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Low-frequency wave instabilities in magnetoactive plasma with spatial inhomogeneity of temperature

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Abstract. The dispersion relation (DR) has been considered for low-frequency waves in a hot magnetoactive collisional plasma with 'longitudinal' large-scale electric field $\mathbf{E}_0(t)$ and weak spatial inhomogeneity of temperature. The electric field was assumed to be extremely weak ('subdriecer' field) and the quasi-static time variation of its amplitude allows us to use the 'direct initiation' mechanism to describe the instability development. Taking into account the ion thermal motion has not changed the form of the DR, which has remained the polynomial of the fourth order, but it has essentially complicated the form of its coefficients. As compared with previous investigations, these changes have resulted in the changes of the solutions of the DR. Two of the roots of the DR have remained of the same type ('kinetic Alfvén-like' waves), but another two have become quite different and get to the range of slow magneto-acoustic waves. Investigation of these new waves has shown that corresponding instability has clearly expressed threshold character, and the value of this threshold depends essentially on the reduced value of the electric field amplitude $E_0(t)$ and ratio of electron and ion temperatures.

1. Introduction

In 1975, Foukal first reported about the 'coronal loops with cool cores' in solar active regions (ARs) (Foukal 1975). Later, Tsap (1997) in his detailed analysis of magneto-hydrostatic equilibrium of the post-flare loops in ARs made the conclusion about a principal possibility of the existence of configurations with anomalous temperature gradients, i.e. with an increase of temperature in the cross section of the loop in the direction from the center to the periphery. In accordance with the modern view, the plasma instabilities are the main sources of the waves in the solar atmosphere in ARs (Somov 1994). Among all of these instabilities, the instabilities with extremely low thresholds of excitation (and, first of all, the drift ones) are the most interesting because of their specific role in plasma dynamics in preflare situations in ARs (Zaitsev et al. 1994; Somov 1994).

In the framework of the well-known Heyvaerts–Priest–Rust (HPR) theory (Heyvaerts et al. 1977) the situation seems obviously typical, when the rise of 'new' magnetic flux to the surface of the photosphere in ARs results in an increase in time of the amplitude of the large-scale quasi-static electric field $\mathbf{E}_0(t)$ in the current circuit of 'old' loop. In this case, the ratio of the amplitude $E_0(t)$ to the

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value of the local Dreicer field E_{D} is the real threshold value in the process of instability development. A weak electric field with

$$\varepsilon_{\mathrm{R}} \equiv \frac{E_0}{E_{\mathrm{D}}} \ll 1, \tag{1.1}$$

for which the boundary value $(\varepsilon_{\mathrm{R}})_{\mathrm{bound}}$ essentially does not exceed the 'equilibrium' value $(\varepsilon_{\mathrm{R}})_{\mathrm{EQ}}$, which corresponds to the 'steady state' in the circuit of the 'old' loop, is now called the 'subdreicer' field (Miller et al. 1997). There is observational evidence of correlation between the moments of the time, one of which corresponds to the maximum value of the $E_0(t)$ amplitude and another to the moment of the first energy release in a flare (Poletto and Kopp 1986; Zaitsev et al. 1994). So, if the time of instability development t_{inst} is much less than the time of increase of the electric field amplitude up to its maximum value, i.e.

$$t_{\mathrm{inst}} \ll t_{\mathrm{max}}, \tag{1.2}$$

then the lower values $\varepsilon_{\mathrm{R}} = (\varepsilon_{\mathrm{R}})_{\mathrm{bound}}$ correspond to the greater times for prediction of a flare. Taking into account that $t_{\mathrm{max}} \in [0.5; 30]$ minutes (Foukal and Hinata 1991; Polletto and Kopp 1986), it is clear that here we mean the case of short-time prediction of a flare. In contradistinction to our previous paper (Kryshtal 2002), in this situation it is not important whether the weak electric field effects or the drift effects dominate. We have to investigate them together: the extremely weak ('subdreicer') electric field values for ε_{R} from (1.1) and the influence of the weak spatial gradient of temperature, when the typical inequality

$$\frac{\rho_{\mathrm{i}}}{L} \ll 1 \tag{1.3}$$

is satisfied. Here $\rho_{\mathrm{i}} \equiv v_{\mathrm{Ti}}/\Omega_{\mathrm{i}}$ is the ion cyclotron radius, v_{Ti} is the ion thermal speed and Ω_{i} is their gyrofrequency in magnetic field $\mathbf{B}_0$. We have designated by L the mean scale of inhomogeneity of the plasma temperature, which we assume to be the same for the ions and electrons, i.e. $L_{\mathrm{e}} = L_{\mathrm{i}} = L$. In accordance with observational data, the most considerable inhomogeneities of the main plasma parameters are located in a loop in the direction perpendicular to its magnetic field $\mathbf{B}_0$ (Somov 1994; Zaitsev et al. 1994).

Taking into account this fact together with the results of Foukal (1975) and Tsap (1997), we have investigated the case of quasi-perpendicular propagation of the perturbation, when the condition

$$k_* \equiv \frac{k_z}{k_\perp} \ll 1 \tag{1.4}$$

is satisfied for the components of the wave vector $\mathbf{k}$, where $|\mathbf{k}|^2 = k_\perp^2 + k_z^2$ and $k_\perp^2 = k_x^2 + k_y^2$. Everywhere in the calculations we assumed $|k_x| = |k_y| = k_\perp/\sqrt{2}$.

Actually, we investigate the physical conditions of the rise and development of the instabilities of low-frequency waves in a hot magnetoactive plasma in a loop with extremely weak ('subdreicer') large-scale field $E_0(t)$ and weak spatial inhomogeneity of temperature in ARs. These waves are assumed to be the solutions of the corresponding dispersion relation (DR).

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2. Plasma model and basic equations

As in our previous papers (Kryshtal and Kucherenko 1995; Kryshtal and Gerasimenko 2002), the plasma under consideration is assumed to be fully ionized, and collisions are described by the model Bhatnagar-Gross-Krook (BGK)-integral (Alexandrov et al. 1988). In this plasma the equilibrium velocity distribution function is assumed to be pure Maxwellian for ions, but the same function for the electrons is described by the shifted Maxwellian distribution with electron shift velocity

$$u_{\mathrm{e}} = \frac{eE_0}{m_{\mathrm{e}}\nu_{\mathrm{ei}}}, \tag{2.1}$$

where ν_{ei} is the ion–electron collision frequency. In the case of 'weak inhomogeneity', when there is a one-dimensional (along the x-axis) spatial gradient of temperature (which increases with the increase of x) in plasma and ion–electron collisions dominate, the small parameter ε_{R} can be described in following form:

$$\varepsilon_{\mathrm{R}}\left(\equiv \frac{E_0}{E_{\mathrm{D}}}\right) = \frac{u_{\mathrm{e}}}{\upsilon_{\mathrm{Te}}} \quad (\ll 1), \tag{2.2}$$

where υ_{Te} is the electron thermal speed. For the long-wave perturbations with

$$z_{\mathrm{e}} = \frac{k_\perp^2 \upsilon_{\mathrm{Te}}^2}{\Omega_{\mathrm{e}}^2} \ll z_{\mathrm{i}} = \frac{k_\perp^2 \upsilon_{\mathrm{Ti}}^2}{\Omega_{\mathrm{i}}^2} \ll 1, \tag{2.3}$$

where Ω_{e} is the electron gyrofrequency, we consider only the local (Michailovsky 1963, 1975) solutions of the DR, which allow us to neglect the influence of the boundaries. The z-axis of the local Cartesian coordinate system is directed along the magnetic field $\mathbf{B}_0$ and the perturbations are propagated in a direction almost perpendicular to the direction of $\mathbf{B}_0$. This means that the xy-plane is actually parallel to the surface of the photosphere. One can easily show that condition (1.3) is the result of the 'geometrical optics' condition

$$\frac{\lambda_\perp}{L} \ll 1, \tag{2.4}$$

where $\lambda_\perp$ is the 'perpendicular' wavelength of perturbation: $\lambda_\perp^2 \equiv \lambda_x^2 + \lambda_y^2$. We take the phase velocity of the perturbation to vary in the range

$$\upsilon_{\mathrm{Ti}} \ll \frac{\omega}{k_z} \ll \upsilon_{\mathrm{Te}}, \tag{2.5}$$

and 'plasma $\beta_{\mathrm{e,i}}$' to vary in the range (Michailovsky 1975)

$$\frac{m_{\mathrm{e}}}{m_{\mathrm{i}}} \ll \beta_{\mathrm{e,i}} \equiv \frac{8\pi n_{0\mathrm{e,i}} k_B T_{0\mathrm{e,i}}}{B_0^2} \ll 1, \tag{2.6}$$

where $T_{0\mathrm{e,i}}(x)$ are the equilibrium ('unperturbed') temperatures of electrons and ions, respectively. In this paper we do not require the ratio

$$t \equiv \frac{T_{0\mathrm{e}}}{T_{0\mathrm{i}}} = 1 \tag{2.7}$$

to be satisfied, as was necessary for the analogous condition for the densities ('quasi-neutrality' condition). However, we have used the analogous equation for the spatial gradients

$$\frac{\partial T_{0\mathrm{e}}(x)}{\partial x} = \frac{\partial T_{0\mathrm{i}}}{\partial x} \tag{2.8}$$

as a physical condition and mathematical simplification.

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We call the spatial inhomogeneity of the plasma temperature 'weak' if the following equation is satisfied (Michailovsky 1963):

$$\omega^{*}_{\mathrm{e,i}} \ll \omega, \tag{2.9}$$

where

$$\omega^{*}_{\mathrm{e,i}} \equiv \frac{k_y v^2_{\mathrm{Te,i}}}{\Omega_{\mathrm{e,i}}} \frac{\partial}{\partial x}[\ln T_{0\mathrm{e,i}}(x)] \tag{2.10}$$

are the 'temperature' drift frequencies of the electrons and ions. The conditions similar to (2.9) for the reduced values

$$\Omega^{*}_{\mathrm{e,i}} \equiv \frac{\omega^{*}_{\mathrm{e,i}}}{k_z v_{\mathrm{A}}} \ll 1 \tag{2.11}$$

supplement the requirement of the weakness of the electric field (2.2). Obviously, the conditions (1.1) and (2.2) coincide only in the situation when the turbulence is absent (Tsap 1997) and

$$\nu_{\mathrm{eff}} = \nu_{\mathrm{ei}}, \tag{2.12}$$

where ν_{eff} is the effective collision frequency (Alexandrov et al. 1988). In addition, we have to point to the one important equation. We have assumed that due to the specific conditions of the plasma atmosphere of a flare in the arch structures in ARs we can neglect the gradients of charged-particle densities and the gradient of the magnetic field intensity B_0 in comparison with the temperature gradient. However, the shear of the magnetic field force lines is one of the most important characteristics not only for the laboratory plasmas (Kadomtsev and Pogootse 1967), but also for the plasmas in the loop structures (Heyvaerts et al. 1977; Zaitsev et al. 1994; Somov, 1994). This means that if we assume that the reduced temperature spatial gradient dominates, then we have to require the equation

$$\frac{1}{2}\beta_{\mathrm{i}} \equiv \frac{4\pi n_{0\mathrm{i}} k_B T_{0\mathrm{i}}}{B_0^2} \ll \frac{z_{\mathrm{i}} k_*^2}{\beta_{\mathrm{A}}^2 \varepsilon_{\mathrm{R}}^2} \tag{2.13}$$

to be satisfied (Kadomtsev and Pogootse 1967). It can be easily shown that the analogous condition for the 'electron plasma β' has the form

$$\tfrac{1}{2}\beta_{\mathrm{e}}\varepsilon_{\mathrm{R}} \ll \mu\sqrt{z_{\mathrm{i}}t}k_*. \tag{2.14}$$

In (2.13) and (2.14) we have used the following designations:

$$\beta_{\mathrm{A}} \equiv \frac{v_{\mathrm{Te}}}{v_{\mathrm{A}}} \tag{2.15}$$

and

$$\mu \equiv \sqrt{\frac{m_{\mathrm{e}}}{m_{\mathrm{i}}}} \cong 2.33 \times 10^{-2}.$$

Obviously, the inequalities (2.13) and (2.14) can only be satisfied for the extremely small values of ε_{R} and a 'happy' choice of plasma equation of state. The values

$$n_{0\mathrm{i}} = n_{0\mathrm{e}} = 6.89 \times 10^{10}\ \mathrm{cm}^{-3}; \quad T_{0\mathrm{i}} (= T_{0\mathrm{e}}) = 4.47 \times 10^{5}\ \mathrm{K} \tag{2.16}$$

from Machado et al. (1980), which correspond to the height $h = 1459$ km above the photosphere and a value of magnetic field amplitude $B_0 = 31.3$ G, provide us with namely the same 'happy' equation of state, which we need.

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Taking into account all of the remarks mentioned above, we can say that the plasma model under consideration is practically the same as the former model with only one essential difference: the spatial gradient of temperature replaces the spatial gradient of density. This leads to the changes in the expressions for the real and imaginary parts of the 'longitudinal' $\varepsilon_\parallel$ and 'transverse' $\varepsilon_\perp$ scalars of dielectric permeability. At the same time the form of DR for low-frequency waves (Michailovsky 1975)

$$\varepsilon_\parallel\left(1-\frac{\omega^2}{c^2k_z^2}\varepsilon_\perp\right)+\frac{k_\perp^2}{k_z^2}\varepsilon_\perp=0 \tag{2.17}$$

remains without any changes. As formerly, in present paper the functions $\varepsilon_\parallel$ and $\varepsilon_\perp$ have the following form:

$$\varepsilon_\parallel=1+\left(\frac{k_\perp}{k_z}\right)^2\frac{\omega_{*\mathrm{i}}^2}{z_\mathrm{i}}(\mathrm{Re}\,\varepsilon_\parallel^0+i\,\mathrm{Im}\,\varepsilon_\parallel^0), \tag{2.18}$$

$$\varepsilon_\perp=1+\frac{\omega_{*\mathrm{i}}^2}{z_\mathrm{i}}(\mathrm{Re}\,\varepsilon_\perp^0+i\,\mathrm{Im}\,\varepsilon_\perp^0), \tag{2.19}$$

where

$$\omega_{*\alpha}\equiv\frac{\omega_{\mathrm{P}\alpha}}{\Omega_\alpha},\quad(\alpha=\mathrm{e},\mathrm{i}) \tag{2.20}$$

and $\omega_{\mathrm{P}\alpha}$ is the plasma frequency for charged particles of type α. However, the expressions for $\mathrm{Re}\,\varepsilon_\parallel^0$, $\mathrm{Re}\,\varepsilon_\perp^0$, $\mathrm{Im}\,\varepsilon_\parallel^0$ and $\mathrm{Im}\,\varepsilon_\perp^0$ become more complicated:

$$\begin{aligned}\mathrm{Re}\,\varepsilon_\parallel^0=\frac{1}{t\Omega(\Omega-\beta_\mathrm{A}\varepsilon_\mathrm{R})^2}\Bigg\{&\Omega^3(1+\tau_*)+\Omega^2\left[-2\beta_\mathrm{A}\varepsilon_\mathrm{R}(1+\tau_*)-\frac{7}{4}\Omega^*\tau_*\right]\\&+\Omega\left[\beta_\mathrm{A}^2\varepsilon_\mathrm{R}^2(1+\tau_*)-\beta_\mathrm{A}^2\mu^2\left(1-\frac{\tau_*^2}{2\pi\mu^2}\right)+\frac{7}{4}\beta_\mathrm{A}\varepsilon_\mathrm{R}\Omega^*\tau\right]\\&-\beta_\mathrm{A}^2\mu^2\left[2\beta_\mathrm{A}\varepsilon_\mathrm{R}+\frac{\Omega^*}{t}(1-2z_\mathrm{i})\right]\Bigg\};\end{aligned}$$

$$\begin{aligned}\mathrm{Im}\,\varepsilon_\parallel^0=\frac{\sqrt{\pi/2}}{t\beta_\mathrm{A}\Omega^2(\Omega-\beta_\mathrm{A}\varepsilon_\mathrm{R})}\Bigg\{&\Omega^4+\left(\frac{\Omega^*}{2}-2\beta_\mathrm{A}\varepsilon_\mathrm{R}\right)\Omega^3\\&+\left[\beta_\mathrm{A}\varepsilon_\mathrm{R}\left(\beta_\mathrm{A}\varepsilon_\mathrm{R}+\frac{\Omega^*}{2}\right)+\frac{\tau_*\beta_\mathrm{A}^2z_\mathrm{i}t}{\pi\sqrt{2}}\right]\Omega^2+\frac{\tau_*\beta_\mathrm{A}^2}{\pi\sqrt{2}}(2\beta_\mathrm{A}^2\mu^2-\beta_\mathrm{A}\varepsilon_\mathrm{R}z_\mathrm{i}t)\Omega\\&+\frac{2\tau_*\beta_\mathrm{A}^2}{\pi\sqrt{2}}\beta_\mathrm{A}^2\mu^2\frac{\Omega^*}{t}\left[1-2z_\mathrm{i}+\frac{3(1-z_\mathrm{i})}{4\sqrt{t}}\right]\Bigg\};\end{aligned}$$

$$\mathrm{Re}\,\varepsilon_\perp^0=z_\mathrm{i}\left(1+\frac{\Omega^*}{t\Omega}\right) \tag{2.21}$$

$$\mathrm{Im}\,\varepsilon_\perp^0=z_\mathrm{i}\frac{\nu}{\Omega}, \tag{2.22}$$

where

$$\Omega\equiv\frac{\omega}{k_zv_\mathrm{A}}, \tag{2.23}$$

$$\nu\equiv\frac{\nu_\mathrm{ei}}{\Omega_\mathrm{i}}\sigma_\mathrm{eff};\quad\sigma_\mathrm{eff}\in[1;1.5] \tag{2.24}$$

$$\tau_*=\sqrt{\frac{2\pi}{z_\mathrm{i}t}}\frac{\nu\mu}{k_*}. \tag{2.25}$$

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In the expressions for $\mathrm{Re}\,\varepsilon_\perp^0$ and $\mathrm{Im}\,\varepsilon_\perp^0$ we have used the inequality (2.3), i.e. $z_\mathrm{e} \ll z_\mathrm{i}$, which is obviously satisfied for the 'equation of state' (2.16) and the mentioned value of B_0.

By means of the factor σ_eff in (2.24) we approximately take into account the contributions (to the effective collision frequency (2.12)) of the mutual collisions of the particles with the same sign of charge (Kryshtal and Kucherenko 1995; Kryshtal 1998).

In (2.18)–(2.22) we have used the designation:

$$\Omega^* = \Omega_\mathrm{e}^* = -t\Omega_\mathrm{i}^*. \tag{2.26}$$

For the one-dimensional temperature function of sufficiently general form

$$T_0(x) = \text{constant} \times \exp\left[+\frac{x}{L}\right] \tag{2.27}$$

the expression for Ω^* can be written as

$$\Omega^* = \sqrt{\frac{t}{2}}\beta_\mathrm{A}\frac{\mu\rho}{k_*}, \tag{2.28}$$

where

$$\rho \equiv \frac{\rho_\mathrm{i}}{L}. \tag{2.29}$$

Obviously, in this case Ω^* is negative.

In fact, the procedure of the transformation of DR (2.17) to the modified dispersion relation (MDR) of the polynomial form remains the same. However, for the extremely small values of ε_R and taking into account in (2.23) and (2.24) the terms which describe the ion thermal motion (the terms with factor $k_z^2 v_\mathrm{Ti}^2/\omega^2$, formally negligible!), the situation essentially changes. It is impossible in this case to obtain the analytical solutions of the MDR and we have to use the Euler method to solve it in a numerical way (Kryshtal and Gerasimenko, 2002). For the extremely small values of ε_R and the moderately non-isothermal plasma with

$$1 < t < 20, \tag{2.30}$$

and with the 'equation of state' (2.16), the MDR for low-frequency waves has the following form:

$$\sum_{i=0}^{4} P_\mathrm{i}\Omega^i = 0 \tag{2.31}$$

with

$$\begin{aligned} P_0 = -\beta_\mathrm{A}^2\mu^2 \Bigg\{ & -1 + \frac{\tau_*^2}{2\pi\mu^2} + \beta_\mathrm{A}^2\varepsilon_\mathrm{R}\frac{\nu\tau_*}{2\sqrt{\pi}} \\ & + \frac{\Omega^*}{t}\left[2\beta_\mathrm{A}\varepsilon_\mathrm{R} + \frac{\Omega^*}{t}(1-2z_\mathrm{i}) + \frac{\beta_\mathrm{A}\nu\tau_*}{\sqrt{\pi}}\left(1 - -2z_\mathrm{i} + \frac{3(1-z_\mathrm{i})}{4\sqrt{t}}\right)\right] \\ & + \beta_\mathrm{A}\varepsilon_\mathrm{R}\left[\Omega^*\left(\frac{7}{4}\tau_* - 2z_\mathrm{i}\right) + \beta_\mathrm{A}\varepsilon_\mathrm{R}\left[1 + \tau_* + z_\mathrm{i}t\left(1 + \frac{\beta_\mathrm{A}\nu\tau_*}{2\sqrt{\pi}}\right)\right]\right]\Bigg\}; \end{aligned} \tag{2.32}$$

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$$P_1 \equiv 2\beta_A\varepsilon_R(1+z_i t+\tau_*)+\sqrt{\frac{\pi}{2}}\beta_A\nu\left[\frac{\sqrt{2}}{\pi}\tau_*\beta_A(\varepsilon_R z_i t-\beta_A\mu^2)+\varepsilon_R^2\frac{\Omega^*}{2}\right]$$
$$-\Omega^*\left(z_i-\frac{7}{4}\tau_*\right)+\beta_A\varepsilon_R\frac{\Omega^*}{t}\left[\frac{7}{4}\Omega^*\tau_*+\beta_A\varepsilon_R(1+\tau_*)\right]$$
$$-2\beta_A^2\mu^2\left[\beta_A\varepsilon_R+\frac{\Omega^*}{t}\left(1-z_i-\frac{\tau_*^2}{4\pi\mu^2}\right)\right]; \tag{2.33}$$

$$P_3 \equiv (1+\tau_*)\left(\frac{\Omega^*}{t}-2\beta_A\varepsilon_R\right); \tag{2.34}$$

$$P_4 \equiv 1+\tau_*-\sqrt{\frac{\pi}{2}}\frac{\nu}{\beta_A}; \tag{2.35}$$

$$P_2 \equiv -1-z_i t-\tau_*-\frac{\Omega^*}{t}\beta_A\varepsilon_R(1+\tau_*)+\left(\beta_A\varepsilon_R-\frac{\Omega^*}{t}\right)\left[\beta_A\varepsilon_R(1+\tau_*)+\frac{7}{4}\Omega^*\tau_*\right]$$
$$-\beta_A^2\mu^2\left(1-\frac{\tau_*^2}{4\pi\mu^2}\right)-\sqrt{\frac{\pi}{2}}\beta_A\nu\left(3\varepsilon_R^2+\frac{\tau_* z_i t}{\pi\sqrt{2}}\right). \tag{2.36}$$

Now a very important question arises: what are the values of ε_R we have to consider as 'extremely small'?

3. Solutions of the modified dispersion relation

Earlier we had pointed to the fact that the distinguishing characteristic of the current circuit of post-flare loop is the existence of a 'weak' current (and corresponding weak 'dc' electric field) in a stationary, 'steady' state of the system, when interaction between the 'old' and 'new' magnetic fluxes is absent (Heyvaerts et al. 1977; Foukal and Hinata 1991; Zaitsev et al. 1994). When the electron stream velocity is described by expression (2.1) and the electron–ion collisions dominate, one can easily obtain the 'equilibrium' value of ε_R as follows:

$$(\varepsilon_R)_{EQ}=\frac{I}{\pi R_0^2 n_0 e v_{Te}}. \tag{3.1}$$

Here I is the current in the loop and R_0 is the radius of its small cross section. For $I=10^{12}$ A and $R_0=10^8$ cm (Zaitsev et al. 1994; Miller et al. 1997) we obtain $(\varepsilon_R)_{EQ}=1.1\times10^{-5}$. Next, we call the values of ε_R 'extremely small' if they vary in the range

$$10^{-5}\leqslant\varepsilon_R\leqslant10^{-4}. \tag{3.2}$$

In accordance with HPR-theory (Heyvaerts et al. 1977) these values correspond to the earliest stage of interaction of the fluxes. In the framework of this theory it seems natural to require the values of parameter t to be minimum as well as for the values of parameters z_i and k_* to be maximum. In the following calculations we shall assume that they vary in the ranges

$$1\leqslant t\leqslant 6; \tag{3.3}$$
$$0.10\leqslant z_i\leqslant 0.15; \tag{3.4}$$
$$0.09\leqslant k_*\leqslant 0.12. \tag{3.5}$$

We call this case the 'case of weak inhomogeneity'. Taking into account the one-dimensional spatial inhomogeneity of temperature, we can say that the most

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distinguishing characteristics of the case under consideration are the maximum values of k_* and the fact that electron–ion collisions dominate. Relatively small values of t and relatively 'large' values of z_i can be considered as additional distinguishing characteristics.

Next we turn our attention on the only real roots of MDR (2.31). This means that we do not consider the processes which are not periodic (so-called 'aperiodic' instabilities or 'aperiodic' damping (Michailovsky 1975)). It is a well-known fact (Van der Waerden 1930) that all of the roots of the polynomial of the fourth order are real if

$$D < 0, \tag{3.6}$$

where D is the discriminant of the resolvent equation of the third order, i.e. the known function of the expressions (2.32) ÷ (2.36). The calculations have shown that condition (3.6) imposes sufficiently rigid restrictions on the possible values of z_i and k_* as well as on the values of $(\varepsilon_\mathrm{R})_\mathrm{bound}$ and t_bound at which the corresponding waves are generated. This results in the fact that 'permissible bands' of z_i and k_* at the constant values of $(\varepsilon_\mathrm{R})_\mathrm{bound}$ and t_bound turn out to be narrow. Of course, the only wave will be generated that has positive growth rate, so the values of $(\varepsilon_\mathrm{R})_\mathrm{bound}$ and t_bound we determine not only from the condition (3.6), but from the condition for the growth rate too:

$$\gamma \geqslant 0. \tag{3.7}$$

The 'narrow bands' for z_i and k_* lead to the result that practically for all four roots of MDR (2.31) the surfaces of the reduced longitudinal phase velocities Ω_i ($i = 1, 2, 3, 4$) as functions of z_i and k_* (i.e. $\Omega_\mathrm{i} = \Omega_\mathrm{i}(z_\mathrm{i}, k_*)$) are almost flat. In a sense this fact has no great significance. The information seems far more important for us about the dependence of the functions Ω_i on the parameters z_i and k_* in addition to the ranges of the change of the values $\|\Omega_\mathrm{i}\|$. We assume that this information can help us to identify (more or less exactly) the type of plasma wave, which corresponds to the certain root Ω_i of MDR (2.31).

4. Analysis of stability: expression for the growth rate and numerical simulation

In the framework of the approximations employed we can use the well-known formula for the reduced growth rate (Michailovsky 1975; Kadomtsev and Pogootse 1967)

$$\Gamma(\Omega_m) \equiv \left(\frac{\gamma}{\omega}\right)_{\Omega=\Omega_m} = -\left(\frac{1}{\omega}\frac{\operatorname{Im} D(\omega, \mathbf{k}, x)}{(\partial/\partial\omega)\operatorname{Re} D(\omega, \mathbf{k}, x)}\right)_{\Omega=\Omega_m} \qquad (m = 1\text{–}4). \tag{4.1}$$

Here we take $\operatorname{Im} D(\omega, \mathbf{k}, x)$ and $\operatorname{Re} D(\omega, \mathbf{k}, x)$ to be the imaginary and real parts of the MDR (2.31). The final form of the reduced growth rate is the following:

$$\Gamma(\Omega_m) = \sqrt{\frac{\pi}{2}}\frac{(\Omega - \beta_\mathrm{A}\varepsilon_\mathrm{R})}{\beta_\mathrm{A}\cdot\Omega^2}\frac{F_1}{F_2}\bigg|_{\Omega=\Omega_m} \tag{4.2}$$

with

$$F_1 = \sum_{i=0}^{6} C_\mathrm{i}\Omega^i, \tag{4.3}$$

$$F_2 = \sum_{i=0}^{4} S_\mathrm{i}\Omega^i, \tag{4.4}$$

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where

$$C_0 \equiv \frac{\tau_* \beta_{\rm A}^4}{\pi\sqrt{2}}\left\{-2\mu^2\left[\frac{\Omega^*}{t}\left(1-2z_{\rm i}+\frac{3(1-z_{\rm i})}{4\sqrt{t}}\right)-\beta_{\rm A}\varepsilon_{\rm R}\right]-\varepsilon_{\rm R}^2 z_{\rm i} t\right\}-\sqrt{\frac{2}{\pi}}\beta_{\rm A}^3\varepsilon_{\rm R}^2\nu z_{\rm i} t; \tag{4.5}$$

$$C_1 \equiv \beta_{\rm A}^2\left\{\frac{\sqrt{2}\tau_*\beta_{\rm A}}{\pi}(\varepsilon_{\rm R} z_{\rm i} t-\beta_{\rm A}\mu^2)+\varepsilon_{\rm R}^2\frac{\Omega^*}{2}+\sqrt{\frac{8}{\pi}}\nu\left[\varepsilon_{\rm R} z_{\rm i} t-\beta_{\rm A}\mu^2\frac{\Omega^*}{2t}(1-2z)\right]\right\} \tag{4.6}$$

$$\begin{aligned} C_2 \equiv{}& \sqrt{\frac{2}{\pi}}\beta_{\rm A}\nu\left[-z_{\rm i} t+\beta_{\rm A}^2\varepsilon_{\rm R}^2(1+\tau_*)-\beta_{\rm A}^2\mu^2\left(1-\frac{\tau_*^2}{2\pi\mu^2}\right)+\frac{7}{4}\beta_{\rm A}\varepsilon_{\rm R}\Omega^*\tau_*\right] \\ &-\beta_{\rm A}^2\varepsilon_{\rm R}^2\left(3+\frac{\Omega^{*2}}{t^2}\right)+\frac{\tau_*\beta_{\rm A}^2}{\pi\sqrt{2}}\left\{4\beta_{\rm A}^2\mu^2\frac{\Omega^*}{t}\left[1-z_{\rm i}+\frac{3(1-z_{\rm i})}{4\sqrt{t}}\right]\right. \\ &\left.-z_{\rm i} t-\beta_{\rm A}\varepsilon_{\rm R}(2z_{\rm i}\Omega^*+2\beta_{\rm A}^2\mu^2-\beta_{\rm A}\varepsilon_{\rm R} z_{\rm i} t)\right\}; \end{aligned} \tag{4.7}$$

$$\begin{aligned} C_3 \equiv{}& 3\beta_{\rm A}\varepsilon_{\rm R}-\frac{\Omega^*}{2}-\sqrt{\frac{2}{\pi}}\beta_{\rm A}\nu\left[2\beta_{\rm A}\varepsilon_{\rm R}(1+\tau_*)+\frac{7}{4}\Omega^*\tau_*\right]+\beta_{\rm A}^2\varepsilon_{\rm R}^2\Omega^*\left(\frac{3}{t}-\frac{1}{2}\right) \\ &+\frac{\tau_*\beta_{\rm A}^2}{\pi\sqrt{2}}(2\beta_{\rm A}^2\mu^2+z_{\rm i}\Omega^*-2\beta_{\rm A}\varepsilon_{\rm R} z_{\rm i} t); \end{aligned} \tag{4.8}$$

$$C_4 \equiv -1-\beta_{\rm A}^2\varepsilon_{\rm R}^2+\sqrt{\frac{2}{\pi}}\beta_{\rm A}\nu(1+\tau_*)+\frac{\Omega^{*2}}{2t}+3\beta_{\rm A}\varepsilon_{\rm R}\frac{\Omega^*}{t}+\frac{\tau_*\beta_{\rm A}^2}{\pi\sqrt{2}}z_{\rm i} t; \tag{4.9}$$

$$C_5 \equiv -3\beta_{\rm A}\varepsilon_{\rm R}+\Omega^*\left(\frac{1}{t}+\frac{1}{2}\right); \quad C_6 \equiv 1; \tag{4.10}$$

$$\begin{gathered} S_0 \equiv 2P_0+\beta_{\rm A}\varepsilon_{\rm R}P_1; \quad S_1 \equiv P_1+2P_2\beta_{\rm A}\varepsilon_{\rm R}; \quad S_2 \equiv 3\beta_{\rm A}\varepsilon_{\rm R}P_3; \\ S_3 \equiv 4P_4\beta_{\rm A}\varepsilon_{\rm R}-P_3; \quad S_4 \equiv -2P_4. \end{gathered} \tag{4.11}$$

All four roots of MDR (2.31) have been obtained in a numerical way based on the Euler method (Van der Waerden 1930; Kryshtal and Gerasimenko 2002). The corresponding values of the reduced growth rates have been obtained by making use of the formulas (4.2)–(4.11). In the process of calculation, all of the above-mentioned conditions have been supplemented by the condition

$$\Gamma_m \ll 1 \quad (m = 1\text{–}4), \tag{4.12}$$

which actually means the 'legitimacy' of linear approximation of the perturbation theory. Practically, it makes the bands for $z_{\rm i}$ and k_* at constant values of $(\varepsilon_{\rm R})_{\rm bound}$ and $t_{\rm bound}$ even more narrow. We have performed the numerical simulation of the behavior of the functions $\Omega_m = \Omega_m(z_{\rm i}, k_*)$ for the all real roots $\Omega_m (m = 1\text{–}4)$ of MDR (2.31) as well as the functions $\Gamma(\Omega_m)$. It has been performed for the plasma equation of state (2.16) (Machado et al. 1980) with $B_0 = 31.3$ G and $\sigma_{\rm eff} = 1.5$ (see (2.24)). We have obtained for the parameters $\varepsilon_{\rm R}$ and t the following boundary values:

$$t_{\rm bound} = 4.3; \tag{4.13}$$

$$(\varepsilon_{\rm R})_{\rm bound} = 6.5 \times 10^{-5}. \tag{4.14}$$

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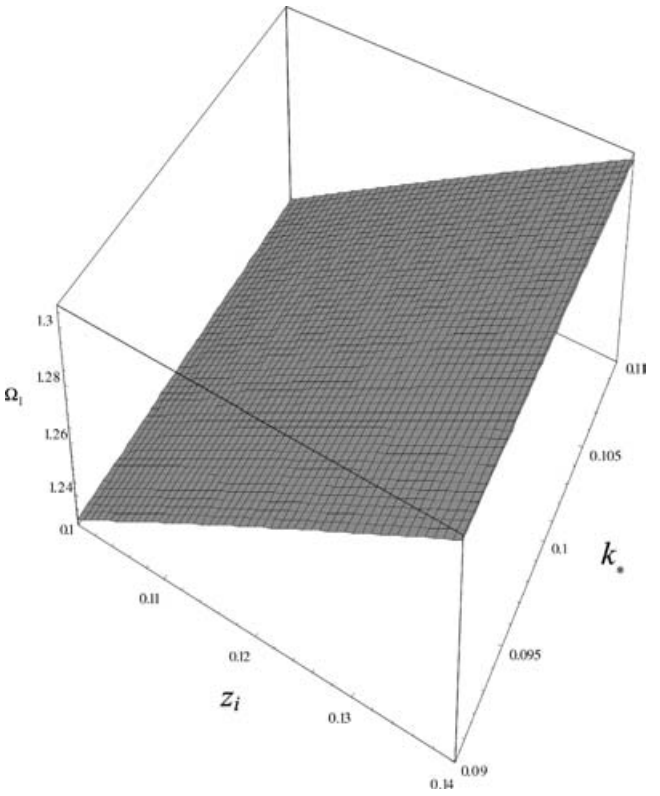


Figure 1. The reduced 'longitudinal' phase velocities as the functions of the perturbation wave vector tilt angle k_* and the ion kinetic parameter z_i at $t = t_\mathrm{bound} = 4.3$ and $\varepsilon_\mathrm{R} = (\varepsilon_\mathrm{R})_\mathrm{bound} = 6.5 \times 10^{-5}$ for the 'direct kinetic Alfvén-like' waves.

Only for $t \geqslant t_\mathrm{bound}$ and $\varepsilon_\mathrm{R} > (\varepsilon_\mathrm{R})_\mathrm{bound}$ the instability under consideration first appears at some specific values of z_i and k_*, i.e. the function $\Gamma_k = \Gamma_k(z_\mathrm{i}, k_*)$ $(k = 1, 2, 3, 4)$ first emerges to the 'positive' half-space $\Gamma_k > 0$, namely at these values of z_i and k_* with $t \geqslant t_\mathrm{bound}, (\varepsilon_\mathrm{R}) \geqslant (\varepsilon_\mathrm{R})_\mathrm{bound}$. Of course, we imply that conditions (3.6) and (4.12) are satisfied in this situation. Figures 1–4 show the behavior of the roots of MDR (2.31) when parameters z_i and k_* vary in the intervals which correspond to the case of 'weak inhomogeneity'. The relatively low value of t_bound shows that ion-acoustic instability could hardly be possible in this case. Here it seems natural to note that the difference between the electron temperature and that of the ions is essential for the dynamics of high-temperature plasma of solar ARs. It is well known (Somov 1994) that real heat fluxes in the framework of two-temperature models with $t > 1$ can be an order of magnitude ($2^{7/2} \approx 10$) larger than the fluxes calculated in one-temperature ($T_\mathrm{e} = T_\mathrm{i}$) models. Because of this, the one-temperature models are much less dynamic ('more sluggish') than one would expect. 'Classical' flare models (see the references in Somov (1994)) have, as a rule, $1 < t \leqslant 2, 5$. In this sense the model of preflare plasma under investigation is much closer to the well-known model of Duijveman et al. (1981), where the authors take t to vary in the range $3, 1 \leqslant t \leqslant 6, 4$. At the same time the value of $(\varepsilon_\mathrm{R})_\mathrm{bound}$ shows the fact that we are really studying the early stage of the flare process, because in a preflare situation the initial electric field amplitude can increase by more than

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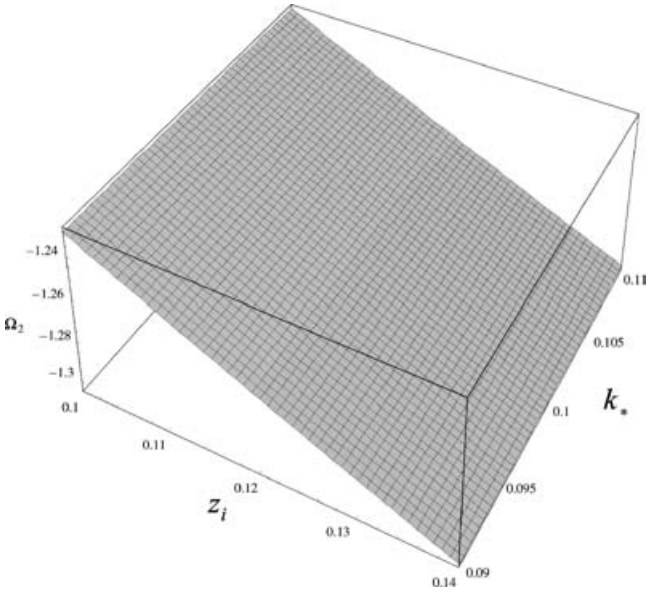


Figure 2. The reduced 'longitudinal' phase velocities as the functions of the perturbation wave vector tilt angle k_* and the ion kinetic parameter z_i at $t = t_\mathrm{bound} = 4.3$ and $\varepsilon_\mathrm{R} = (\varepsilon_\mathrm{R})_\mathrm{bound} = 6.5 \times 10^{-5}$ for the 'inverse kinetic Alfvén-like' waves.

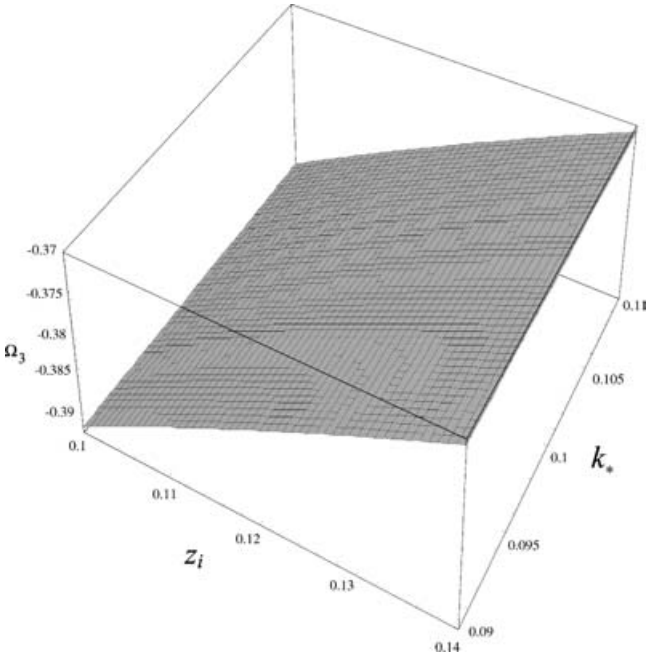


Figure 3. The reduced 'longitudinal' phase velocities as the functions of the perturbation wave vector tilt angle k_* and the ion kinetic parameter z_i at $t = t_\mathrm{bound} = 4.3$ and $\varepsilon_\mathrm{R} = (\varepsilon_\mathrm{R})_\mathrm{bound} = 6.5 \times 10^{-5}$ for the 'inverse' SMAWs with $\Omega = \Omega_3$.

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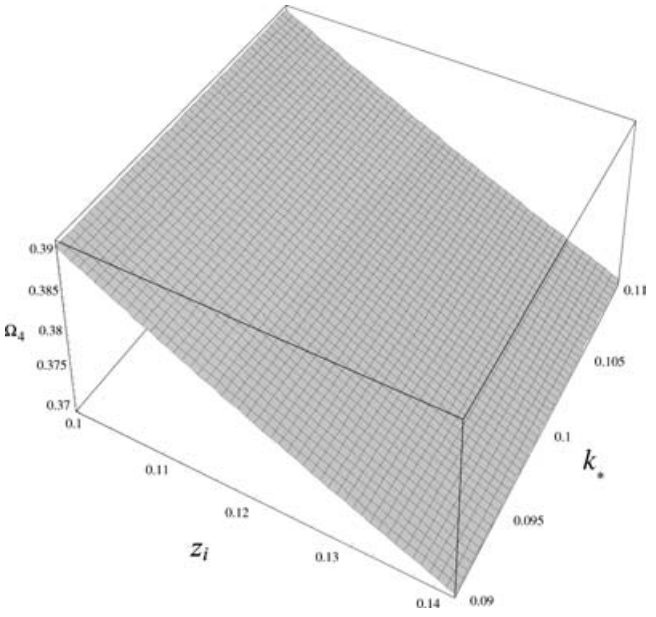


Figure 4. The reduced 'longitudinal' phase velocities as the functions of the perturbation wave vector tilt angle k_* and the ion kinetic parameter z_i at $t = t_\mathrm{bound} = 4.3$ and $\varepsilon_\mathrm{R} = (\varepsilon_\mathrm{R})_\mathrm{bound} = 6.5 \times 10^{-5}$ for the 'direct' SMAWs with $\Omega = \Omega_4$.

10^2–10^3 times (Poletto and Kopp 1986; Zaitsev et al. 1994). It is not so hard to identify the waves that correspond to the roots Ω_1 and Ω_2. Next we call them, in a brief form, 'Ω_1-wave' and 'Ω_2-wave'. For these waves, the values $\|\Omega_1\|$ and $\|\Omega_2\|$ do not depend on k_* and increase when z_i increases. Taking into account the numerical values of Ω_1 and Ω_2, we can say that they are kinetic Alfvén waves (KAWs), that of 'direct' (Ω_1 with $k_z > 0$) and 'inverse' (Ω_2 with $k_z < 0$). The dispersion laws of these waves have small corrections due to the existence in a plasma of the weak large-scale electric field $\mathbf{E}_0(t)$, pair Coulomb collisions and weak spatial inhomogeneity of the temperature. However, for 'Ω_3-wave' and 'Ω_4-wave' the problem of identification is not so simple. They are definitely not the fast magneto-acoustic waves (FMAWs), because the dependence of the values Ω_3 and Ω_4 on k_* is absent. At the same time the values $\|\Omega_3\|$ and $\|\Omega_4\|$ are relatively small for the FMAWs. Essentially because of the same reasons we cannot identify these waves as KAWs. So, we can only propose to consider them as 'direct' (Ω_4 with $k_z > 0$) and 'inverse' (Ω_3 with $k_z < 0$) waves from the range of the slow magneto-acoustic ones (SMAWs), which are 'modified' by the same corrections as the mentioned above KAWs. At the same time we cannot consider these waves exactly as 'modified SMAWs', because these corrections are not so small in this case and (what is more important) because of the fact that the values $\|\Omega_3\|$ and $\|\Omega_4\|$ considerably decrease, when z_i increases. By the way, this dependence also does not allow us to identify the 'Ω_3-wave' and 'Ω_4-wave' as 'modified KAWs'.

Figures 5–8 show the behavior of the reduced growth rates $\Gamma_k \equiv \Gamma(\Omega_k) = \Gamma_k(z_\mathrm{i}, k_*)$ $(k = 1, 2, 3, 4)$ at the same values of t_bound and $(\varepsilon_\mathrm{R})_\mathrm{bound}$ as in (4.13) and (4.14). The values of z_i and k_* vary in the area on (z_i, k_*)-plane, which corresponds to the case of 'weak inhomogeneity' (Kryshtal and Gerasimenko 2002). As we can

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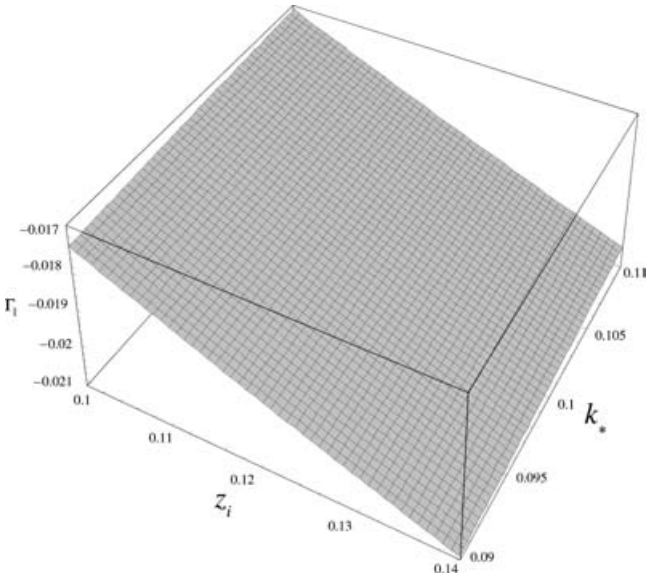


Figure 5. The reduced growth rates $\Gamma_k = \Gamma(\Omega_k)$ $(k = 1, 2, 3, 4)$ of the four types of the waves with $\Omega = \Omega_k$ as the functions of the perturbation wave vector tilt angle k_* and the ion kinetic parameter z_i at $t = t_\mathrm{bound} = 4.3$ and $\varepsilon_\mathrm{R} = (\varepsilon_\mathrm{R})_\mathrm{bound} = 6.5 \times 10^{-5}$ for the 'direct kinetic Alfvén-like' waves.

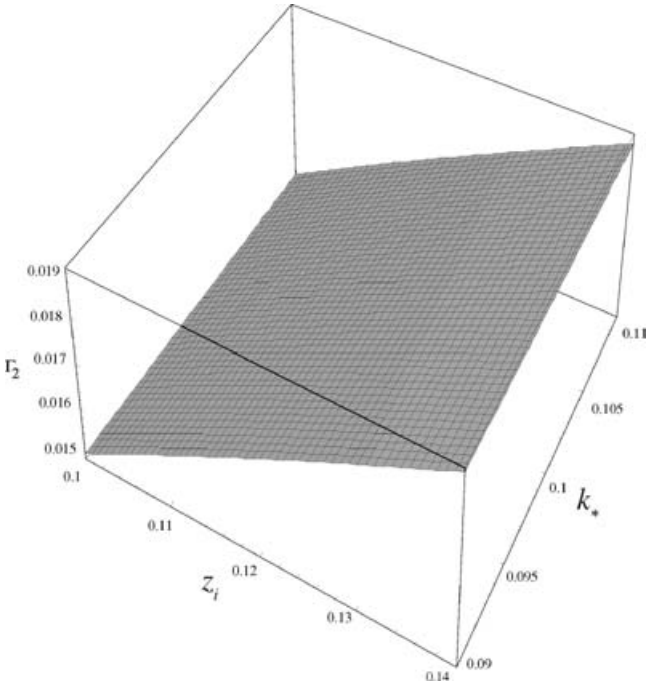


Figure 6. The reduced growth rates $\Gamma_k = \Gamma(\Omega_k)$ $(k = 1, 2, 3, 4)$ of the four types of the waves with $\Omega = \Omega_k$ as the functions of the perturbation wave vector tilt angle k_* and the ion kinetic parameter z_i at $t = t_\mathrm{bound} = 4.3$ and $\varepsilon_\mathrm{R} = (\varepsilon_\mathrm{R})_\mathrm{bound} = 6.5 \times 10^{-5}$ for the 'inverse kinetic Alfvén-like' waves.

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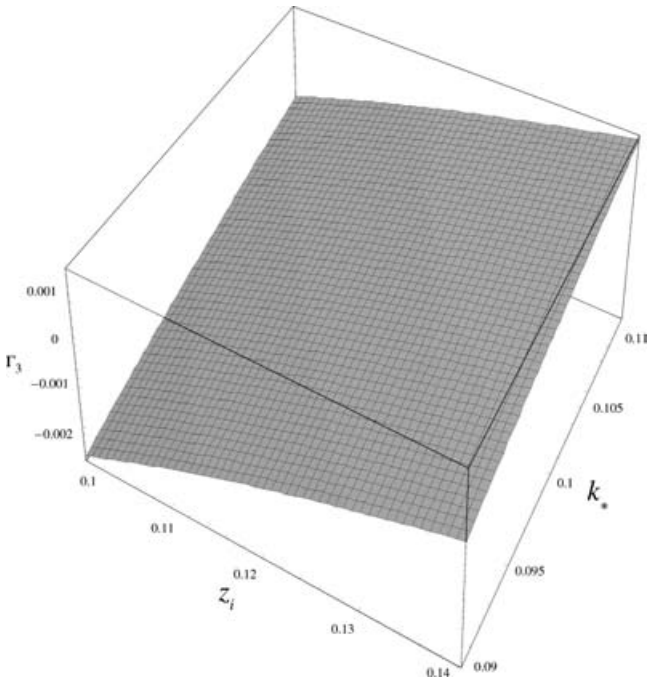


Figure 7. The reduced growth rates $\Gamma_k = \Gamma(\Omega_k)$ $(k = 1, 2, 3, 4)$ of the four types of the waves with $\Omega = \Omega_k$ as the functions of the perturbation wave vector tilt angle k_* and the ion kinetic parameter z_i at $t = t_\mathrm{bound} = 4.3$ and $\varepsilon_\mathrm{R} = (\varepsilon_\mathrm{R})_\mathrm{bound} = 6.5 \times 10^{-5}$ for the 'inverse' SMAWs with $\Omega = \Omega_3$.

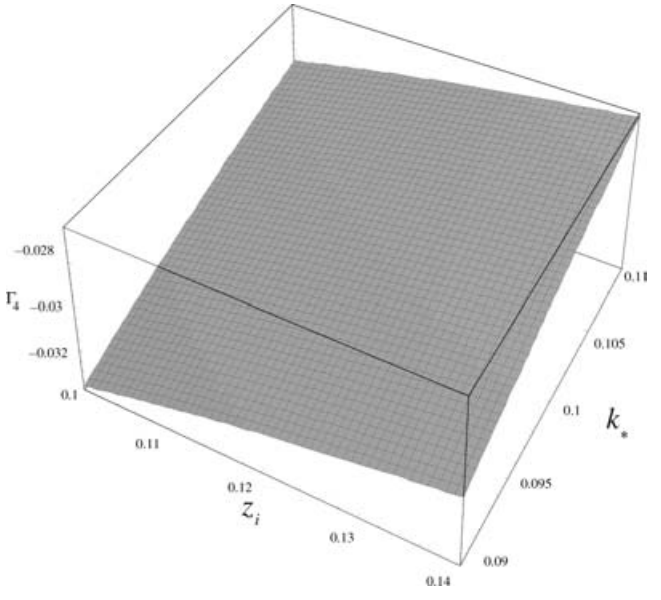


Figure 8. The reduced growth rates $\Gamma_k = \Gamma(\Omega_k)$ $(k = 1, 2, 3, 4)$ of the four types of the waves with $\Omega = \Omega_k$ as the functions of the perturbation wave vector tilt angle k_* and the ion kinetic parameter z_i at $t = t_\mathrm{bound} = 4.3$ and $\varepsilon_\mathrm{R} = (\varepsilon_\mathrm{R})_\mathrm{bound} = 6.5 \times 10^{-5}$ for the 'direct' SMAWs with $\Omega = \Omega_4$.

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see, only two of the four waves, which correspond to the roots Ω_k $(k = 1, 2, 3, 4)$, can be generated in the investigated situation due to the rise and development of corresponding instabilities, namely the 'Ω_2-wave' and 'Ω_3-wave'. The 'Ω_1-wave' and 'Ω_4-wave' have the negative growth rates Γ_1 and Γ_4, so, as a result they damp. The growth rate Γ_2 of 'modified KAW' Ω_2 practically always sufficiently exceeds the growth rate Γ_3 of the 'Ω_3-wave'. However, numerical simulation shows that the value of Γ_2 does not depend on the value of ε_{R}. This means that the modified KAW instability rises and develops only due to the existence of magnetic field $\mathbf{B}_0$ in the loop and does not 'feel' the changes of the value of the large-scale electric field amplitude $\mathbf{E}_0(t)$. At the same time the instability of the 'Ω_3-wave' has clearly expressed threshold character and rises at certain values of z_{i} and k_* only when $t = t_{\mathrm{bound}}$ and $\varepsilon_{\mathrm{R}} = (\varepsilon_{\mathrm{R}})_{\mathrm{bound}}$ (see (4.13) and (4.14)). This means that in the framework of HPR-theory (Heyvaerts et al. 1977) this instability can be considered as an 'indicator' of changes of the quasi-static field amplitude and, in principle, it is possible to make use of it for the short-time prediction of a flare in the loop structures. The small values of Γ_2 and Γ_3 point to the fact that we investigate the real wave process, when many wave periods are generated during the time of instability development.

The main results of present work can be briefly summarized as follows.

1. Consideration of the terms that describe the ion thermal motion (which are small by definition) in the expressions for the $\mathrm{Re}\,\varepsilon_{\parallel}^0$ and $\mathrm{Im}\,\varepsilon_{\parallel}^0$ has led to the sufficient complication of the expressions for the coefficients of the MDR and has made it impossible to obtain the solutions of the MDR in analytical form. So, we had to use the numerical calculations based on the Euler method. However, this way provides us with the possibility of consideration of the extremely small values (up to the zero ones) of ε_{R}. It allows to obtain the threshold ("Boundary") values $(\varepsilon_{\mathrm{R}})_{\mathrm{bound}}$ for ε_{R} and $\mathrm{t}_{\mathrm{bound}}$ for the ratio of the electron and ion temperatures. In the present case, there are not instabilities that rise at $t = 1$. The first instability only appears at the corresponding value of $t = t_{\mathrm{bound}} > 1$.

2. For the extremely small values of ε_{R}, which do not differ greatly from the 'equilibrium' and consideration of ion thermal motion, the four 'kinetic Alfvén wave-like' roots of MDR, which exist for the 'extremely large' values of ε_{R} (Kryshtal 2002), transform into two pairs. The first represents the kinetic Alfvén waves, modified by the weak quasi-static large-scale field $\mathbf{E}_0(t)$, the pair Coulomb collisions and drift motions. The second pair represents the waves from the range of the SMAWs, which are modified by the same above-mentioned effects. The frequencies of these waves have non-traditional (in comparison with the KAWs) dependence on kinetic ion parameter z_{i} and do not depend on parameter k_*.

3. The 'direct' KAW ('Ω_1-wave' with $k_z > 0$) and the 'direct Ω_4-wave' (with $k_z > 0$) have negative growth rates Γ_1 and Γ_4, so, naturally, they damp. At the same time the 'inverse' KAW ('Ω_2-wave' with $k_z < 0$) as well as the 'inverse Ω_3-wave' (with $k_z < 0$) have positive growth rates, hence they are generated during the initial (linear) stage of the process of instability development. The growth rate Γ_2 does not depend on the value of ε_{R} and corresponding instability exists even at $t = 1$. The main reason of this instability is the magnetic field $\mathbf{B}_0$ of the 'old' loop and it does not 'feel' the changes of the electric field amplitude. The 'Ω_3-wave' instability has clearly expressed threshold character and Γ_3 becomes

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positive at certain values of z_i and k_* when $t \geqslant t_{bound}$ and $(\varepsilon_R) \geqslant (\varepsilon_R)_{bound}$. The main reason of this instability is the increase of the electric field amplitude $\mathbf{E}_0(t)$. In this situation consideration of the drift effects due to the presence of spatial inhomogeneity of temperature (and also pair Coulomb collisions) results in the appearance of the corrections in the dispersion law. In a sense it seems unusual, because at $\varepsilon_R = 0$, $t = 1$ and $\nu_{ei} = 0$ (Michailovsky 1963, 1975) the lowest threshold among all of the roots of the DR has 'very slow' drift Alfvénic wave. In the present case the lowest threshold has 'Ω_3-wave' from the range of SMAWs. The drift effect in this wave does not dominate.

4. Some characteristics of low-frequency instabilities remain unchanged under transition from the relatively 'large' values of ε_R to the extremely small values and taking into account the ion thermal motion. As in our previous paper (Kryshtal 2002) the investigated instabilities can arise only at some specific values of plasma density and temperature, i.e. on a certain floor of the loop structure with a certain plasma equation of state. The instabilities at $\Omega = \Omega_2$ and $\Omega = \Omega_3$ are both of marginal type (Kryshtal and Kucherenko 1995), because the maximums of the growth rates are reached at the borders of the intervals.

5. Due to the extremely high qualities of the investigated waves (very small values of Γ_2 and Γ_3) the above-mentioned instabilities at $\Omega = \Omega_2$ and $\Omega = \Omega_3$ are not as fast as the early investigated instabilities (Kryshtal 2002). Some of the times of instability development for $\Omega = \Omega_3$ can reach the values $\tau_{inst} \approx 0.1$–0.2 s.

6. The obtained values of $(\varepsilon_R)_{bound}$, t_{bound} and Γ_3 show to the fact that the clearly expressed wave process, i.e. the generation of low-frequency waves that are close to the SMAWs, can arise at a very early stage of development of preflare situation in the loop structure in solar ARs. Of course, it is not enough for the short-time prediction of a flare, but exactly it is one of the most important necessary conditions of such prediction.

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