Generation of kinetic Alfvén waves by upper-hybrid pump waves

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(Received 17 September 1997 and in revised form 22 November 1997)

The nonlinear mechanism for kinetic-Alfvén-wave (KAW) excitation by upper-hybrid waves (UHWs) is discussed. Taking into account perpendicular dispersion of KAWs, caused by effects of finite ion Larmor radius and electron inertia, we examine a new channel for UHW decay, in which a pump UHW decays into another UHW and an ultralow-frequency wave, KAW: UHW \( \rightarrow \) KAW+UHW. A nonlinear dispersion relation is derived, and the growth rate of the parametric decay instability is calculated for a pump UHW propagating at an arbitrary angle to the background magnetic field. We find that the resulting KAWs often have a two-peaked spectrum with different perpendicular dispersions. Using satellite observations, the analytical results are applied to show that the considered process represents an effective mechanism for KAW generation and the consequent spreading of the UHW spectrum in the Earth’s magnetosphere and solar corona.

1. Introduction

It follows from one-fluid magnetohydrodynamic (MHD) theory that in a cold magnetized plasma, ultralow-frequency \( \omega \ll \omega_{Bi} \) (where \( \omega_{Bi} = eB_0/m_i c \) is the ion-cyclotron frequency) Alfvén waves with a linear dispersion law can exist:

\[
\omega^2 = k_z^2 V_A^2,
\]

where \( V_A = B_0^2/4\pi n_0 m_i \) is the Alfvén speed; the \( z \) axis is chosen to be along the external magnetic field \( B_0 \). Taking into account the finite values of the Larmor ion radius or the electron inertia in the same frequency region, there also exists a much shorter-scale mode with components of the wave vector \( k_L \gg k_z \). These waves were called kinetic Alfvén waves (KAWs) by Hasegawa and Chen (1976). Though KAWs retain the main properties of magnetohydrodynamic (MHD) Alfvén waves, they have some important new properties as well, including (i) a dependence of the wave dispersion on the transverse wave vector component, and (ii) the presence of a non-zero longitudinal component of the electric field \( E_z \). Owing to these properties, KAWs can interact with plasma particles and other kinds of waves more effectively than MHD Alfvén waves.
Excitation of KAWs by linear mode conversion (Hasegawa and Chen 1976) and by nonlinear decay of lower-hybrid (Shukla and Mamedov 1978), ion-cyclotron (Patel et al. 1985) and ion-Bernstein (Sharma and Tripathi 1988) waves has been investigated. In the present paper we examine a new channel for KAW generation in a plasma in the presence of upper-hybrid waves (UHWs). It has been shown by Murtaza and Shukla (1984) that lower-hybrid waves can be generated as a result of UHW decay. Three-wave interaction including UHWs has been considered by Kotsarenko et al. (1993) and Yukhimuk and Yukhimuk (1994), where the pump wave was an ordinary electromagnetic wave. Parametric excitation of MHD (non-dispersive) Alfvén (and/or magnetoacoustic) and electron-cyclotron waves has been considered by Stenflo and Shukla (1995). Here we show that parametric decay of the UHW into another UHW and a KAW may be an effective mechanism of KAW excitation.

We apply the results of our calculations to investigate a new way of KAW generation in the magnetosphere. Investigations of parametric processes in a magnetized plasma are of particular theoretical and practical interest, and many articles have been devoted to this problem (Shukla and Mamedov 1978; Murtaza and Shukla 1984; Patel et al. 1985; Sharma and Tripathi 1988; Kotsarenko et al. 1993, 1994; Yukhimuk and Yukhimuk 1994). We show that in those magnetosphere regions where there exists a high level of UHWs, the decay UHW→UHW+KAW can take place, resulting in KAW generation.

The long-period geomagnetic pulsations Pc5, and in some cases also Pc4, have been considered as MHD Alfvén waves. Ground-based observations have allowed investigation of the properties of geomagnetic pulsations (Hasegawa 1975; Nishida 1980). The sources of the geomagnetic pulsations can be waves penetrating into the magnetosphere from the solar wind, as well as plasma instabilities inside the magnetosphere itself. The generation of geomagnetic pulsations by MHD resonances and by development of plasma instabilities (Hasegawa 1975; Nishida 1980) have been investigated. Kinetic Alfvén waves have been observed in the magnetosphere by, for example, Hughes and Southwood (1976), Gurnett et al. (1984) and Volokitin and Dubinin (1989), but from the Earth’s surface it is possible to observe only the long-scale part of their spectrum because of the screening influence of the ionosphere and atmosphere. In the work by Voitenko et al. (1990) the origin of these waves was dealt with in the presence of intense longitudinal currents and KAW current-driven instability. At the same time, as was shown by Watanabe and Oya (1993) and Mishin et al. (1989), at the precipitation sites of energetic particles, the UHWs are easily generated. Recent high-quality investigations with the help of the EXOS-D satellite have shown that the UHW is the most stable wave mode in the high-frequency region \( \omega \gg \omega_{Be} \), where \( \omega_{Be} \) is an electron-cyclotron frequency. A high level of UHWs is continuously observed at the heights over 1000 km at all latitudes from the equator to the auroral zone, and enhancements of the oscillation intensity near the magnetic equator and in the precipitation region of energetic particles have been pointed out by Oya et al. (1990).

In Section 2 the main problem is stated and the basic equations are established. In Section 3 the dispersion equation for KAWs is derived. In Section 4 the dispersion equation for scattered UHWs is obtained. In Section 5 the dispersion equation describing three-wave interaction is found and analysed. Section 6 is devoted to the application of the obtained results to the magnetosphere plasma. Section 7 contains a discussion, and Section 8 contains the conclusions.
2. Basic equations

We consider a homogeneous magnetized plasma \( B_0 = B_0 e_z \) in which UHWs propagate:

\[
E_1 = (E_{1x} e_x + E_{1z} e_z) \exp(i\psi) + \text{c.c.,} \tag{2}
\]

where the wave phase

\[
\psi_1 = -\omega_1 t + k_{1x} x + k_{1z} z, \\
\omega_1^2 = \frac{1}{2} \left[ \omega_h^2 + (\omega_h^2 - 4 \omega_{pe}^2 \cos^2 \vartheta_1) \right]^{1/2},
\]

\[
\omega_h^2 = \omega_{pe}^2 + \omega_{Be}^2, \quad \omega_{pe}^2 = \frac{4\pi n_0 e^2}{m_e}, \quad \omega_{Be} = \frac{eB_0}{mc},
\]

\( E_{1x} \) and \( E_{1z} \) are transverse and longitudinal components of the wave electric field, and \( \vartheta_1 \) is the angle between the direction of the external magnetic field \( B_0 \) and the wave vector \( k_1 \).

Let us discuss the case when the UHW pump decays into a KAW and another UHW. In this case synchronism conditions that are necessary for the effective three-wave interaction must hold:

\[
\omega_1 = \omega_2 + \omega_3, \quad k_1 = k_2 + k_3, \tag{3}
\]

where \( \omega_2, \omega_3 \) and \( k_2, k_3 \) are the frequencies and wave vectors of the UHW and KAW.

To describe three-wave interaction, we shall use a set of two-fluid MHD equations:

\[
\frac{\partial V_{\alpha}}{\partial t} = \frac{1}{m_{\alpha}} (e_{\alpha} E + F_{\alpha}) + (V_{\alpha} \times \omega_{B_{\alpha}}) - V_{T_{\alpha}}^2 \nabla n_{\alpha} / n_{\alpha}, \tag{4a}
\]

\[
\frac{\partial n_{\alpha}}{\partial t} = -\nabla \cdot (n_{\alpha} V_{\alpha}), \tag{4b}
\]

\[
\nabla \times B = \frac{4\pi}{c} j + \frac{1}{c} \frac{\partial E}{\partial t}, \tag{4c}
\]

\[
\nabla \times E = -\frac{1}{c} \frac{\partial B}{\partial t}, \tag{4d}
\]

\[
\nabla \cdot E = 4\pi \rho_e, \tag{4e}
\]

where

\[
F_{\alpha} = \frac{e_{\alpha}}{c} (V_{\alpha} \times B) - m_{\alpha} (V_{\alpha} \cdot \nabla) V_{\alpha},
\]

\[
j = e (n_i V_i - n_e V_e), \quad \rho_e = e (n_e - n_i).
\]

Let us represent the electron density, electron speed, and the electric and magnetic fields as

\[
n_e = n_0 + n_1 + n_2 + n_3, \tag{5a}
\]

\[
V = V_1 + V_2 + V_3, \tag{5b}
\]

\[
E = E_1 + E_2 + E_3, \tag{5c}
\]

\[
B = B_0 e_z + B_3, \tag{5d}
\]

where \( n_0 \) is the average equilibrium plasma density value, \( n_1 \) is the electron density perturbation due to the UHW pump, \( n_2 \) and \( n_3 \) are the density perturbations due to decay products, \( V_1 \), is the electron velocity in the pump-wave electromagnetic field.
3. Dispersion equation for KAWs

Since Alfvén waves are ultralow-frequency, in order to obtain a dispersion equation, we can use the plasma approximation:

\[ n'_e = n'_i, \]

where \( n'_e \) and \( n'_i \) are the perturbations of ion and electron number densities.

From the equations of motion and the continuity equation, we can find expressions for \( n'_e \) and \( n'_i \):

\[ \frac{n'_e}{n_0} = \frac{e}{T_e} \left( 1 - \frac{\omega_3^2}{k_{Ax}^2 V_{Te}^2} \right)^{-1} \left( \varphi_3 - A_3 + \frac{\kappa^2 \omega_3^2}{k_{Ax}^2 V_{Te}^2} \frac{F_{3x}}{ie k_{3z}} - \frac{k_{Ax} \omega_3}{ie k_{3z}} \frac{F_{3y}}{ie k_{3z}} + \frac{F_{3z}}{ie k_{3z}} \right), \]

\[ \frac{n'_i}{n_0} = -\frac{e}{T_i} \kappa \left( \varphi_3 + \frac{\omega_B^2}{\omega_3^2} \frac{k_{Ax}^2}{k_{3x}^2} A_3 \right). \]

Here \( \kappa_n = k_{Ax} \rho_n, \rho_n = V_B/\omega_{Bo}, \kappa = \kappa_i, A_3 = (\omega_3/k_{3z} c)A_{3z}, \) and \( \varphi_3 \) and \( A_3 \) are the scalar and vector potentials.

To find the relation between \( \varphi_3 \) and \( A_3 \), we use Ampère’s law

\[ \nabla^2 A_3 = \frac{4\pi e}{c} (n_{3x} V_{3x} - n_{3e} V_{3e})_z, \]

and \( z \) is the component of the equation of motion for electrons, (4).

From (4) and (9), we obtain

\[ \frac{n_{3x}}{n_0} = \frac{e}{T_e} \left[ \varphi_3 - A_3 \left( 1 + \frac{k_{Ax}^2 \delta_e^2}{c^2} \right) + \frac{F_{3y}}{ie k_{3z}} \right], \]

where \( \delta_e = c/\omega_{pe} \) is the electron inertia length.

From (7), (8) and (10), we can find a dispersion equation for KAWs:

\[ \varepsilon_3 (\omega_3, k_3) \varphi_3 = \mu_3 \varphi_1 \varphi_2^*, \]

where

\[ \varepsilon_3 (\omega_3, k_3) = \omega_3^2 - \omega_{3A}^2, \]

\[ \omega_{3A}^2 = k_{3A}^2 V_A^2 \left[ 1 + \frac{k_{Ax}^2 \rho_0^2}{1 + k_{3x}^2 \delta_e^2} \right] \]

\[ \rho_0 = \left( 1 + \frac{T_e}{T_i} \right)^{1/2} \rho_i, \quad \text{where} \quad \rho_i \text{ is the Larmor ion radius}, \]

\[ \mu_3 = \frac{e}{m_e} k_{3A}^2 V_A^2 \left[ 1 + \frac{k_{Ax}^2 \rho_0^2}{1 + k_{3x}^2 \delta_e^2} \right] \left[ -1 + (1 + \frac{k_{3x}^2 \rho_0^2}{1 + k_{3x}^2 \delta_e^2}) \frac{\omega_{Bi}}{\omega_3} \frac{\omega_{Be} k_{1x} + k_{2x}}{\omega_1 k_{3z}} \right] \frac{k_1 \cdot k_2}{\omega_1 \omega_2}. \]

In the calculation of the coupling coefficient \( \mu_3 \), we have taken account of the ponderomotive force as well as the transverse current created by the incident and scattered UHWs. And since in the lower magnetosphere the plasma parameter \( \beta < m_e/m_i \), in the dispersion equation for KAWs we also keep terms with electron-inertia effects.
4. Dispersion equation for UHWs

In order to obtain a dispersion equation for UHWs, we use the equations for the electron component of the plasma:

\[
\frac{\partial V_e}{\partial t} = \frac{1}{m_e} (eE + F_e) - \frac{V_e \times \omega_{Be}}{m_e} - \frac{T_e}{m_e} \nabla n_e, \quad \frac{\partial n_e}{\partial t} = -\nabla \cdot (n_e V_e), \quad \nabla^2 \varphi = 4\pi e n_e. \tag{12a, b, c}
\]

From the equations of motion and continuity, we find the electron density perturbation due to UHWs:

\[
\frac{n'_e}{n_0} = -\frac{e}{m_e} \left[ \frac{k_{2e}^2}{\omega_e^2 - \omega_{Be}^2} \left( \varphi_2 + i \frac{F_{2e}}{\omega_{Be}} \right) + \frac{k_{2e}^2}{\omega_e^2} \right]. \tag{13}
\]

Substituting (13) into Poisson's equation, we obtain a dispersion equation for the UHWs:

\[
\varepsilon_2(\omega_2, k_2) \varphi_2 = \mu_2 \varphi_1 \varphi_3^*, \tag{14}
\]

where

\[
\varepsilon_2(\omega_2, k_2) = \omega_2^2 - \omega_{2e}^2,
\]

\[
\omega_{2e}^2 = \frac{1}{2} \left[ \omega_h^2 + \omega_{pe}^2 \omega_{Be}^2 \cos^2 \vartheta_2 \right]^{1/2},
\]

\[
\mu_2 = \frac{e_k}{m_e} \frac{k_{3e}}{m_e} \frac{\omega_{pe}}{\omega_1} \frac{|k_1 \times k_2|}{k_2^2 k_3^2 (1 + \kappa^2)}.
\]

Here \(\vartheta_2\) is the angle between the wave vector \(k_2\) and the external magnetic field \(B_0\).

5. Nonlinear dispersion equation and decay growth rate

From (11) and (14), we can find a nonlinear dispersion expression describing three-wave interaction

\[
\varepsilon_3^*(\omega_3, k) \varepsilon_3^*(\omega_3, k) = \mu_3 \mu_2^* |\varphi_1|^2. \tag{15}
\]

In the absence of the pump wave, \(\varphi_1 = 0\), in the plasma there exist two kinds of waves, with frequencies \(\omega_3 = \omega_{LA}\) and \(\omega_2 = \omega_{2e}\). In the presence of the pump wave, \(\varphi_1 \neq 0\), energy transfer from the pump UHW to the other plasma modes occurs, and amplitudes increase with the growth rate.

Putting \(\omega_1 = \omega_{3f} + i\gamma_{NL}\) and \(\omega_2 = \omega_{2e} + i\gamma_{NL}\) \((|\gamma_{NL}| \ll \omega_{2e}, \omega_{3f})\) into (15) and expanding \(\varepsilon_3(\omega_3, k_3)\) and \(\varepsilon_2(\omega_2, k_2)\) in Taylor series, we obtain an expression for the wave growth rate:

\[
\gamma_{NL}^2 = \frac{\mu_3 \mu_2^*}{4 \omega_{2e} \omega_{LA}} |\varphi_1|^2. \tag{16}
\]

Substituting the expressions for the coupling coefficients \(\mu_3\) and \(\mu_2\) into (16), we find:

\[
\gamma_{NL} = \frac{1}{2} \left( \frac{W}{2} \right)^{1/2} \sqrt{\frac{\rho_{ei}}{\delta_e}} \left( \frac{m_e \omega_{Be}}{m_e \omega_{pe}} \right)^{1/2} \tau_{NL}(\vartheta_1, \vartheta_2) \omega_{BI}, \tag{17}
\]

where

\[
\tau_{NL}(\vartheta_1, \vartheta_2) = \left| \text{sign}(k_{3e}) \sin(2\vartheta_2 - 2\vartheta_1) (k_1 d_e \sin \vartheta_1 + k_2 d_e \sin \vartheta_2) \right|^{1/2}.
\]
\[ W = \frac{|E_0|^2}{4\pi n_0 T_e}, \]

and

\[ d_e = \left( \frac{T_e}{4\pi n_0 e^2} \right)^{1/2} \]

is the electron Debye radius.

It follows from (19) that instability growth rate depends on the parameters of the pump UHW as well as on the direction of secondary UHW propagation with respect to the direction of the external magnetic field. The wavenumber \( k_2 \), appearing in \( \gamma_{NL} \) is to be found from the frequency-matching condition.

As follows from the general three-wave interaction theory, the squared decay increment (16) may be written explicitly as a squared quantity (Stenflo 1994). Without going into complicated derivations of such an explicit form of the \( \gamma_{NL}^2 \), we only note that sign \((k_3 z) \sin(2 \vartheta_2 - 2 \vartheta_1) (k_1 d_e \sin \vartheta_1 + k_2 d_e \sin \vartheta_2) > 0 \), provided that matching conditions are satisfied, and therefore (16) can indeed be written as a squared quantity.

6. Application

Let us apply the obtained results to the magnetosphere plasma. Typical parameters for the lower magnetosphere are

\[ T_e \approx T_i \approx 1 \text{ eV}, \quad \frac{\omega_{pe}}{\omega_B} \approx 10, \quad n_0 \approx 10 \text{ cm}^{-3}, \]

\[ \omega_B / \omega_{pe} \approx 10^2 - 10^3 \text{ s}^{-1}, \quad \omega_{pe} / \omega_{pe} \approx 10^6 - 10^7 \text{ s}^{-1}, \quad V_{Ti} \approx 10^6 \text{ cm s}^{-1}. \]

Substituting these values of the plasma parameters and a UHW intensity of \( W \approx 10^{-5} \) (Oya et al. 1990; Watanabe and Oya 1993) into (17), we obtain

\[ \gamma_{NL} \approx (1-10) \tau_{NL} (\vartheta_1, \vartheta_2) \text{ s}^{-1}. \]  

(18)

Numerical calculations dependence of the maximum growth rate on the angle \( \vartheta_1 \) for \( k_1 d_e = 0.3 \), shown in Fig. 1, give sufficiently short values of the time for development of instability in a wide range of \( \vartheta_1 \):

\[ \tau \approx \gamma_{NL}^{-1} \approx (0.1-1) \text{ s}. \]  

(19)

The angular dependence of the decay growth rate and the corresponding wave-numbers of the excited KAWs are plotted in Fig. 2. Figure 3 shows the dependence of the maximum growth rate on the pump-wave wavenumbers \( k_{1x} \) and \( k_{1z} \). In all plots the regions where \( \vartheta_1 < \vartheta_2 \) are unphysical.

Taking into account that near the equator UHW emission enhancement is induced by the non-equilibrium electron velocity distribution in the transverse direction, the angle \( \vartheta_1 > \frac{1}{2} \pi \), and, as a result of parametric decay, KAW turbulence with \( k_{3x} \lambda_{De} \lesssim 0.35 \) will be generated. In the auroral regions, where UHW generation is caused by field-aligned electron beams, the angle \( \vartheta_1 < \frac{1}{2} \pi \) (Wahlund et al. 1994a,b), and parametric instability of such UHWs generates turbulence of KAWs with \( k_{3x} \lambda_{De} \lesssim 0.1 \). That is why, in the satellite system of coordinates, the typical frequencies of KAW turbulence near the geomagnetic equator and in the auroral regions can be different. Kinetic Alfvén waves generated by parametric instability will be dissipated by electron Landau resonance with consequent heating of the
Generation of kinetic Alfvén waves by upper-hybrid pump waves

0.6
0.4
0.2
0

$\theta_1$

0°
30°
60°
90°

7. Discussion

The investigations described here have shown that in a magnetized plasma with a low value of the plasma parameter $\beta$, it is possible for an upper-hybrid pump wave to decay into a kinetic Alfvén wave and another upper-hybrid wave (together with scattering of the upper-hybrid wave on the kinetic Alfvén wave). This decay channel, as far as we know, has not been discussed in the literature before. It is possibly due to finite-ion-temperature effects for Alfvén waves with $k_\perp \gg k_z$.

The physical mechanism of the instability considered here is similar to that considered by Drake et al. (1974). The nature of this mechanism is as follows. The scattered upper-hybrid wave field, added to the upper-hybrid pump field, leads to slow modulation of the field intensity and to a ponderomotive force. This ponderomotive force acts on plasma electrons and intensifies low-frequency kinetic Alfvén waves. Then the kinetic Alfvén waves modulate the pump and increase the intensity of the scattered upper-hybrid wave. So, the upper-hybrid pump, interacting with the scattered upper-hybrid wave, generates ultralow-frequency kinetic Alfvén waves with frequency $\omega_1 - \omega_2$ and wave vector $\mathbf{k}_1 - \mathbf{k}_2$. High-frequency upper-hybrid waves with frequency $\omega_1 - \omega_3$ and wave vector $\mathbf{k}_1 - \mathbf{k}_3$ will be generated in the same way.

In the present work we have considered three-wave parametric interaction for the case of a homogeneous plasma. Inhomogeneity can lead to localization of the wave interaction region due to phase mixing. That is why the waves produced by decay must be intensified on going through the interaction region or the generation source, provided that regular oscillations in the upper-hybrid resonance layer exist.

Data from EXOS-D satellite observations (Oya et al. 1990) show that a high level of UHWs is always observed at altitude above 1000 km from the equator to the auroral zone.
Let us find the threshold value of UHW intensity for which the considered decay is possible under the magnetosphere plasma conditions. In a non-uniform plasma, the wave vectors $k_1$ and $k_2$ of interacting upper-hybrid waves undergo the same behaviour, and hence the mismatch $\Delta k$ is defined by the KAW wave-vector deviation: $\Delta k = \Delta (k_1 - k_2 - k_3) = -\Delta k_3$. For effective KAW generation, the time of parametric decay $\gamma_{NL}^{-1}$ must be less than the mismatch time $(\partial \ln k_3/\partial t)^{-1}$. Since $(\partial \ln k_3/\partial t) = (V_A/L_A)^{-1}$ and the shortest longitudinal scale on which the Alfvén velocity changes at an altitude of $10^3 - 10^4$ km is $L_A = 10^3$ km, we take (18) into account, put $V_A = 10^3$ km s$^{-1}$, and find the threshold value of UHW amplitude for decay instability: $W \approx 10^{-7}$. Usually $L_A > 10^3$ km, and hence the threshold is even lower.

In general, the threshold also depends on the product of the linear decrements of the decay products: $W \approx \gamma_L^2 \gamma_L^3$. But the phase velocity $V_{f2}$ of the decaying UHW is close to that of the pump wave, $\Delta V_{f2} \approx 10^{-3}$, and, in the region of phase velocities $V_{f2} \approx V_{f1}$, UHWs are almost undamped, $\gamma_L^2 \approx 0$ (and even can increase), because this is a region of UHW generation. Čerenkov damping of KAWs is also not essential ($\gamma_3^L \approx 0.01\omega_0$), and ion-cyclotron decay is negligible because of the anisotropic character of KAWs ($k_\perp \gg k_2$); even with $k_\perp \rho_i \gg 1$, the cyclotron reso-
Figure 3. Dependence of maximum growth rate on the normalized wavenumbers $k_{1z}d_e$ and $k_{1z}d_e$.

The resonance condition

$$\omega - k_z V_{1\text{res}} - n\omega_{B1} \approx -k_z V_{1\text{res}} - n\omega_{B2} = 0$$

holds for an exponentially small number of resonant particles ($V_{1\text{res}} = n\omega_{B1}/k_z \approx (\omega_{B1}/\omega)V_A$). That is why, in this case, $\gamma^L_2, \gamma^L_3 \approx 0$.

Nonlinear interaction investigations have been carried out on the basis of two-fluid MHD theory for the case of coherent waves. The use of two-fluid MHD is justified for the description of KAW dynamics when details of the particle distribution function in velocity space are not very important. The KAW dispersion obtained by means of the two-fluid MHD approximation is similar to that obtained with the help of kinetic theory with $k_{\perp} \rho_i < 1$ (changing 1 to $3/4$ in the coefficient of $k_{\perp} \rho_i$), and they coincide when $k_{\perp} \rho_i > 1$. Let us note that the restriction
$k_{⊥}\rho_{i} < 1$ does not exist for KAWs since even when $k_{⊥}\rho_{i} > 1$, the wave remains in the ultralow-frequency region ($\omega \ll \omega_{Bi}$) because of small $k_{z}$. The applicability of the coherent UHW interaction to magnetosphere conditions has been explained by Murtaza and Shukla (1984) using both theoretical and experimental data. Recent theoretical investigations have also shown that even warm electron beams should generate coherent upper-hybrid wave packets (Sigov and Levchenko 1996). Note also that, as observed by Oya et al. (1990), UHW intensity in magnetosphere is concentrated at the top of the spectrum. So, we can suppose that in the Earth’s magnetosphere coherent upper-hybrid wave packets are generated, and downspreading of the UHW spectrum can be caused by the considered process. For a UHW spectrum of bounded width, we can use (17) and (18), considering $W$ as part of the energy concentrated in the frequency interval $\Delta \omega \approx \gamma_{NL}$, i.e. the maximum growth rate is defined by the energy density in this part of the spectrum. The value used above $W = 10^{-5}$, which greatly exceeds the threshold, just represents the part of the energy of the spectrum with maximum intensity, lying near the upper margin of the UHW spectrum.

Similar processes of KAW generation by UHWs can also take place in the solar corona, where UHWs are supposed to be excited by electron beams.

8. Conclusions

Our investigations have shown that the parametric process $\text{UHW} \rightarrow \text{KAW} + \text{UHW}$ may be an effective mechanism of KAW generation in a plasma where upper-hybrid waves are excited by a plasma instability. Broadening of the UHW spectrum is expected in that case.

This process can take place in the equatorial and auroral regions of the Earth’s magnetosphere, where intense UHW emission is observed. At the same time, although Alfvén waves are often observed at the sites of enhanced UHW emission, the final solution of this problem is possible only after simultaneous satellite measurements of UHW and KAW amplitudes in narrow spectral intervals.

Acknowledgement

This work was supported in part by INTAS Grant 96-530.

References


Hughes, W. J. and Southwood, D. J. 1976 The screening of micropulsation signals by the atmosphere and ionosphere. J. Geophys. Res. 81, 3234.


Generation of kinetic Alfvén waves by upper-hybrid pump waves


