falls off rapidly. On the other hand, there is a general conjecture that magnetic reconnection (or merging) is the dominant plasma transport mechanism at the magnetopause, although no comprehensive theoretical study has been performed. However, ULF transverse waves are commonly observed with dominant wave power at the magnetopause [*Song et al.*, 1993a, 1993b; *Rezeau et al.*, 1989, 1993], and plasma transport resulting from ULF waves has not been studied. It is well known [*Cheng et al.*, 1994] that in laboratory plasmas MHD waves are a leading cause of plasma transport. In this paper, we will address transport which results from the observed ULF waves which have dominant wave power at the magnetopause.

Our proposed scenario for magnetopause wave activity is that compressional MHD waves in the magnetosheath propagate to the magnetopause and couple with transverse KAWs at the magnetopause. General magnetopause features of the background magnetic field, density, flow, and temperature have been surveyed for both high and low magnetic shear cases [*Paschmann et al.*, 1986, 1993]. In both cases background gradients are large and the Alfvén velocity, v_A, can monotonically increase across the magnetopause from the magnetosheath side to the magnetosphere side by about a factor of 10. For a typical parallel wavelength of 1 R_E and 400km/s $< v_A <$ 4000km/s at the magnetopause, the Alfvén field line resonance frequency, $\omega_A = k_{\|}v_A$, forms a continuous spectrum and spans the range 60 mHz $< f <$ 600 mHz which overlaps with the frequency band of magnetosheath compressional type waves of maximum wave power. Near the resonance location where the compressional wave frequency, ω, matches ω_A, the compressional wave couples strongly with the shear Alfvén wave which has a logarithmic singularity in transverse wave power according to the MHD theory. However, this resonance singularity is resolved by the ion gyroradius effects [*Hasegawa and Chen*, 1976], which also cause MHD mode conversion into KAWs. When the perpendicular wavelength is on the order of the gyroradius, ion motion decouples from the electron motion due to polarization drift effects leading to charge separation. Typically $v_e > \omega/k_{\|} > v_i$ so that electrons can maintain charge quasi-neutrality which provides a finite parallel electric field for KAWs. In addition, KAWs exhibit two properties in contrast to the incoming compressional waves: their polarization is primarily transverse and

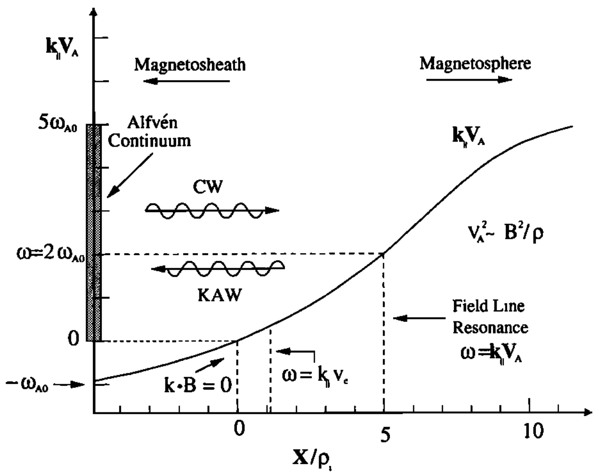

Figure 1. Alfvén frequency for a typical magnetopause crossing from the magnetosheath to the magnetosphere for southward IMF. The position is shown scaled to ρ_i, and the magnetopause scale length is 10 ρ_i. The Alfvén velocity is enhanced by a factor of 10 across the magnetopause. Compressional type waves (CW) with frequency ω propagate to the field line resonance location ($\omega = k_\parallel v_A$ at $x = 5\rho_i$) and mode convert to backward propagating kinetic Alfvén waves. Because of the magnetic drift effects, KAWs can propagate to the $k_\parallel = 0$ location ($x = 0$) and drive time-dependent (patchy) magnetic reconnection.

their wave amplitude is strongly enhanced (by $\mathcal{O}(10)$) over the level of the compressional waves in the magnetosheath.

The KAW can cause significant plasma transport. Based on the quasilinear theory, resonant particles are accelerated (or decelerated) along the field lines by the parallel electric field of KAWs and their perpendicular magnetic drift motion is modified by the perpendicular electric field. Resonating ions can thus decouple from the magnetic field lines and contribute to a radial particle flux across the field lines. For a diffusion coefficient of $\mathcal{O}(10^9)\mathrm{m}^2/\mathrm{s}$, these ions can move across the magnetopause in about $\mathcal{O}(10^2)$ wave periods.

These wave and transport processes at the magnetopause depend on the IMF direction. For northward IMF, the KAW has largest amplitude near the resonance location and radiates back into the magnetosheath. Based on the quasilinear theory, transport due to broadband KAWs will be on the order of 10^9 m^2/s for typical magnetopause parameters. For southward IMF, strong magnetic shear in the current layer gives rise to $k_\parallel = 0$ locations as illustrated in Figure 1 where KAWs form magnetic islands. Because the wave spectrum is broadband, it is likely that many such islands form at different locations and overlap leading to massive transport well above the quasilinear level.

Kinetic Alfvén Waves at Magnetopause

We investigate the KAW structure at the magnetopause using the kinetic-MHD model [*Cheng*, 1991]. This model incorporates kinetic effects into the MHD equations using Ohm's law and the momentum equation obtained from the gyrokinetic equation. The essential approximations required by this model are that $\omega \ll \Omega_i$, $k_\perp L_\perp \gg 1$, $k_\parallel L_\parallel \gg 1$, and $\rho_i \ll L_\perp$, where $\rho_i = v_i/\Omega_i$, $v_i^2 = T_i/m_i$, Ω_i is the ion cyclotron frequency, and $L_{\perp,\parallel}$ are the magnetopause perpendicular and parallel equilibrium scale lengths, respectively. At the magnetopause, these approximations are satisfied because $\rho_i \sim 50$ km, $L_\perp \sim 10\ \rho_i$ and $L_\parallel \sim 10\ \mathrm{R_E}$.

This model is appropriate for studying this problem because it: (a) includes full Larmor radius effects for ions, (b) includes diamagnetic, ∇B, and curvature drifts, (c) is valid in a high β plasma, and (d) includes wave-particle resonance. The ion Larmor radius effects are essential in describing mode conversion of MHD waves into KAWs. Because $k_\perp \rho_i \sim 1$, full ion Larmor radius effects must be included. Diamagnetic and magnetic drift frequencies can be on the order of the Alfvén frequency and their effects must be included in the region of large plasma and magnetic field gradient. The wave particle resonance is very important for determining the KAW structure and plasma transport. The magnetic drift not only contributes to the wave-particle resonance but also dominates plasma transport for typical magnetopause conditions.

For simplicity, we take a one-dimensional model with variation along the radial direction, x, from the magnetosheath to the magnetopause. We employ the Padé approximation, $(1 - I_0(b)e^{-b})/b \approx 1/(1+b)$ where $b = k_\perp^2 \rho_i^2$, so that our analysis is a good approximation for studying full ion Larmor radius effects. Given these approximations, the wave equation which couples the KAW with the compressional wave may be obtained from *Cheng* [1991] and is given by

$$k_\parallel^2 v_A^2 (1+\Delta)\rho_i^2 \frac{d^4\phi}{dx^4} + \frac{d}{dx}(\tilde{\omega}^2 - k_\parallel^2 v_A^2)\frac{d\phi}{dx} + (\tilde{\omega}^2 - k_\parallel^2 v_A^2)(\tilde{\omega}^2 - k_y^2 v_A^2 - k_\parallel^2 v_A^2)\phi / v_A^2 = 0, \quad (1)$$

where $\tilde{\omega}^2 = \omega(\omega - \omega_{*i})$ where $\omega_{*i} = \mathbf{k}\cdot\mathbf{b}\times\nabla \ln P_i/\Omega_i$ is the ion diamagnetic drift frequency

$$\Delta = \frac{T_e(1 - \omega_{*i}/\omega)}{T_i W_e + T_e W_i}, \quad (2)$$

$$W_{e,i} = \frac{T}{n}\int d^3v \frac{\partial F}{\partial \mathcal{E}} \frac{(\omega - \omega_{*T})k_\parallel v_\parallel J_0^2(k_\perp v_\perp/\Omega)}{\omega(\omega - k_\parallel v_\parallel - \omega_d)}, \quad (3)$$

where $\omega_{*T} = \mathbf{k}\cdot\mathbf{b}\times\nabla F/(\Omega \partial F/\partial \mathcal{E})]$, and $\omega_d = \mathbf{k}\cdot(\mathbf{b}/\Omega)\times[(v_\perp^2/2)\nabla \ln B + v_\parallel^2 \kappa]$ is the magnetic drift frequency.

We solve Eq. (1) using asymptotic analysis. Away from the field line resonance location, two spatial scales characterize the solution: an MHD scale which is the balance of the second and third terms of Eq. (1),

$$v_A^2 \frac{d}{dx}(\tilde{\omega}^2 - k_\parallel^2 v_A^2)\frac{d\phi}{dx} + (\tilde{\omega}^2 - k_\parallel^2 v_A^2)(\tilde{\omega}^2 - k_y^2 v_A^2 - k_\parallel^2 v_A^2)\phi = 0, \quad (4)$$

and a kinetic scale which is the balance of the first and second terms of Eq. (1)

$$\rho_i^2(1+\Delta)\frac{d^2\phi}{dx^2} + (\tilde{\omega}^2/k_\parallel^2 v_A^2 - 1)\phi = 0. \quad (5)$$

Near the Alfvén resonance (at x_0) an inner equation can be found by linearizing $\tilde{\omega}^2 - k_\parallel^2 v_A^2 \approx -\tilde{\omega}_0^2(\bar{x}/L_B)(1+\mathcal{O}(\bar{x}/L_B)^2)$ where $\bar{x} = (x - x_0)$. In terms of $X = \bar{x}/\delta$ with $\delta^3 = \rho_i^2 L_B(1+\Delta)$, Eq. (1) reduces to the inhomogeneous Airy equation,

$$\frac{d^2}{dX^2}\frac{d\phi}{dX} - X\frac{d\phi}{dX} = E_{0x}\left(\frac{L_B}{\rho_i}\right)^{2/3}, \quad (6)$$

where E_{0x} is amplitude of the MHD wave, and has the well known solution [*Hasegawa and Chen*, 1976]

$$\delta E_x = \pi E_{0x}\left(\frac{L_B}{\rho_i}\right)^{2/3}\left[c_1 \mathrm{Ai}(\frac{\bar{x}}{\delta}) + c_2 \mathrm{Bi}(\frac{\bar{x}}{\delta}) + \mathrm{Gi}(\frac{\bar{x}}{\delta})\right], \quad (7)$$

where Ai, Bi, and Gi are Airy functions defined in *Abramowitz*

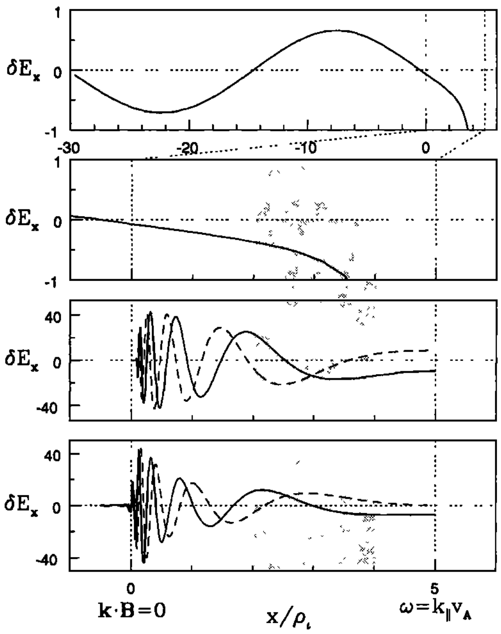

Figure 2. The upper two panels show the outer MHD solution for δE_x which satisfies the boundary condition of an incoming compressional wave. The third panel shows the kinetic Alfvén wave solution without including diamagnetic and magnetic drift effects. The dashed and dotted lines correspond to the real and imaginary parts of the solution. The bottom panel shows the solution of KAW by including realistic magnetic drift effects. The outer solutions are matched to the inner analytic solution in the shaded region.

and Stegun [1970], c_1 and c_2 are determined by the matching to the outer WKB solutions. Note that there is a overlap region for matching the inner solution, valid for $x < L_B$, with the outer solutions, valid for $x > (\rho_i^2 L_B/4)^{1/3}$ (for typical magnetopause parameters $10 \gg |\bar{x}/\rho_i| \gg 1$). Eqs. (4) and (5) are solved numerically to constrain the outer solution to have only an incoming MHD wave. The numerical match to the inner solution determines the coefficients, c_1 and c_2. The relative amplitudes of the two outer solutions is $(\delta E_{KAW}/\delta E_{MS})_{out} \sim \sqrt{\pi}(L_B/\rho_i)^{2/3} \sim 10$.

For weak magnetic shear at the magnetopause (northward IMF cases), Eq. (7) shows that a KAW will propagate backward towards the magnetosheath. The strong magnetic shear case is more interesting, and we perform numerical solutions of Eq. (1) for a model magnetopause equilibrium based on the observations of a southward IMF crossing by *Song et al.* [1993b]. We choose the incoming compressional wave frequency satisfying the Alfvén field line resonance condition at $x_0 = 5\ \rho_i$. We also choose $k_\parallel = 0$ to be at $x = 0$. The gradient scale length of the magnetopause width is approximately 10 ρ_i.

Figure 2 shows the numerical solutions of Eq. (1) in the outer region. The top two panels show the outer MHD solution for the electric field, $\delta E_x = -d\phi/dx$ which satisfies the boundary condition of an incoming compressional Alfvén wave. The solution behaves as $1/x$ near the Alfvén resonance location. In the third panel we plot the KAW outer solution without including diamagnetic and magnetic drift effects. The dashed and dotted lines correspond to the real and imaginary parts of the KAW solution. Notice that as the KAW propagates from the Alfvén resonance location ($x = 5\rho_i$) toward the $k_\parallel = 0$ location ($x = 0$) its amplitude is about a factor of 40 larger than the outer MHD solution because the perpendicular phase velocity is much smaller than that of the incoming MHD wave. At the electron Landau damping point, where $\omega = k_\parallel v_e$, the wave power is absorbed by strong electron Landau damping and the wave amplitude is negligible at $k_\parallel = 0$. However, with realistic magnetic drift effects appreciable KAW amplitude is found near the $k_\parallel = 0$ location as shown in the bottom panel. The KAW propagates to the $k_\parallel = 0$ location because the magnetic drift shifts the resonance and reduces electron Landau damping. When broadband compressional waves propagate from the magnetosheath to the magnetopause, broadband transverse KAWs will be excited and distributed over the entire magnetopause with enhanced amplitude. The above numerical wave solutions are consistent with the observed wave features in the magnetosheath and magnetopause: transverse waves at the magnetopause are broadband and are strongly enhanced (by a factor of 10) over the level of the compressional waves in the magnetosheath.

Transport due to Kinetic Alfvén Waves

Based on the quasilinear theory, KAWs can cause substantial transport. From the gyrokinetic equation, the density evolves approximately according to

$$\frac{\partial n}{\partial t} = \frac{\partial}{\partial x} D_\perp \frac{\partial n}{\partial x}. \tag{8}$$

This quasilinear transport equation results from decoupling of the plasma from field lines due to both the parallel electric field as well as the perpendicular electric field that couples to the particle magnetic drift motion.

The diffusion coefficient including magnetic drift effects and full Larmor radius effects can be obtained assuming a Maxwellian distribution

$$D_\perp^{(1)} \sim \sum_k \sqrt{\pi}\frac{v_A^2}{\omega}\left|\frac{\delta B_{xk}}{B_o}\right|^2 \times$$
$$\left(\eta^2 \mathscr{I}_{00}^+ + \alpha^2(1+\eta^2)\mathscr{I}_{40}^+ - 2\alpha\eta(1+\eta)\mathscr{I}_{20}^-\right), \tag{9}$$

where $E_\parallel = \mathbf{b}\cdot\nabla(\eta\phi)$ defines η, $\alpha = (|k_y|v_i/\omega)(\rho_i/L_B)$, and

$$\mathscr{I}_{nl}^\pm = 2\zeta\int_0^\infty dx x^{n+1} J_l^2(\lambda x)\exp(-x^2)\times$$
$$\left[\exp(-\zeta^2(1-\alpha x^2)^2) \pm \exp(-\zeta^2(1+\alpha x^2)^2)\right]. \tag{10}$$

For kinetic Alfvén waves, the parameters take the specific form $\eta \approx (T_e/T_i)k_\perp^2\rho_i^2/(1+k_\perp^2\rho_i^2)$, $\lambda = \sqrt{2}k_\perp\rho_i$, $\zeta = \omega/(\sqrt{2}k_\parallel v_i)$, $\alpha \approx (k_y V_A/\omega)\sqrt{\beta_i/2}(\rho_i/L_B)$, and the dispersion relation $\omega^2 \approx k_\parallel^2 v_A^2(1+\delta k_\perp^2\rho_i^2)$, with $\delta \approx 3/4 + T_e/T_i$ for $k_\perp\rho_i \ll 1$ and $\delta \approx 1 + T_e/T_i$ for $k_\perp\rho_i \gg 1$. Note that as $\zeta \gg 1$, $\mathscr{I}_{nl}^\pm \sim 1/k_\perp\rho_i$, and thus $D_\perp^{(1)}$ is proportional to $1/k_\perp\rho_i$ for KAWs. Physically, the first term of the diffusion coefficient arises from the parallel electric field [*Hasegawa and Mima*, 1978] and peaks at $k_\perp\rho_i \sim 1$. The second term in the diffusion coefficient results from magnetic drift effects and has not been described previously. In comparison, the second term is typically at least an order of magnitude larger than the first term over the entire range of $k_\perp\rho_i$. This is because at the magnetopause $T_e/T_i \ll 1$ and the drift factor $\alpha \sim 1$ for a high β_i plasma. The third term results from the cross coupling between the parallel electric field and perpendicular drift effects. The parallel magnetic field also contributes to the diffusion giving a contribution

$$D_\perp^{(2)} \sim \sum_k \sqrt{\pi}\beta_i \left|\frac{k_y}{k_\perp}\right|^2 \frac{v_A^2}{\omega} \left|\frac{\delta B_{\parallel k}}{B_o}\right|^2 \mathscr{I}_{21}^+, \tag{11}$$

which is much smaller than $D_\perp^{(1)}$ for the entire range of $k_\perp \rho_i$. Using typical magnetopause parameters, $T_e/T_i \sim 0.2$, $k_x^2\rho_i^2 \sim 5$, $\omega = 100$mHz, $\delta B_x/B_o = 0.2$, and $v_A = 500$ km/s, and replacing the summation over the perturbed magnetic field with an average value, we find that the diffusion coefficient is $D_\perp^{(1)} \sim 10^9$ m^2/s.

For southward IMF conditions, plasma transport can be further enhanced over the quasilinear prediction by magnetic reconnection resulting from KAWs. Reconnection can result from magnetic perturbations with finite radial amplitude at the location where $\mathbf{k}\cdot\mathbf{B}_o = k_z B_{zo}(x) + k_y B_{yo}(x) = 0$, $\mathbf{k}$ is the wave vector. For a radial magnetic field perturbation with $\delta B_x = \delta\psi(x)\cos(k_y y + k_z z - \omega t)$, a magnetic island appears in the $x - s$ plane with a width $\Delta x \simeq (L_s \delta B_x / k_y B_o)^{1/2}$, where $s = y + k_z z/k_y$, $L_s = | d\ln(\mathbf{k}\cdot\mathbf{B}_o)/dx |^{-1}$ is the magnetic shear scale length. Because there is a broadband spectrum of compressional waves in the magnetosheath, KAWs are present at different $k_\parallel = 0$ locations where magnetic islands form. Particle orbit islands, corresponding to the magnetic islands, will form in the particle phase space. When the KAW fluctuation level is above a threshold value, particle orbit islands overlap, particle orbits are stochastic, and particle transport can be enhanced over the quasi-linear level [*Cheng et al.*, 1994, and references therein].

Summary

In summary, we have presented a physical process of generating transverse Alfvén waves at the magnetopause: compressional waves propagating from the magnetosheath can mode convert into transverse KAWs at the magnetopause with fluctuation level enhanced over the incoming compressional wave level. We have shown that the radial structure of KAWs depends sensitively on the magnetic shear at the magnetopause. Moreover, quasi-linear theory predicts that KAWs can cause plasma transport with a diffusion coefficient $D_\perp \sim 10^9$ m^2/s and a plasma convection on the order of 1 km/s. However, for southward IMF additional transport can occur because KAWs form magnetic islands at $k_\parallel = 0$ locations. Due to the broadband nature of KAWs these islands can overlap leading to stochastic transport which is much larger than that due to quasilinear effects. Massive stochastic particle transport can lead to local flattening of the plasma density profile where islands overlap.

We emphasize that there are three particular features of wave observations in the magnetosheath and magnetopause that support the KAW scenario. First, most of the wave power in the magnetosheath and near the magnetopause is at low frequencies (10–500 mHz). Second, large gradients in the background plasma density, magnetic field, and pressure are observed at the magnetopause so that the Alfvén velocity can vary by up to a factor of 10. Third, large peaks in the $\delta B_\perp$ component are observed by the high resolution instrument during magnetopause crossings while the analytical result predicts the KAW wave amplitude to be enhanced by an order of magnitude.

Finally, we conclude by pointing out that quantitative studies of plasma transport due to KAWs are yet to be performed. In addition, plasma transport due to competing mechanisms such as resistive magnetic reconnection (or merging) has not been performed and must also be addressed before a conclusion can be drawn on the transport processes at the magnetopause.

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