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KINETIC THEORY OF GEOMAGNETIC PULSATIONS 2.
ION FLUX MODULATIONS BY TRANSVERSE WAVES

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**Kinetic Theory of Geomagnetic Pulsations 2.
Ion Flux Modulations by Transverse Waves**

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Abstract

Ion flux modulations by ultra-low-frequency radially polarized geomagnetic pulsations are examined theoretically based on the gyrokinetic analysis of Chen and Hasegawa [1991]. The theoretical results thus contain important effects such as plasma anisotropy and inhomogeneities, finite Larmor radii, realistic magnetic field, magnetic trapping, and wave mode structures. The predicted properties are consistent with the satellite observations [Takahashi et al., 1990] and further support the drift-Alfvén ballooning mode as a primary instability candidate. Our analysis, furthermore, demonstrates that, in the case of highly energetic ions, it is crucial to include the finite-Larmor-radius effects self-consistently in order to properly analyze and compare with the satellite observations.

I. Introduction

Studies of particle flux modulations associated with ultra-low-frequency (ULF) geomagnetic pulsations can provide valuable information on the oscillation mode structures and possible wave excitation mechanism(s). Recently, Takahashi et al. [1990] have carried out a detailed study of an ion flux pul-

sation event associated with a radially polarized transverse Pc 5 geomagnetic pulsation (period ~ 200 s) near the magnetic equator, at a radial distance of $5.7 - 7.0 R_E$, and near 1300 magnetic local time. The observations are consistent with a second-harmonic antisymmetric standing wave with westward propagation and a large azimuthal wave number, $|m| \sim 110$. These wave properties further support the theoretical picture that the pulsation event corresponds to a low- β drift-Alfvén ballooning mode (DABM) destabilized by the inward pressure gradient of energetic ions energy ~ 100 keV and the drift-bounce resonance ($|\omega_b| \sim |\bar{\omega}_d|$) [Southwood 1976; Chen and Hasegawa, 1988, 1991]. Here, $\beta = 8\pi P/B^2$ is the ratio of the plasma pressure to the magnetic pressure, ω_b and $\bar{\omega}_d$ are, respectively, the bounce and the bounce-averaged magnetic drift frequency. Note that $\bar{\omega}_d$ is proportional to m .

Theories for analyzing particle flux modulations by ULF waves were pioneered by Southwood and Kivelson in an extensive series of papers [Southwood and Kivelson, 1981, 1982; Kivelson and Southwood, 1983, 1985]. While the theoretical results contain effects of realistic plasma equilibria and wave mode structures, the drift-kinetic approximation employed in the analyses renders the validity regime to $|k_\perp \rho| < 1$; where $k_\perp$ is the perpendicular (to

$\tilde{B}$) wavenumber and ρ is the Larmor radius. More specifically, since, in the drift-kinetic approximation, charged particles are approximated to be guiding centers, finite ion-Larmor-radius (FILR) effects on the ion dynamics are ignored. For DABM instabilities with $|m| \sim O(10^2)$, the drift-kinetic approximation breakdown for energy ($\mathcal{E}$) $\gtrsim O(10^2 keV)$ near the synchronous orbit. It is thus of fundamental interest to extend the theories to regimes with $k_{\perp}\rho_i \sim O(1)$; i.e. where the FILR effects are included in a self-consistent fashion. An additional motivation is that in the observations of Takahashi et al [1991], the east-west phase lag does not show a systematic increase with energy ($60keV < \mathcal{E} < 430keV$) as expected from the simple drift-kinetic theory. On the otherhand, for $\mathcal{E} \sim O(400keV)$, we have $k_{\perp}\rho_i > 1$ and FILR effects could become crucial such that the drift-kinetic theory may not be applicable.

The purpose of this work is to examine particle flux modulations by ULF waves including the self-consistent FILR effects. We shall adopt the theoretical results developed by Chen and Hasegawa [1991; hereafter referred as Paper 1] based on the gyrokinetic equations [Antonsen and Lane, 1980; Catto *et al.*, 1981]. In order to test specific theoretical predictions against known

observations, we shall limit this present work to radially polarized transverse waves and observations near the geomagnetic equator. As noted in Paper 1, finite- β effects introduce, self-consistently, compressional magnetic perturbations. Thus, again, we emphasize that the present analysis is applicable to low- β plasmas during geomagnetically quiet periods. Considerations on high- β plasmas will be given in a future publications.

II. Theoretical Model and Wave Structures

As in Paper I, we consider a magnetospheric plasma consisting of a core (C) and energetic (E) components. In the low- β regime of interest here, the equilibrium dipole magnetic field is given as $\underline{B} = \underline{\nabla}\psi \times \underline{\nabla}\phi$ with ϕ being the azimuthal (east-west) symmetry coordinate. Consistent with observations of radially polarized transverse waves with $|m| \sim O(10^2)$, we shall adopt WKB ansatz in the perpendicular (to $\underline{B}$) direction; i.e.,

$$\delta\hat{\psi}(\underline{x}, t) = \delta\psi(l) \exp[i(\int^{\underline{x}} \underline{k}_{\perp} \cdot d\underline{x}_{\perp} - \omega t)]; \quad (1)$$

where l is the coordinate along $\underline{B}$, $\underline{k}_{\perp} \simeq m\underline{\nabla}\phi$, $c\partial\delta\psi/\partial l = i\omega\delta A_{\parallel}$ and $\delta A_{\parallel}$ is the vector potential (c.f. Paper 1). We have, furthermore, $\delta E_{\parallel} \simeq 0$ and

$\delta B_{\parallel} \simeq 0$ due to, respectively, the ideal magnetohydrodynamic (MHD) and low- β conditions.

In the low- β limit, $\delta\psi$ satisfies approximately the following eigenmode equation along $\underline{B}$ [Chen and Hasegawa, 1988; Paper 1];

$$\left[B \frac{\partial}{\partial l} \frac{k_{\perp}^2}{B} \frac{\partial}{\partial l} + \frac{\omega^2 k_{\perp}^2}{V_A^2} \right] \delta\psi(l) = 0. \quad (2)$$

Note, in Eq. (2), all the coefficients in principle depend on l . Assuming that $\delta\psi$ has an antisymmetric second-harmonic (odd) mode structure, we can further approximate $\delta\psi$ as

$$\delta\psi(l) = \psi_0 \sin \theta(l) \quad (3)$$

where $\theta(l) = \int_0^l k_{\parallel}(l') dl'$ and $k_{\parallel} \simeq \omega/V_A$. The boundary conditions of perfectly conducting ionospheres then dictate

$$\theta(l_0) = \pi \quad (4)$$

with l_0 being the location of the northern ionosphere.

The radial component of the wave magnetic field, δB_r , is then given by

$$\delta B_r = -\frac{ck_{\phi}}{\omega} \frac{\partial \delta\psi}{\partial l} = -\left(\frac{c}{\omega} k_{\phi} \psi_0 k_{\parallel}\right) \cos \theta(l). \quad (5)$$

Here, we note that δB_r is even in l and $k_\phi > 0$ for westward propagating wave. Meanwhile, $|\delta B_\phi/\delta B_r| \sim |k_\psi/k_\phi| \ll 1$ and, hence, the wave is indeed predominantly radially polarized standing transverse MHD shear Alfvén wave.

III. Ion-Flux Modulations at the Equator

Since the differential particle flux is in phase with the particle distribution function [Parks, 1991], it is then sufficient for the present work to examine the relative phase of the perturbed ion distribution function $\delta \hat{F}$ with respect to $\delta \hat{B}_r$. Adopting the same WKB ansatz for $\delta \hat{F}$; i.e.,

$$\delta \hat{F}(\mathbf{x}, v, t) = \delta F(l, v) \exp[i(\int^{\mathbf{x}_\perp} \mathbf{k}_\perp \cdot d\mathbf{x}_\perp - \omega t)], \quad (6)$$

δF can then be shown to be (c.f. Paper 1)

$$\delta F = \delta F_a + \delta K \exp(iL_k); \quad (7)$$

where δF_a is the so called adiabatic response given by

$$\begin{aligned} \delta F_a(l, v) = \frac{q}{m} \left\{ \delta \psi \frac{\partial F_0}{\partial \mathcal{E}} - \frac{Q}{\omega} F_0 J_0 \delta \psi e^{iL_k} \right. \\ \left. + \frac{1}{B} \frac{\partial F_0}{\partial \mu} \left[\left(\delta \psi + \frac{iv_\parallel}{\omega} \frac{\partial \delta \psi}{\partial l} \right) (1 - J_0 e^{iL_k}) \right. \right. \\ \left. \left. - \frac{v_\perp J_1}{k_\perp c} \delta B_\parallel e^{iL_k} \right] \right\}, \end{aligned} \quad (8)$$

$$L_k = \tilde{k} \cdot \tilde{\varepsilon}_{\parallel} \times \tilde{v}/\Omega, \quad (9)$$

$$Qf_0 = \left(\omega \frac{\partial}{\partial \mathcal{E}} + \tilde{k} \times \frac{\tilde{\varepsilon}_{\parallel}}{\Omega} \cdot \tilde{\nabla} \right) F_0, \quad (10)$$

$\mu = v_{\perp}^2/2B$, $\tilde{\varepsilon}_{\parallel} \equiv B/B$, J_0 and J_1 are Bessel functions with argument $\lambda = k_{\perp}\rho$, and the rest of the notations is standard. Noting that $\delta B_{\parallel}$ and $\delta\psi$ are odd modes, we have, at the equator ($l = 0$), $\delta B_{\parallel}$, $\delta\psi \rightarrow 0$ and δF_a reduces to

$$\delta F_a(l = 0, \tilde{v}) \equiv \delta F_{a0} = \frac{q}{m} \frac{1}{B} \frac{\partial F_0}{\partial \mu} \frac{i v_{\parallel}}{\omega} \frac{\partial \delta\psi}{\partial l} (1 - J_0 e^{iL_k}). \quad (11)$$

It is worth noting that $\delta F_{a0} \rightarrow 0$ as $\lambda \rightarrow 0$ and, thus, δF_{a0} is **absent** in the drift-kinetic theories.

δK , meanwhile, satisfies the following gyrokinetic equation

$$\left[\hat{\sigma} |v_{\parallel}| \frac{\partial}{\partial l} - i(\omega - \omega_d) \right] \delta K = i \frac{q}{m} Q F_0 (\delta S_1 - i \hat{\sigma} \delta S_2), \quad (12)$$

where $\hat{\sigma} = \text{sgn}(v_{\parallel})$,

$$\delta S_1 = \omega_d J_0 \delta\psi / \omega, \quad (13)$$

$$\delta S_2 = (|v_{\parallel}| \delta\psi / \omega) \partial J_0 / \partial l, \quad (14)$$

and, again, we have neglected $\delta E_{\parallel}$ and $\delta B_{\parallel}$. We can readily solve Eq. (12)

and obtain formal expressions of δK (referring to Paper 1). For the present

analysis, it is sufficient to examine δK 's corresponding to various ranges in energy.

A. Near wave-particle resonances

Here, we have, for $K = \pm 1$,

$$\omega - \bar{\omega}_d - K\omega_b \simeq 0. \quad (15)$$

Thus, $K = 1$ corresponds to bounce resonances due to low energy ions; while $K = -1$ corresponds to drift-bounce resonances due to highly energetic ions. Since the left-hand-side propagator of Eq. (12) nearly vanishes for the resonant particles, $|\delta K| \gg |\delta F_a|$ and we find, at $l = 0$,

$$\delta F_0 \simeq -i\hat{\sigma}K\left(\frac{q}{m}QF_0\right)\frac{\overline{\delta S_1|s_0^l|}}{K\omega_b + \bar{\omega}_d - \omega}e^{iL_k}, \quad (16)$$

where $\bar{A}$ denotes the bounce average of A , $s_0^l = \sin I_0^l$,

$$I_0^l = \int_0^l \frac{dl}{|v_{||}|}(\omega - \omega_d), \quad (17)$$

and we note that δS_2 contribution, typically, can be neglected with respect to δS_1 .

It is worth noting the following interesting properties of Eq. (16). First, the $\hat{\sigma}$ dependency indicates flux oscillations at pitch angle α and its conjugate, $180^\circ - \alpha$, are 180° out of phase. Second, noting that both δS_1 and s_0^l are odd functions and

$$\bar{A} \equiv \int_0^{l_T} A dl / |v_{||}| / \int_0^{l_T} dl / |v_{||}|,$$

it is then obvious that for particles with 90° pitch angle, $l_T \rightarrow 0$, and, hence, $\delta F_0(\alpha = 90^\circ) = 0$. Meanwhile, particles with moderate pitch angles such that $l_T \approx l_m$ with l_m being the location at which δS_1 is peaked will tend to have maximum flux oscillations. Both properties are consistent with satellite observations [Takahashi *et al.* 1990].

As to the energy dependence of the phases of ion flux oscillations, let us consider the $\hat{\sigma} > 0$ case. From Eq. (16), the phase variations depends on the value of K . Note also that $\delta S_1 \propto \psi_0 \propto -\delta B_{r0}$. Thus, for the $K = -1$ drift-bounce resonance of energetic ions, we have

$$\delta F_0(\hat{\sigma} > 0, K = -1) \propto -i\delta B_{r0} \frac{Q F_0}{\bar{\omega}_d - \omega_b - \omega} \exp(iL_k). \quad (18)$$

Let $\omega = \omega_0 + i\Delta$ where ω_0 and Δ correspond to the frequency peak and the associated finite bandwidth, respectively. Note that Δ may be due to effects such as nonlinear resonance broadening and/or finite detector energy “band-

width" [Southwood and Kivelson, 1981]. Noting $\exp(-i\omega t)$ time dependence, the relative phase between δF_0 and δB_{r0} , Φ , is then given by

$$\Phi = \begin{cases} \pi/2 + \tilde{k} \cdot \rho & \text{for } \mathcal{E} > \mathcal{E}_{rd} \\ \tilde{k} \cdot \rho & \text{for } \mathcal{E} \approx \mathcal{E}_{rd}, \\ -\pi/2 + \tilde{k} \cdot \rho & \text{for } \mathcal{E} < \mathcal{E}_{rd}; \end{cases} \quad (19)$$

where $\mathcal{E}_{rd}$ is the resonance energy; i.e., $\bar{\omega}_d = \omega_b + \omega \simeq \omega_b$ at $\mathcal{E} = \mathcal{E}_{rd}$, and we have assumed the resonance is **destabilizing** ($QF_0 > 0$) and recalled $L_k = -\tilde{k} \cdot \rho$. For the drift-bounce resonance, we have $k_\phi \rho \approx 1$ and, hence, $\tilde{k} \cdot \rho \approx \pm 0.34 k_\phi \rho \approx \pm 20^\circ$ for the November 20, 1985 event reported by Takahashi *et al.* [1990]. Here + and - denote, respectively, eastside and westside guiding centers. Note the theoretical predictions of Eq. (19) are the same as those based on the drift-kinetic model [Takahashi *et al.*, 1990] and are consistent with satellite observations.

In the case of the $K = 1$ ($\omega \approx \omega_b$) bounce resonance of lower-energy ions, Eq. (16) indicates that the phase difference, Φ , will be similar to that given by Eq. (19) if $\mathcal{E}_{rd}$ is replaced by $\mathcal{E}_{rb}$, $|k_\perp \rho| \ll 1$, and one assumes that $QF_0 < 0$ (i.e., **stabilizing**). Since, for $\mathcal{E} \approx \mathcal{E}_{rb}$, $QF_0 \approx \omega \partial F_0 / \partial \mathcal{E}$, this will provide an interesting observational test on the stability role of these bounce-resonance ions and whether the stability role is consistent with the property of the

velocity distribution function, F_0 .

B. Highly energetic ions

As discussed in Sec. I, one of the motivations of the present work is to examine flux modulations of highly energetic ($|k_\perp \rho| > 1$) ions such that the drift-kinetic theories are definitely not applicable. For example, with $\mathcal{E} \approx 400 \text{ keV}$ and $|m| \approx 100$, we have $|k_\phi \rho| \approx 2$. Noting that $|k_\phi \rho| > 1$ implies $|\omega_d| > |\omega_b| > \omega$, it is easier to derive δK from Eq. (12) via expansions in terms of ω/ω_d and ω_b/ω_d . Letting $\delta K = \delta K_0 + \delta K_1 + \dots$ in Eq. (12), we have

$$\delta K_0 = \frac{q}{m} \frac{Q F_0}{\omega} J_0 \delta \psi, \quad (20)$$

and

$$\delta K_1 = i \frac{q}{m} \frac{Q F_0}{\omega} J_0 \frac{\hat{\sigma} |v_{||}|}{\omega_d} \frac{\partial}{\partial l} \delta \psi. \quad (21)$$

At $l = 0$, $\delta K_0 = 0$, and

$$\delta K_{10} = i \frac{q}{m} \frac{Q F_0}{\omega} J_0 \frac{\hat{\sigma} |v_{||}| k_0}{\omega_{d0}} \psi_0. \quad (22)$$

For $|k_\perp \rho| > 1$, δF_{a0} , meanwhile, becomes [c.f. Eq. (11)]

$$\delta F_{a0} \simeq i \frac{q}{m} \frac{1}{B} \frac{\partial F_0}{\partial \mu} \frac{\hat{\sigma} |v_{||}| k_0}{\omega} \psi_0. \quad (23)$$

Noting that Eq. (23) is strictly a result of large ion Larmor radii; i.e., it is absent in the usual drift-kinetic analysis.

Combining Eqs. (22) and (23) yields

$$\delta F_0 \simeq i \frac{q}{m} \frac{\hat{\sigma} |v_{\parallel}| k_0}{\omega} \psi_0 \left(\frac{1}{B} \frac{\partial F_0}{\partial \mu} + \frac{J_0 Q F_0}{\omega_d} e^{iL\kappa} \right); \quad (24)$$

where all quantities are to be evaluated at $l = 0$. That $\delta F_0 \propto \hat{\sigma} |v_{\parallel}|$ readily implies the two observational characteristics noted above; that is, (1) the oscillations at pitch angles α and $180^\circ - \alpha$ are out of phase, and (2) the flux modulations are minimal for $\alpha = 90^\circ$ ($|v_{\parallel}| \rightarrow 0$) particles. Since $\mathcal{E} |Q F_0| \sim |\omega_d F_0|$ and $|J_0| < 1$ for $|k_{\perp} \rho| > 1$, the second term in Eq. (24) is, typically, smaller than the first term, $\partial F_0 / B \partial \mu$, for **anisotropic** plasmas. That is, for $|k_{\perp} \rho| > 1$ and $|\partial F_0 / \partial \mu| \neq 0$, we have, approximately,

$$\delta F_0 \approx i \frac{q}{m} \frac{\hat{\sigma} |v_{\parallel}| k_0}{\omega} \psi_0 \frac{1}{B} \frac{\partial F_0}{\partial \mu}. \quad (25)$$

Note that Eq. (25) predicts the interesting property that there is negligible east-west phase difference. The physical explanation is that, in the short-wavelength large FILR limit, particle dynamics is decoupled from that of the “distant” guiding centers and, thus, depends only on local wave perturbations. For F_0 peaked about $\alpha = 90^\circ$ [c.f. Fig. 9 of Takahashi *et al.*, 1990], we

have $\partial F_0/\partial\mu > 0$ for moderate pitch angles (e.g., $\alpha \approx 30^\circ$ or 150°), and the phase difference Φ is given by

$$\Phi \approx \hat{\sigma} 90^\circ; \quad (26)$$

which is consistent with the satellite observations and, indeed, provides a theoretical explanation of the observation that the east-west phase lag at 430 keV is smaller than that at 59 keV.

IV. Summary and Discussion

In this work we have applied the results of gyrokinetic analysis of Paper 1 to the study of ion flux modulations by ULF geomagnetic pulsations. As an initial application, we have limited the present study to radially polarized transverse waves relevant to low- β plasmas as well as the satellite observations of Takahashi *et al.* [1990]. By examining the perturbed ion distribution function, δF , at the equator, we have found that the theoretical results exhibit properties consistent with satellite observations. It further supports the suggestion that these radially polarized transverse waves correspond to low- β drift-Alfvén ballooning modes (DABM) driven unstable by inward energetic ion pressure gradient and drift-bounce resonances. Since the gy-

rokinetic analysis contains finite ion Larmor-radius effects self-consistently. It is particularly interesting to study flux modulations of highly energetic ions. Our results indicate that, for these highly energetic ions, wavelengths become shorter than the Larmor radii such that the particle dynamics becomes decoupled from the corresponding guiding centers and responds only to **local** wave perturbations. A direct implication of the large ion Larmor-radius results in that there exists negligible east-west phase lag; which, while in contrast to the drift-kinetic predictions, is nonetheless consistent with satellite observations.

It is obviously desirable to extend the present work to the $\beta \sim 0(1)$ regime where one needs to include effects due to the finite magnetic compressional component, $\delta B_{\parallel}$. In this respect, it will be interesting to examine the finite $\delta B_{\parallel}$ effects with and without the self-consistent constraints imposed by the drift-Alfvén ballooning modes. Such studies could, perhaps, provide additional insights to the role (if any) of DABM instabilities.

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