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Nonlinear Interaction of Energetic Ring Current Protons with Magnetospheric Hydromagnetic Waves

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Abstract

In order to study nonlinear wave-particle interactions in the Earth's magnetosphere we have derived Hamiltonian equations for the gyrophase-averaged nonrelativistic motion of charged particles in a perturbed dipole magnetic field. We assume low frequency (less than the proton gyrofrequency) fully electromagnetic perturbations, and we retain finite Larmor radius effects. Analytic and numerical results for the stochastic threshold of energetic protons ($\gtrsim 100$ keV) in compressional geomagnetic pulsations in the Pc 5 range of frequencies (150-600 seconds) are presented. These protons undergo a

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drift-bounce resonance with the Pc 5 waves which breaks the second (longitudinal) and third (flux) adiabatic invariants, while the first invariant (the magnetic moment) and the proton energy are approximately conserved. The proton motion in the observed spectrum of waves is found to be strongly diffusive, due to the overlap of neighboring primary resonances.

I. Introduction

For some time it has been recognized that magnetic drift and bounce resonances with ring current ions may be an important mechanism for the excitation of magnetospheric hydromagnetic waves [Southwood *et al.*, 1969]. Specifically, it has recently been shown [Chen and Hasegawa, 1988] that adiabatically injected energetic protons (≥ 100 keV) can destabilize these waves by the magnetic drift-bounce resonance $\omega - m\bar{\omega}_d = p\omega_b$. Here ω_b is the bounce frequency, $\bar{\omega}_d$ is the bounce-averaged magnetic drift frequency, positive westward for ions, m is the azimuthal (east-west) wave number, p is a negative integer, and ω ($\ll \omega_b, m\bar{\omega}_d$) is the wave frequency. The predicted instability properties are consistent with satellite observations [Takahashi *et al.*, 1985; Takahashi *et al.*, 1987; Lin *et al.*, 1988] of day-side compressional Pc 5 pulsations. We note that the drift-bounce resonance condition corresponds to $k_\perp \rho = O(1)$, where $k_\perp$ is the azimuthal wave number and ρ is the equatorial gyroradius; i.e., finite Larmor radius (FLR) effects must be considered for the drift-bounce resonance.

In this paper we investigate the effect of the compressional Pc 5 waves

on the drift-bounce resonant protons. The drift-bounce resonance breaks the second (longitudinal) and third (flux) adiabatic invariants, while the first invariant (the magnetic moment) and the particle energy are approximately conserved. In particular, using the paradigm of resonance overlap [Chirikov, 1979; Lichtenberg and Lieberman, 1983], we wish to determine whether nonlinear interaction with the waves results in large scale chaotic motion. This is an important issue because large scale radial diffusion will alter the spatial gradients which are assumed to provide the free energy for the instability, thereby contributing to nonlinear saturation of the wave growth. Also, the corresponding pitch angle diffusion may enhance particle precipitation.

Early studies of the radial diffusion of radiation belt particles (see Schulz and Lanzerotti [1974] for a review) considered the breaking of the third adiabatic invariant (with the first and second invariants conserved) by large-scale ($m \simeq 1$), random impulses in the geomagnetic field or the convection electric field. Dungey [1964] and Cornwall [1966] were among the first to consider radial diffusion due to bounce resonances ($\omega \simeq \omega_b$) with hydromagnetic waves. In this case the second and third invariants are broken, the first invariant is constant, and energy is not conserved. Calculations of radial diffusion coefficients for these processes used a random phase approximation which was justified by appealing to the random nature of the impulses (or waves), or to a phase-averaging due to the finite bandwidth of measuring instruments. In this work we use the example of the drift-bounce resonance to point out that the random phase approximation can be justified by the intrinsic stochasticity (i.e., chaos) of the nonlinear wave-particle interaction.

More recent studies of the motion of charged particles in low frequency electromagnetic waves [Hasegawa, 1979; Kivelson and Southwood, 1985] have emphasized the prediction of particle fluxes expected at spacecraft locations and used linear approximations to calculate changes in the particle distribution function due to the waves. Except for Hasegawa [1979], these authors chose to neglect FLR effects.

Our approach is to derive gyrophase-averaged, Hamiltonian test particle equations of motion which we then solve numerically. In this way we take into account the *nonlinear* interaction of the particles with a spectrum of low frequency fully electromagnetic waves in a magnetic field of general geometry. The equations we derive are significant in their own right because they extend the nonrelativistic guiding center theory to include FLR effects. The Hamiltonian formulation provides conserved quantities (even for time-dependent systems) which are essential for making Poincaré surface of section plots [Lichtenberg and Lieberman, 1983] and are very useful for checking the accuracy of numerical calculations.

II. Hamiltonian Gyrocenter Equations of Motion

In this section we outline the derivation of the equations of motion, details of the calculations will be given in a future publication.

We consider the motion of a charged particle in a dipole magnetic field $B_0(\mathbf{x})$ subject to electromagnetic perturbations $\delta\mathbf{B}(\mathbf{x}, t)$ and $\delta\mathbf{E}(\mathbf{x}, t)$. Based on satellite observations, we adopt the following orderings in our perturbation

analysis :

$$\frac{\rho}{L_0} = O(\epsilon_0), \quad \frac{\delta B}{B_0} \simeq \frac{\delta v_E}{v_t} \simeq \frac{k_{\parallel}}{k_{\perp}} \simeq \frac{\omega}{\omega_c} = O(\epsilon), \quad k_{\perp} \rho = O(1). \quad (1)$$

Here $L_0 = |\nabla \ln B_0|^{-1}$ is the scale length of the unperturbed magnetic field, δv_E is the magnitude of the $\delta \mathbf{E} \times \mathbf{B}_0$ drift velocity, $v_t = \sqrt{2K/m_p}$ (where K and m_p are the particle kinetic energy and mass), $\omega_c = eB_0/m_p c$, and $k_{\parallel}$, $k_{\perp}$, and ω are characteristic parallel and perpendicular wave numbers and frequency of the perturbation. ϵ_0 is the usual guiding center small parameter and ϵ orders the perturbation quantities. We assume $\epsilon_0 \ll \epsilon$; for example, $\epsilon_0 \simeq 10^{-2}$ and $\epsilon \simeq 10^{-1}$ for 100 keV protons at $L = 6.62$ in Pc 5 waves. Note that we retain FLR effects by allowing $k_{\perp} \rho = O(1)$.

With these orderings the effect of the waves can be treated as a perturbation of the guiding center motion in $\mathbf{B}_0$, as follows: The unperturbed motion is described by the phase-space Lagrangian for guiding center motion in the dipole magnetic field [Littlejohn, 1983]. Adding the perturbations $\delta \mathbf{B} = \nabla \times \delta \mathbf{A}$ and $\delta \mathbf{E} = -\nabla \delta \Phi - (1/c)\partial \delta \mathbf{A}/\partial t$ introduces gyrophase dependence. We remove the gyrophase dependence by a Lie transformation from the guiding center coordinates to new coordinates which we call the gyrocenter coordinates [Brizard, 1989]. They differ from the guiding center coordinates by terms of $O(\epsilon)$ which are purely oscillatory in gyrophase. We assume an eikonal form for the perturbation quantities and adopt magnetic coordinates (χ, ψ, ζ) defined by $\mathbf{B}_0 = \nabla \chi = \nabla \psi \times \nabla \zeta$. (In terms of spherical polar coordinates (r, θ, ϕ) : $\chi = M \cos \theta / r^2$, $\psi = M \sin^2 \theta / r$, and $\zeta = -\phi$, where $M = B_E R_E^2$ with $B_E = 0.31$ gauss and $R_E = 6380$ km. χ is related to s , the distance along a field line, by $d\chi = B_0 ds$, and ψ, ζ are the Euler po-

tentials [Stern, 1970].) The transformed Lagrangian is called the gyrocenter phase-space Lagrangian

$$\mathcal{L} = \left(\frac{1}{\epsilon_0} \dot{\psi} + \epsilon J_0 \gamma \right) \dot{\zeta} + (\bar{p}_{\parallel} + \epsilon J_0 \alpha) \dot{\chi} + \epsilon J_0 \beta \dot{\psi} + \epsilon_0 \bar{\mu} \dot{\xi} - \mathcal{H} \quad (2)$$

$$\mathcal{H} = \frac{1}{2} \bar{p}_{\parallel}^2 B_0^2 + \bar{\mu} B_0 + \epsilon (J_0 \delta \Phi + \frac{2J_1}{k_{\perp} \rho} \bar{\mu} \delta B_{\parallel}). \quad (3)$$

Here $\bar{p}_{\parallel}$ and $\bar{\mu}$ are the gyrocenter coordinates corresponding to the guiding center coordinates $\rho_{\parallel} = v_{\parallel}/B_0$ and $\mu = v_{\perp}^2/2B_0$, where $v_{\parallel}$ and $v_{\perp}$ are the parallel and perpendicular components of the guiding center velocity and $\bar{\xi}$ is the new gyrophase. J_0 and J_1 are Bessel functions with argument $k_{\perp} \rho$ which show the reduction in the wave amplitudes due to the FLR effect. In deriving Eqs. (2) and (3), we chose the vector potentials $\mathbf{A}_0 = \psi \nabla \zeta$, $\delta \mathbf{A} = \alpha \nabla \chi + \beta \nabla \psi + \gamma \nabla \zeta$, and we defined $\delta B_{\parallel} = \dot{\mathbf{b}} \cdot \delta \mathbf{B}$, where $\dot{\mathbf{b}} = \mathbf{B}_0/B_0$.

$\mathcal{L}$ is written in dimensionless form by normalizing lengths by R_E , magnetic field quantities by B_E , mass by the proton mass m_p , and time by the inverse gyrofrequency $\omega_{ce}^{-1} = m_p c / e B_E$. In these units $\psi = 1/L$, where L is the dipole L-shell parameter.

The Euler-Lagrange equations for $\mathcal{L}$ yield the desired equations of motion. By construction $\mathcal{L}$ is independent of $\bar{\xi}$, so $\dot{\bar{\mu}} = 0$ and $\bar{\mu}$ may be treated as a constant in the equations for $\bar{p}_{\parallel}$, $\dot{\chi}$, $\dot{\psi}$, and $\dot{\zeta}$. To lowest order in ϵ_0 , $(\bar{p}_{\parallel}, \chi, \psi, \zeta)$ are canonical coordinates for the unperturbed (guiding center) motion [White and Chance, 1984]. For later use we note that, to first order in ϵ :

$$\dot{\psi} = \epsilon (J_0 r \bar{p}_{\parallel} B_0 \delta B_{\psi} - \frac{2J_1}{k_{\perp} \rho} \bar{\mu} \frac{\partial \delta B_{\parallel}}{\partial \zeta} + J_0 r \delta E_{\zeta}), \quad (4)$$

where δB_{ψ} and δE_{ζ} are components in the ψ and ζ directions.

III. Magnetic Drift-Bounce Resonance in Compressional Pc 5 Waves

Based on spacecraft measurements of compressional Pc 5 waves at synchronous orbit, near the equator (latitudes $\lesssim 15^\circ$) [Takahashi *et al.*, 1985; Takahashi *et al.*, 1987], we use the following model for the compressional component:

$$\delta B_{\parallel} = \sum_{m=m_0-\sigma_m}^{m_0+\sigma_m} \beta_m \sin k_{\parallel} s \cos(m\zeta - \omega t). \quad (5)$$

Here s is the length along the field line measured from the equator. This model describes a spectrum of westward travelling compressional waves with an antisymmetric standing-wave structure about the equator. The azimuthal mode numbers are centered at $m = m_0$ with a spread given by σ_m [Takahashi *et al.*, 1985, Figure 8]. We use the following typical values: $m_0 = 70$, $\sigma_m = 10$, $\beta_m = \beta_0/2\sigma_m$, $\beta_0 = 10$ nT, $k_{\parallel} = 2\pi/3R_E$, and $\omega = 2\pi/(500 \text{ sec})$. From radial measurements [Lin *et al.*, 1988] we assume that the wave amplitudes are constant for $6 \lesssim L \lesssim 8$.

Since the wave frequency ω is small compared to ω_b and $m\bar{\omega}_d$, the particle essentially sees a static wave and we let $\omega = 0$ in the following analysis. Also, to isolate the effect of the compressional component, we set the wave electric field components and the transverse magnetic components to zero. We describe the effect of including nonzero wave frequency and other wave components at the end of this section. The effect of corotation and convection electric fields is negligible for energetic protons [Schulz and Lanzerotti, 1974].

Firstly, consider the motion of a drift-bounce resonant proton in one

mode; i.e., one term in the sum (5).

Through its s dependence, each mode contains an infinite number of bounce harmonics corresponding to $p = -1, -3, \dots$ in $m\bar{\omega}_d = -p\omega_b$. The $p = -1$ resonance is the most important, by the following reasoning: Consider the drift-bounce resonance condition $m\bar{\omega}_d = -p\omega_b$. Schulz and Lanzerotti [1974] give expressions for ω_b and $\bar{\omega}_d$ as functions of energy K , the sine of the equatorial pitch angle $y = \sin\eta_e$, and L . For $m = 70$ and $p = -1$ the resonance condition is satisfied by a 371 keV proton at $L = 6.62$ with equatorial pitch angle $\eta_e = 70^\circ$ (for this particle the unperturbed mirror point latitude is $\alpha_M = 10^\circ$). Now, since the proton gyrocenter motion conserves $\bar{\mu}$ and K , substituting $y = \sqrt{\bar{\mu}/KL^3}$ into $m\bar{\omega}_d(K, y, L) = -p\omega_b(K, y, L)$ gives a polynomial equation in L which can be solved to obtain the resonant L value for a given m and $p \leq -3$. For $m = 70$ and $p \leq -3$ we find that the resonance condition requires $L \geq 10.7$, where the earth's magnetic field is no longer well approximated by a dipole and where regular Pc 5 waves are not observed [Lin *et al.*, 1988].

We have solved the test particle equations of motion numerically using a fifth order Runge-Kutta algorithm. In Figure 1 we plot ψ against $\zeta \pmod{2\pi}/2\pi$ whenever the gyrocenter trajectory crosses the equatorial plane travelling northward. The initial conditions (plotted as crosses) correspond to particles with the same $\bar{\mu}$ and energy as the $m = 70$, $p = -1$ resonant proton mentioned above. Thus, the chain of islands in Figure 1 corresponds to the primary resonance $70\bar{\omega}_d = \omega_b$.

Using secular perturbation theory [Lichtenberg and Lieberman, 1983], we

find the island half-width for a single m, p primary resonance is given approximately by

$$\Delta\psi_{m,p} = 2\sqrt{\frac{2J_1(k_\perp\rho)\beta_m J_p(k_\parallel s_M)}{3k_\perp\rho\psi_{m,p}}}. \quad (6)$$

Here s_M is the unperturbed mirror point distance and the Bessel function $J_p(k_\parallel s_M)$ shows the reduction in the particle's response due to averaging over parallel wavelengths. All quantities on the right-hand side of (6) are to be evaluated at the m, p resonance. As discussed above, the resonant ψ surface for a given m and p , $\psi_{m,p} \approx 1/L_{m,p}$, is given by the solution of the drift-bounce resonance condition assuming constant $\bar{\mu}$ and energy. However, for $m \approx m_0$, with $m_0 = O(100)$, and $p = -1$, we can ignore the effect of changing pitch angle and $\psi_{m,-1}$ is given approximately by

$$\psi_{m,-1} = 6.67 \times 10^{-3} \sqrt{m\sqrt{K^*}Q(y_0)}, \quad (7)$$

where K^* is the kinetic energy in keV, and y_0 is the pitch angle sine corresponding to the $m = m_0$, $p = -1$ resonance, with $Q(y_0) = 0.455 - 0.084y_0^2$. The resonance width and location predicted by Eqs. (6) and (7) are in good agreement with the numerical results, as shown in Figure 1.

Now consider the effect of adding a second mode. Using Eqs. (6) and (7), and solving for the wave amplitude, we find the Chirikov condition for the overlap of the $m = m_0$ resonance with a neighboring resonance $m' = m_0 \pm 1$ is

$$\beta_0 > (3.1 \times 10^4) \frac{3\sigma_m k_\perp \rho}{4J_1(k_\perp \rho)J_1(k_\parallel s_M)} \psi_m \psi_{m'} (\sqrt{\psi_m} - \sqrt{\psi_{m'}})^2, \quad (8)$$

where 3.1×10^4 is a numerical factor which converts the amplitude to nanoTesla and the p subscript has been dropped with the understanding that

$p = -1$. For $m = 70$ and $m' = 71$, Eqs. (7) and (8) give $\psi_m = 1/6.62$, $\psi_{m'} = 1/6.57$, and $\beta_0 > 0.05$ nT. Compared with the observed value of $\beta_0 = 10$ nT, we conclude that the Chirikov condition is easily satisfied. Figure 2 shows the resulting chaotic motion in the two modes $m = 70$ and $m' = 71$ for $\beta_0 = 10$ nT. A proton with initial coordinates in the stochastic region can diffuse in L over $\sim 0.53R_E$ and between latitudes of $\pm 13^\circ$.

We now briefly describe the effect of other wave components and nonzero wave frequency on the resonant radial motion. From Eq. (4) we see that only δB_ψ and δE_ζ contribute directly to the radial motion. Because δB_ψ has even symmetry about the equator, resulting in resonance with the $p = -2, -4, \dots$ harmonics, it will produce a small effect compared to the $p = -1$ resonance. If we estimate the magnitude of δE_ζ from Faraday's law (together with the ideal MHD approximation $\delta E_\parallel = 0$, the observation that $k_\psi \ll k_\zeta$ [Takahashi *et al.*, 1985], and the measured values of ω , $k_\parallel$, and δB_ψ) and compare the $\delta B_\parallel$ term with the δE_ζ term in Eq. (4), we find that δE_ζ is antisymmetric about the equator and it will produce resonant radial motion of the same order of magnitude as $\delta B_\parallel$. The remaining wave components (δB_ζ and δE_ψ) can be estimated from combining satellite measurements with $\nabla \cdot \delta \mathbf{B} = 0$ and Faraday's law; they are expected to make small, nonresonant contributions to the radial motion. Finally, when ω is nonzero the resonant ψ surface will shift slightly (from $L = 6.62$ to $L = 6.73$ for $\omega = 2\pi/(500\text{sec})$, for 371 keV protons), energy will no longer be exactly conserved, and Arnold diffusion [Chirikov, 1979] will take place. We note that, because the Arnold diffusion is slow compared to diffusion in the Chirikov overlap regime, the latter is

expected to dominate.

IV. Conclusion

In summary, we have derived gyrophase-averaged Hamiltonian equations of motion for charged particles in a dipole magnetic field subject to low frequency (less than the proton gyrofrequency) electromagnetic perturbations, including finite Larmor radius effects. These equations have been applied to investigate the effect of compressional Pc 5 waves on energetic ($\gtrsim 100$ keV) ring current protons. We find that, for drift-bounce resonant protons in the measured spectrum of Pc 5 waves, the Chirikov overlap criterion is greatly exceeded. This leads to pitch angle diffusion and large scale radial diffusion, and possibly provides a nonlinear saturation mechanism for the growing waves.

A proper treatment of the wave saturation involves a self-consistent calculation of the wave-particle system. The fact that the Chirikov criterion is greatly exceeded suggests that quasilinear theory is applicable [Kaufman, 1972]. However, we point out that the range of the diffusion and the quantitative effect of the diffusion on the wave saturation will depend on the exact radial structure of the waves, which is not known. Future work will focus on these issues.

Finally, the methods we present here should be useful in the interpretation of spacecraft particle flux measurements. Simulations for a spectrum of energies and a range of pitch angles can give the expected test particle

response to a given geomagnetic pulsation. We emphasize that, although the results in this paper are appropriate for eikonal perturbations on a static dipole magnetic field, our formulation is not restricted to these assumptions.

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Figures

FIG. 1. Poincaré surface of section plot for protons in a wave consisting of just the $m = 70$ term of the sum (5). The resonance width and location predicted by Eqs. (6) and (7) are shown to the right.

FIG. 2. Same as Fig. 1 except $m = 70$ and $m = 71$ modes are present together.

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Fig. 1

Fig. 2

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