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Weakly nonlinear transport in warm multicomponent dusty magnetoplasmas

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Abstract. A general nonlinear theory for low-frequency electromagnetic field generation due to the action of non-uniform and non-stationary high-frequency electromagnetic radiation (ponderomotive force) in uniform multicomponent dusty magnetoplasmas is developed. The anisotropic thermal pressure of all species in slow and fast oscillations is taken into account. This theory permits us to consider the nonlinear action of all waves that can exist in such plasmas. As an example, the equations for dust particle velocities arising due to nonlinear electric fields induced by electromagnetic cyclotron waves travelling along an external magnetic field are obtained. The dust grains are considered to be magnetized as well as unmagnetized. Various regimes of dust grain dynamics depending on the spatio-temporal variation of a fast wave amplitude are discussed. It is shown that induced nonlinear electric fields can have both an electrostatic and an electromagnetic nature. In some cases these fields across the background magnetic field in warm plasmas are considerably larger than those in cold plasmas. The longitudinal nonlinear electric field remains the same in warm and cold plasmas in its order of magnitude. Regimes are found for which the transverse acceleration of dust grains is much larger than the longitudinal one.

1. Introduction

The presence of charged dust grains is a common phenomenon in various low-temperature laboratory and technology plasmas, in the Earth's lower ionosphere and in astrophysical objects such as planetary and stellar atmospheres, planetary rings, the solar system, cometary tails, interstellar clouds, intergalactic media, dust discs in the galactic nuclei [1–4]. Radio-location data [5] and direct rocket probe measurements [6] have revealed meteoric dust in the Earth's mesosphere (at heights from 80 up to 100 km). The spatial distribution of dust grains is non-uniform. For example, some rings show radial and azimuthal optical depth modulations [2].

Dusty plasmas are an ensemble of micron-sized extremely massive (approximately 10^9 times heavier than protons) and highly charged (up to several thousand times that of the electron charge) dust grains usually immersed in a fully or partially ionized low-temperature electron–ion plasma. Dust particles may have negative or positive charges due to various processes: absorption of electrons and ions from the background plasma, the presence of electron and ion currents, photoemission, thermionic emission, secondary ionization and so on [7–12]. It has been

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shown by Chow et al. [7] that the smaller (larger) sized particles must be charged positively (negatively) due to the secondary ionization effect.

The presence of dust grains modifies considerably the dispersion relations for the electron–ion plasma because of the change in the quasi-neutrality condition, dust charge variations and dust dynamics [13–19]. The important effects related to dust charging are also the change in the distributions of electrons and ions taking part in the charging process, and the definition of the background state of dusty plasmas [20]. Interest in the physics of collective processes in dusty plasmas has grown considerably owing to the discovery of dust acoustic waves [13] and dust plasma crystals [10, 21, 22].

It is obvious that taking dust grain dynamics into account is necessary for an adequate description of plasma phenomena in the laboratory, in technological applications, and in geophysical and astrophysical plasmas. In particular, dust grain movements influence the cosmic microwave background [23, 24]. This is a consequence of interaction of dust with the background electron–ion plasma and electromagnetic fields.

One of the important interaction processes between particles and waves is the mass transport and particle energization due to ponderomotive forces. There are many works that deal with these and related processes in the laboratory (e.g. [25–28]), and in magnetospheric ([29] and references therein, [30–33]) and solar [34–36] electron–ion plasmas. The investigations have shown that the ponderomotive forces of finite-amplitude waves can considerably influence the plasma energy and its distribution. Lately, some theoretical studies on this subject have been continued for dusty plasmas (e.g. [37–39]). Accelerated dust winds and dust fluxes are observed in the cosmic environment and late-type stars [1, 4, 40]. Such grains can in turn generate low-frequency dust waves. Experiments have been carried out on the interaction of radar beams with the mesospheric dusty plasma [5, 41]. The ponderomotive force of high-frequency waves can be used for various purposes in laboratory devices [42], and, in particular, for the removal of impurities (dust grains) from the plasma. Thus, the study of the dust dynamics in high-frequency electromagnetic fields is of importance both for practical uses and in explaining natural phenomena.

If the ponderomotive force arises owing to the high-frequency electron or ion electromagnetic waves, dust grains experience the action of the low-frequency electric field induced by the ponderomotive force. The differential acceleration of dust grains along and across the external magnetic field under the action of that electric field has been considered in [37–39] for the case of a high-frequency wave with quasi-stationary amplitude. The induced electric field was assumed to be an electrostatic one. In nature, however, waves generally exist in non-uniform and non-stationary media. Therefore their amplitudes change dynamically. Spatial inhomogeneity of the wave amplitude can originate from inhomogeneity of the medium or nonlinear effects. In experiments external radiation also has spatio-temporal variations. Strong ponderomotive effects occur in standing Alfvén waves, which are observed in the Earth's magnetosphere [43]. In the general case the relations between the spatial and temporal changes of the wave amplitude may be various, i.e. not only the quasi-stationary regime can be realized. This general case is considered here.

In the present paper we develop a general weakly nonlinear theory for low-frequency electromagnetic field generation under the action of high-frequency electromagnetic radiation (ponderomotive force) in warm uniform multicomponent

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dusty plasmas in an external magnetic field. This theory permits us to consider the nonlinear action of all waves that can exist in such plasmas. The low-frequency electric field generated by the ponderomotive force is calculated from the full system of equations describing the warm multicomponent plasma. We will show that the induced nonlinear field has an electromagnetic nature in the general case. As a concrete example for high-frequency waves we consider electromagnetic cyclotron waves travelling along the background magnetic field. We derive the equations for dust particle velocities arising due to the action of the nonlinear electric field. Magnetized as well as unmagnetized dust grains are considered. We discuss various regimes of the dust grain dynamics depending on the spatio-temporal variations of the fast wave amplitude.

In Sec. 2 we derive the general equation for induced nonlinear electric fields in multicomponent plasmas, taking into account the anisotropic thermal pressure of all species in slow and fast oscillations. Then, we find the solution of this equation for the case when the high-frequency waves are circularly polarized electromagnetic waves travelling along the background magnetic field. In Sec. 3 we obtain the expressions for the velocities of dust grains in various regimes. The equation for the evolution of the dust number density is also given. In Sec. 4 we consider the nonlinear action of the electron and ion cyclotron waves separately. In Sec. 5 the results obtained are summarized.

2. Basic equations and expressions

2.1. General relations

We start with the equations of motion

$$\frac{\partial \mathbf{v}_j}{\partial t} + \mathbf{v}_j \cdot \nabla \mathbf{v}_j = -\frac{T_{aj} \nabla n_j}{m_j n_j} + \frac{e_j}{m_j}\mathbf{E} + \frac{q_j}{m_j c}\mathbf{v}_j \times (\mathbf{B}_0 + \mathbf{B}), \tag{1}$$

continuity

$$\frac{\partial n_j}{\partial t} + \nabla \cdot n_j \mathbf{v}_j = 0 \tag{2}$$

and Maxwell's equations

$$\nabla \times \mathbf{E} = -\frac{1}{c}\frac{\partial \mathbf{B}}{\partial t}, \tag{3}$$

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

where the current density is $\mathbf{j} = \sum_j q_j n_j \mathbf{v}_j$. Here q_j and m_j are the charge and mass of species j, where $j =$ e, i, d denotes electrons, ions and dust grains, respectively, $\mathbf{v}_j$ is the hydrodynamic velocity, n_j is the number density, T_{aj} is the anisotropic temperature (see below), $\mathbf{E}$ and $\mathbf{B}$ are the perturbed electric and magnetic fields, $\mathbf{B}_0$ is the background magnetic field and c is the speed of light in vacuum.

Below we will be interested in non-oscillatory nonlinear movements of particles induced by electromagnetic waves. It is obvious that a nonlinear electric field must be found for this purpose. Let us consider, at first, equations (1)–(4) up to the first order in the wave amplitude. We suppose here that the background magnetic field $\mathbf{B}_0$ and number density n_{j0} are uniform, and the temperature T_{aj} is constant. Let $\mathbf{B}_0$ be directed along the z-axis. Then, index $a = (\perp, \parallel)$, where the signs '$\perp$' and '$\parallel$'

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denote the directions perpendicular to and along $\mathbf{B}_0$, respectively. For simplicity, all variable values are considered to depend on the coordinates y and z only. In this case in the linear approximation we obtain the following expressions for the velocity and number density perturbations in the Fourier representation

$$\mathbf{v}_{j1k} = \frac{q_j}{m_j}\frac{i\omega}{\omega^2-\omega_{cj}^2}\mathbf{L}_{jk}(\mathbf{E}_{1k}), \tag{5}$$

where

$$L_{jxk}(\mathbf{D}) = (1+c_j)D_x + i\frac{\omega_{cj}}{\omega}\left(1+\frac{\omega^2}{\omega_{cj}^2}c_j\right)D_y + i\frac{\omega_{cj}}{\omega}d_{\perp j}D_z,$$

$$L_{jyk}(\mathbf{D}) = \left(1+\frac{\omega^2}{\omega_{cj}^2}c_j\right)\left(-i\frac{\omega_{cj}}{\omega}D_x + D_y\right) + d_{\perp j}D_z, \tag{6}$$

$$L_{jzk}(\mathbf{D}) = d_{\parallel j}\left(-i\frac{\omega_{cj}}{\omega}D_x + D_y\right) + \frac{\omega^2-\omega_{cj}^2}{\omega^2(1-\delta_{jk})}D_z$$

and

$$n_{j1k} = \frac{q_j n_{j0}}{m_j}\frac{i}{1-\delta_{jk}}\left[\frac{k_y}{\omega^2-\omega_{cj}^2}\left(-i\frac{\omega_{cj}}{\omega}E_{1xk}+E_{1yk}\right)+\frac{k_z}{\omega^2}E_{1zk}\right]. \tag{7}$$

Here $\omega_{cj} = q_j B_0/m_j c$ is the cyclotron frequency, $k = \{\mathbf{k}, \omega\}$, where $\mathbf{k}$ is the wave vector and ω is the frequency, the index '1' denotes the linear approximation. The following notation is introduced in (6) and (7):

$$c_j = \frac{T_{\perp j}}{m_j}\frac{\omega_{cj}^2 k_y^2}{\omega^2(\omega^2-\omega_{cj}^2)(1-\delta_{jk})},$$

$$d_{\perp(\parallel)j} = \frac{T_{\perp(\parallel)j}}{m_j}\frac{k_y k_z}{\omega^2(1-\delta_{jk})},$$

$$\delta_{jk} = \frac{T_{\parallel j}}{m_j}\frac{k_z^2}{\omega^2}.$$

We consider the case when the inequality $1 \gg (T_{\perp j}/m_j)k_y^2/(\omega^2-\omega_{cj}^2)$ is fulfilled.

Using the expressions (5) and (6), we can find the linear current $\mathbf{j}_{1k}$. We suppose that the background velocity $\mathbf{v}_{j0}$ is equal to zero. Substituting this current into Maxwell's equations, (3) and (4), written in the linear approximation in the Fourier representation, we obtain the equation for the wave electric field,

$$\hat{A}_k\mathbf{E}_{1k} = 0, \tag{8}$$

where the matrix $\hat{A}_k$ is

$$\hat{A}_k = \begin{vmatrix} n^2-\varepsilon_1+c_1 & -i(\varepsilon_2-c_2) & id_{\perp 2} \\ i(\varepsilon_2-c_2) & n_z^2-\varepsilon_1 & -n_y n_z + d_{\perp 1} \\ -id_{\parallel 2} & -n_y n_z + d_{\parallel 1} & n_y^2-\varepsilon_3 \end{vmatrix}. \tag{9}$$

Here $n^2 = n_y^2 + n_z^2$, $n_y = k_y c/\omega$, $n_z = k_z c/\omega$, and the other notation used is

$$\varepsilon_1 = 1 - \sum_j a_j,$$

$$\varepsilon_2 = -\sum_j b_j,$$

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$$\varepsilon_3 = 1 - \sum_j \frac{\omega_{\mathrm{p}j}^2}{\omega^2(1-\delta_{jk})},$$

$$c_1 = \sum_j a_j c_j,$$

$$c_2 = \sum_j \frac{\omega^2}{\omega_{\mathrm{c}j}^2} b_j c_j,$$

$$d_{\perp(\parallel)1} = \sum_j a_j d_{\perp(\parallel)j},$$

$$d_{\perp(\parallel)2} = \sum_j b_j d_{\perp(\parallel)j},$$

where

$$a_j = \frac{\omega_{\mathrm{p}j}^2}{\omega^2 - \omega_{\mathrm{c}j}^2},$$

$$b_j = \frac{\omega_{\mathrm{p}j}^2 \omega_{\mathrm{c}j}}{\left(\omega^2 - \omega_{\mathrm{c}j}^2\right)\omega}.$$

Here $\omega_{\mathrm{p}j} = (4\pi n_{j0} q_j^2/m_j)^{1/2}$ is the plasma frequency.

Let us now consider the equation of motion (1) up to the second order in the wave amplitude. We introduce the nonlinear force

$$\mathbf{F}_{j2} = \frac{T_{aj}\nabla n_{j1}^2}{2q_j n_{j0}^2} - \frac{m_j}{q_j}\mathbf{v}_{j1}\cdot\nabla\mathbf{v}_{j1} + \frac{1}{c}\mathbf{v}_{j1}\times\mathbf{B}_1. \tag{10}$$

The ponderomotive force is $\langle\mathbf{F}_{j2}\rangle$, where (and below) the angled brackets $\langle\cdots\rangle$ denote a time-averaging over fast oscillations. The index '2' means that the values are of second order. Then, we obtain for the velocity $\mathbf{v}_{j2}$ the same equation as that for $\mathbf{v}_{j1}$, if we make the substitutions: '1'→ '2', $\mathbf{E}_1 \to \mathbf{E}_2 + \mathbf{F}_{j2}$. The continuity equation (2) has the following solution up to the second order in the wave amplitude:

$$n_{j2k} = n_{j0}\frac{\mathbf{k}\mathbf{v}_{j2k}}{\omega} + \frac{\mathbf{k}(n_{j1}\mathbf{v}_{j1})_k}{\omega}. \tag{11}$$

The total nonlinear current $\mathbf{j}_2$ is equal to $\mathbf{j}_{21} + \mathbf{j}_{22}$, where

$$\mathbf{j}_{21} = \sum_j q_j n_{j1}\mathbf{v}_{j1}$$

and

$$\mathbf{j}_{22} = \sum_j q_j n_{j0}\mathbf{v}_{j2}.$$

Calculating $\mathbf{v}_{j2}$ and substituting $\mathbf{j}_2$ into Maxwell's equations, we obtain the equation for the determination of the nonlinear electric field $\mathbf{E}_{2k}$

$$\hat{A}_k\mathbf{E}_{2k} = \mathbf{C}_k. \tag{12}$$

Here the vector $\mathbf{C}_k$ equals

$$\mathbf{C}_k = \frac{4\pi i}{\omega}\mathbf{j}_{21k} - \sum_j a_j[\mathbf{L}_{jk}(\mathbf{F}_{j2k}) - i\mathbf{G}_{jk}], \tag{13}$$

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where

$$\mathbf{G}_{jk} = \frac{T_{\perp j}}{q_j n_{j0}} \frac{\mathbf{k}(n_{j1}\mathbf{v}_{j1})_k}{\omega(1-\delta_{jk})} \left(i\frac{\omega_{cj}}{\omega} k_y;\ k_y;\ \frac{T_{\parallel j}}{T_{\perp j}} \frac{\omega^2 - \omega_{cj}^2}{\omega^2} k_z \right). \tag{14}$$

All of the expressions obtained above have a general form and are suitable for various waves that can exist in warm uniform multicomponent magnetoplasmas. Note, that using, for example, (3) and (12), we can find the induced magnetic moment of a medium in various cases.

2.2. *Circularly polarized waves*

Below, as an example, we will consider the generation of the nonlinear electric field $\mathbf{E}_{2k}$ by circularly polarized electromagnetic waves travelling along the background magnetic field $\mathbf{B}_0$. For these waves, as follows from (8) and (9), the following conditions must be fulfilled:

$$\begin{gathered} n_z^2 \gg n_y^2, c_1, c_2, \\ \varepsilon_3 \gg n_y^2, \frac{n_y}{n_z}(d_\perp + d_\parallel), \frac{1}{n_z^2} d_\perp d_\parallel, \end{gathered} \tag{15}$$

where $d_{\perp(\parallel)} = d_{\perp(\parallel)1} + \sigma(\omega/|\omega|) d_{\perp(\parallel)2}$. Note, that in (15) the value of δ_{jk} is arbitrary. For the electromagnetic electron and ion cyclotron waves considered below we may write the second inequality (15) in the following (sufficient) form: $\varepsilon_3 \gg \alpha_j^2 n_y^2$, where $\alpha_e = 1$ and $\alpha_i = \omega_i/\omega$. The inequality $n_z^2 \gg n_y^2$ is sufficient for neglecting c_1 and c_2, when $T_{\perp j} \sim T_{\parallel j}$. Then, the dispersion relation takes the form

$$n_z^2 = \varepsilon_1 + \sigma \frac{\omega}{|\omega|} \varepsilon_2, \tag{16}$$

and the relation between the electric field components is

$$E_{1xk} = i\sigma \frac{\omega}{|\omega|} E_{1yk}. \tag{17}$$

Here $\sigma = \pm 1$ and the sign $|\ |$ denotes the absolute meaning of the magnitude. When $\sigma = +1(-1)$ the wave is the left- (right-) polarized one. Note that the second inequality of (15) means that the contribution of particle oscillations along the background magnetic field to the dispersion relation (16) can be neglected ($E_{1zk} \approx 0$).

Let us now consider the case where the nonlinear plasma movements vary slowly with respect to the wave frequencies. For this purpose we average (12) over fast oscillations. Suppose that the wave amplitudes depend on the coordinates y, z and the time t, and that the wave phases do not depend, for simplicity, on the coordinate y. Let the following conditions for the slow movements be fulfilled:

$$\begin{gathered} \varepsilon_3 \gg n_y^2, \\ n_y n_z \gg d_{\perp(\parallel)1}, \\ n^2 - \varepsilon_1 + c_1 \gg d_{\perp 2} d_{\parallel 2}/\varepsilon_3. \end{gathered} \tag{18}$$

The last two inequalities of (18) are, roughly, justified under the sufficient condition $c_{Ai}^2 \gg c_s^2$, where $c_{Ai} = B_0/(4\pi n_{i0} m_i)^{1/2}$ is the ion Alfvén velocity, $c_s = [\sum_j n_{j0}(T_{\perp j} + T_{\parallel j})/n_{i0} m_i]^{1/2}$ is the ion sound velocity. Then, we obtain the solution of (12) in

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the form

$$\begin{aligned}
\varepsilon\langle E_{2x}\rangle_k &= \left(n_z^2-\varepsilon_1\right)\langle C_x\rangle_k + i\left(\varepsilon_2 - c_2 + d_{\perp 2}\frac{n_y n_z}{\varepsilon_3}\right)\langle C_y\rangle_k \\
&\quad + \frac{i}{\varepsilon_3}\left[\left(n_z^2-\varepsilon_1\right)d_{\perp 2} - (\varepsilon_2-c_2)n_y n_z\right]\langle C_z\rangle_k, \\
\varepsilon\langle E_{2y}\rangle_k &= -i\left(\varepsilon_2 - c_2 + d_{\parallel 2}\frac{n_y n_z}{\varepsilon_3}\right)\langle C_x\rangle_k + (n^2-\varepsilon_1+c_1)\langle C_y\rangle_k \\
&\quad + \frac{1}{\varepsilon_3}[(\varepsilon_2-c_2)d_{\perp 2} - n_y n_z(n^2-\varepsilon_1+c_1)]\langle C_z\rangle_k, \\
\varepsilon_3\langle E_{2z}\rangle_k &= -\langle C_z\rangle_k - id_{\parallel 2}\langle E_{2x}\rangle_k - n_y n_z\langle E_{2y}\rangle_k,
\end{aligned} \tag{19}$$

where

$$\varepsilon = (n^2-\varepsilon_1+c_1)\left(n_z^2-\varepsilon_1\right) - (\varepsilon_2-c_2)\left[\varepsilon_2 - c_2 + (d_{\perp 2}+d_{\parallel 2})\frac{n_y n_z}{\varepsilon_3}\right].$$

Note, that the wave vector $\mathbf{k}$ and frequency ω in (19) characterize the inhomogeneity and non-stationarity of the wave amplitude.

When calculating the nonlinear values $\langle n_{j1}\mathbf{v}_{j1}\rangle$ and $\langle \mathbf{F}_{j2}\rangle$ we assumed that $c_j \ll 1$ for the fast oscillations. Then, the expression for the transverse velocity $\mathbf{v}_{j1\perp}$ does not contain the thermal pressure. The longitudinal electric field E_{1zk} and velocity v_{j1zk} are related to the component E_{1yk} as follows:

$$\begin{aligned}
E_{1zk} &= -\frac{n_y n_z}{\varepsilon_3}(1-\alpha_k)E_{1yk}, \\
v_{j1zk} &= -\frac{q_j}{m_j}\frac{i}{\omega}\frac{n_y n_z}{\varepsilon_{4j}}E_{1yk},
\end{aligned} \tag{20}$$

where

$$\alpha_k = \frac{1}{n_z^2}\sum_j a_j\left(1+\sigma\frac{\omega_{cj}}{|\omega|}\right)\frac{\delta_{jk}}{1-\delta_{jk}},$$

$$\varepsilon_{4j} = \varepsilon_3\frac{1-\delta_{jk}}{1-\alpha_k-\delta_{jk}\varepsilon_3 n_z^{-2}|\omega|(|\omega|-\sigma\omega_{cj})^{-1}}.$$

It is interesting to note that $\varepsilon_{4e} = \varepsilon_{3\mathrm{cold}} = -\omega_{pe}^2/\omega^2$ both for $\delta_{ek} \ll 1$ and $\delta_{ek} \gg 1$ (when $\delta_{ik} \ll 1$). Using (5)–(7), (10), (13), (14), (17) and (20), we find the nonlinear 'driving force' $\langle \mathbf{C}\rangle$. We do not take into account the double derivatives of the wave amplitude, which give a small contribution. Then, the expressions for the Fourier components of $\langle \mathbf{C}\rangle$ take the form

$$\begin{aligned}
\langle C_x\rangle_k &= \frac{k_y}{\omega}\sum_j \frac{q_j}{2m_j\omega_{cj}}\frac{\omega_{pj}^2}{(\sigma\omega'-\omega_{cj})^2} \\
&\quad \times\left[1+\frac{\delta_{jk}}{1-\delta_{jk}}\frac{T_{\perp j}}{T_{\parallel j}}\frac{(\sigma\omega'-\omega_{cj})}{\sigma\omega'}\left(1+\frac{\omega_{cj}}{\sigma\omega'-\omega_{cj}}\frac{k_z'\omega}{k_z\omega'}\right)\right]\left\langle \mathbf{E}_1^2\right\rangle_k, \\
\langle C_y\rangle_k &\simeq 0, \\
\langle C_z\rangle_k &= \frac{ik_z}{\omega^2}\sum_j \frac{q_j}{2m_j}\frac{\omega_{pj}^2}{(1-\delta_{jk})\sigma\omega'(\sigma\omega'-\omega_{cj})}\left(1+\frac{\omega_{cj}}{\sigma\omega'-\omega_{cj}}\frac{k_z'\omega}{k_z\omega'}\right)\left\langle \mathbf{E}_1^2\right\rangle_k.
\end{aligned} \tag{21}$$

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Here and below a prime denotes the wave number and frequency of the high-frequency wave. The frequency ω' (> 0) is assumed to be the definite value. The wave number k'_z is determined from (16). We assume that $\omega^2 \ll \omega_{cj}^2$ for particle species which participate in fast oscillations. Dust grains are considered to take no part in fast oscillations, i.e. $\omega'^2 \gg \omega_{pd}^2, \omega_{cd}^2$. Note, that the second term in the square brackets in the expression for $\langle C_x \rangle_k$ is the new term we have obtained, which appears in warm plasmas.

3. Nonlinear transport of dust grains

Below, we consider various regimes of the spatio-temporal inhomogeneity of the wave amplitude, and the cases of magnetized and unmagnetized dust grains. But we suppose everywhere the dust grains to be cold ($1 \gg c_d, d_{\perp(\parallel)d}, \delta_{dk}$) and the electrons to be warm ($\delta_{ek} \gg 1$) in the time-averaged nonlinear movements. Thus $k_z \neq 0$. This is a good condition for longitudinally inhomogeneous media, wave packets and short impulses. The resonance cases when $\varepsilon = 0$ and/or $\varepsilon_3 = 0$ are of interest, but we do not consider these cases here.

3.1. Magnetized dust grains

In this section we account for the case when dust grains are magnetized: $\omega^2 \ll \omega_{cd}^2$. Then, the quantities ε_1 and ε_2 are equal,

$$\varepsilon_1 = \frac{c^2}{c_A^2},$$
$$\varepsilon_2 = \sum_j \frac{\omega_{pj}^2 \omega}{\omega_{cj}^3},$$

where $c_A = B_0/(4\pi \sum_j n_{j0} m_j)^{1/2}$ is the Alfvén velocity. We see that $\varepsilon_1 \gg \varepsilon_2$.

3.1.1. Stationary regime. Suppose that $\omega^2 \ll k_z^2 c_A^2$. Then, we obtain from (19)

$$\langle E_{2x} \rangle_k = \frac{1}{n^2} \langle C \rangle_k = \frac{1}{n^2} \left(\langle C_x \rangle_k + i \frac{d_{\perp 2}}{\varepsilon_3} \langle C_z \rangle_k \right),$$
$$\langle E_{2y} \rangle_k = -i \frac{\varepsilon_2 - c_2}{n_z^2} \langle E_{2x} \rangle_k + \frac{n_y}{n_z} \langle E_{2z} \rangle_k, \tag{22}$$
$$\langle E_{2z} \rangle_k = -\frac{1}{\varepsilon_3} \langle C_z \rangle_k.$$

Here $\varepsilon_2 - c_2$ may be substituted by $\varepsilon_{2d} = \omega_{pd}^2 \omega / \omega_{cd}^3$. Note that the thermal terms in the expression for $\langle C \rangle_k$ are of the same order of magnitude (for more detail see (30)). In the expressions for $\langle C_x \rangle_k$ and $\langle C_z \rangle_k$ the non-stationary terms proportional to ω must be neglected.

Let us now find the equations for the dust grain velocity $\langle \mathbf{v}_{d2} \rangle$. In (5) we must substitute the index '1' by '2'. Then, using (6) and (22), we obtain

$$\langle v_{d2x} \rangle_k = \frac{c}{B_0} \left(-i \frac{\omega}{\omega_{cd}} \langle E_{2x} \rangle_k + \frac{k_y}{k_z} \langle E_{2z} \rangle_k \right),$$
$$\langle v_{d2y} \rangle_k = -\frac{c}{B_0} \left(\langle E_{2x} \rangle_k + i \frac{\omega}{\omega_{cd}} \frac{k_y}{k_z} \langle E_{2z} \rangle_k \right), \tag{23}$$
$$\langle v_{d2z} \rangle_k = \frac{q_d}{m_d} \frac{i}{\omega} \langle E_{2z} \rangle_k.$$

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In the expressions for $\langle v_{\mathrm{d}2x(y)k}\rangle$ the field $\langle E_{2y}\rangle_k$ enters connected with $\langle E_{2z}\rangle_k$ (see (22)). If the condition $\omega\omega_{\mathrm{ci}} \gg \mathbf{k}^2 c_{\mathrm{Ai}}^2$, where $\mathbf{k}^2 = k_y^2 + k_z^2$, is fulfilled, i.e. the wave is a little non-stationary, we find from (23)

$$\langle v_{\mathrm{d}2y}\rangle = -\frac{c}{B_0}\langle E_{2x}\rangle,$$
$$\langle v_{\mathrm{d}2x}\rangle \ll \langle v_{\mathrm{d}2y}\rangle,$$

where $\langle E_{2x}\rangle \gg \langle E_{2y}\rangle, \langle E_{2z}\rangle$, i.e. this field has an electromagnetic nature. We see that the transverse nonlinear electric field is generated perpendicular to the transverse wave amplitude inhomogeneity, and the transverse dust grain velocity is directed along the latter. In the opposite case, $\omega\omega_{\mathrm{ci}} \ll \mathbf{k}^2 c_{\mathrm{Ai}}^2$, which may be called the purely stationary regime (see below for the case of warm ions), we obtain

$$\langle v_{\mathrm{d}2x}\rangle_k = \frac{c}{B_0}\frac{k_y}{k_z}\langle E_{2z}\rangle_k,$$
$$\langle v_{\mathrm{d}2y}\rangle \ll \langle v_{\mathrm{d}2x}\rangle.$$

In this case $\langle E_{2y}\rangle_k = (k_y/k_z)\langle E_{2z}\rangle_k \gg \langle E_{2x}\rangle_k$. Thus, the induced electric field has an electrostatic nature only in the purely stationary regime, when the field does not depend on time (see (21) and (22)).

The dust particles move along (across) the z-axis with constant (in time) acceleration (velocity).

The expressions (22) for $\langle E_{2z}\rangle_k$ and (23) for $\langle v_{\mathrm{d}2z}\rangle_k$ are the same in all of the cases considered below.

In the results obtained above ions can be both warm ($\delta_{\mathrm{i}k} \gg 1$) and cold ($\delta_{\mathrm{i}k} \ll 1$). For warm ions $\varepsilon_3 (\gg 1) = (k_z r_{\mathrm{D}})^{-2}$, where $r_{\mathrm{D}} = [T_{\|\mathrm{e}} T_{\|\mathrm{i}}/4\pi e^2 (n_{\mathrm{i}0} T_{\|\mathrm{e}} + n_{\mathrm{e}0} T_{\|\mathrm{i}})]^{1/2}$ is the plasma Debye length ($-q_{\mathrm{e}} = q_{\mathrm{i}} = e$). We suppose in this case that $\omega^2 \gg \omega_{\mathrm{pd}}^2 (k_z r_{\mathrm{D}})^2$, i.e. the frequency ω must be larger than the dust sound frequency. For cold ions $\varepsilon_3 = (k_z r_{\mathrm{De}})^{-2}$, where $r_{\mathrm{De}} = (T_{\|\mathrm{e}}/4\pi n_{\mathrm{e}0} e^2)^{1/2}$. In the expression for $\langle C_z\rangle_k$ it is sufficient in this case to take into account the electron contribution only. Note that the longitudinal electric field $\langle E_{2z}\rangle$ in the warm plasma stays of the same order of magnitude as that in the cold plasma.

The equation for the dust number density evolution here and below takes the form (see (11))

$$\frac{\partial \langle n_{\mathrm{d}2}\rangle}{\partial t} = -n_{\mathrm{d}0}\left(\frac{\partial \langle v_{\mathrm{d}2y}\rangle}{\partial y} + \frac{\partial \langle v_{\mathrm{d}2z}\rangle}{\partial z}\right).$$

3.1.2. Dynamic regime. Here we suppose that $\omega^2 \gg \mathbf{k}^2 c_{\mathrm{A}}^2$. This regime is possible with a weak dependence of the wave amplitude on coordinates: $\omega_{\mathrm{cd}}^2 \gg \mathbf{k}^2 c_{\mathrm{A}}^2$. In this case we obtain from (19)

$$\langle E_{2x}\rangle_k = -\frac{1}{\varepsilon_1 - c_1}\langle C\rangle_k,$$
$$\langle E_{2x}\rangle \gg \langle E_{2y}\rangle,\ \langle E_{2z}\rangle, \tag{24}$$

i.e. the field is an electromagnetic one. The transverse velocities of dust grains are equal

$$\langle v_{\mathrm{d}2y}\rangle = -\frac{c}{B_0}\langle E_{2x}\rangle,$$
$$\langle v_{\mathrm{d}2x}\rangle \ll \langle v_{\mathrm{d}2y}\rangle. \tag{25}$$

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From here and (23) we obtain

$$\frac{\langle v_{\mathrm{d}2y}\rangle}{\langle v_{\mathrm{d}2z}\rangle} \sim \frac{k_y}{k_z}\frac{c_{\mathrm{A}}^2}{c_{\mathrm{Ai}}^2}\frac{\omega_{\mathrm{ci}}}{\omega_{\mathrm{cd}}}.$$

The right-hand side of this equality can be very large. Note that when $k_z/\omega \gg k_z'/\omega'$, we have $\langle E_{2x}\rangle_k \sim \omega^{-1}$ (see (21) and (24)). Thus, the transverse electric field increases with time. As a result, the dust grains move in the y-direction with acceleration.

For warm ions the regime considered is possible when $n_{\mathrm{d}0}m_{\mathrm{d}} \gg n_{\mathrm{i}0}m_{\mathrm{i}}$ as long as $c_{\mathrm{Ai}}^2 \gg c_{\mathrm{s}}^2$ (see (18)). For cold ions the quantity c_1 in (24) may be omitted. The necessary condition for this regime has the form $v_{T_{\parallel}e}^2 \gg (\mathbf{k}^2/k_z^2)c_{\mathrm{A}}^2$, where $v_{T_{\parallel}e} = (T_{\parallel\mathrm{e}}/m_{\mathrm{e}})^{1/2}$.

3.2. *Unmagnetized dust grains*

Now we consider the case where $\omega^2 \gg \omega_{\mathrm{cd}}^2$. Then, the quantities ε_1 and ε_2 take the form

$$\varepsilon_1 = \frac{\omega_{\mathrm{pe}}^2}{\omega_{\mathrm{ce}}^2} + \frac{\omega_{\mathrm{pi}}^2}{\omega_{\mathrm{ci}}^2} - \frac{\omega_{\mathrm{pd}}^2}{\omega^2},$$

$$\varepsilon_2 = -\delta\frac{\omega_{\mathrm{pi}}^2}{\omega\omega_{\mathrm{ci}}},$$

where $\delta = q_{\mathrm{d}}n_{\mathrm{d}0}/en_{\mathrm{i}0}$. The third term on the right-hand side of the expression for ε_1 is less than ε_2 by a factor of $\omega/\omega_{\mathrm{cd}} \gg 1$. We suppose below that $n_z^2 - \varepsilon_1 \gg \omega_{\mathrm{pd}}^2/\omega^2$.

3.2.1. Stationary regime. In the frequency range $k_z^2c_{\mathrm{Ai}}^2 \gg \omega^2, \delta\omega\omega_{\mathrm{ci}}$ the solution for the electric field $\langle\mathbf{E}_2\rangle_k$ has the form (22) (c_2 can be omitted). The expressions for the components of the transverse velocity $\langle\mathbf{v}_{\mathrm{d}2\perp}\rangle_k$ are

$$\begin{aligned}\langle v_{\mathrm{d}2x}\rangle_k &= \frac{q_{\mathrm{d}}}{m_{\mathrm{d}}}\frac{i}{\omega}\left(\langle E_{2x}\rangle_k + i\frac{\omega_{\mathrm{cd}}}{\omega}\frac{k_y}{k_z}\langle E_{2z}\rangle_k\right),\\ \langle v_{\mathrm{d}2y}\rangle_k &= \frac{q_{\mathrm{d}}}{m_{\mathrm{d}}}\frac{i}{\omega}\left[-i\left(\frac{\omega_{\mathrm{cd}}}{\omega} + \frac{\varepsilon_2}{n_z^2}\right)\langle E_{2x}\rangle_k + \frac{k_y}{k_z}\langle E_{2z}\rangle_k\right].\end{aligned} \tag{26}$$

When $\mathbf{k}^2c_{\mathrm{Ai}}^2 \ll \omega\omega_{\mathrm{ci}}$ (this case is suitable for $\delta \ll 1$), we obtain from (26) that $\langle v_{\mathrm{d}2x}\rangle \sim \langle E_{2x}\rangle \gg \langle v_{\mathrm{d}2y}\rangle$ (the field $\langle\mathbf{E}_2\rangle$ is the electromagnetic one). In the opposite case, $\mathbf{k}^2c_{\mathrm{Ai}}^2 \gg \omega\omega_{\mathrm{ci}}$, we have $\langle v_{\mathrm{d}2y}\rangle \sim \langle E_{2z}\rangle \gg \langle v_{\mathrm{d}2x}\rangle$ and $\langle v_{\mathrm{d}2y}\rangle_k/\langle v_{\mathrm{d}2z}\rangle_k = k_y/k_z$ (the field $\langle\mathbf{E}_2\rangle$ is the electrostatic one).

3.2.2. Dynamic regime. In the frequency range $\omega^2 \gg \mathbf{k}^2c_{\mathrm{Ai}}^2, \delta\omega\omega_{\mathrm{ci}}$ (here again $\delta \ll 1$) we obtain for $\langle\mathbf{E}_{2\perp}\rangle_k$ the expression (24) (c_1 can be omitted). The components of $\langle\mathbf{v}_{\mathrm{d}2\perp}\rangle_k$ are equal to

$$\begin{aligned}\langle v_{\mathrm{d}2x}\rangle_k &= \frac{q_{\mathrm{d}}}{m_{\mathrm{d}}}\frac{i}{\omega}\langle E_{2x}\rangle_k,\\ \langle v_{\mathrm{d}2y}\rangle &\ll \langle v_{\mathrm{d}2x}\rangle.\end{aligned} \tag{27}$$

In this case

$$\frac{\langle v_{\mathrm{d}2x}\rangle}{\langle v_{\mathrm{d}2z}\rangle} = \frac{\langle E_{2x}\rangle}{\langle E_{2z}\rangle} \sim \frac{k_y}{k_z}\frac{\omega_{\mathrm{ci}}}{\omega}.$$

This ratio is very large for $k_y \sim k_z$ as long as $\omega \ll \omega_{\mathrm{ci}}$. The regime considered is possible when $v_{T_{\parallel}\mathrm{e}}^2 \gg (\mathbf{k}^2/k_z^2)c_{\mathrm{Ai}}^2$.

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3.2.3. Quasi-dynamic regime. Suppose that the following conditions are fulfilled: $\delta\omega_{\rm ci} \gg \omega \gg \mathbf{k}^2 c_{\rm Ai}^2/\delta\omega_{\rm ci}$. Then, the solution for $\langle \mathbf{E}_{2\perp}\rangle_k$ has the form

$$\begin{gathered}\langle E_{2y}\rangle_k = \frac{i}{\varepsilon_2}\langle C\rangle_k,\\ \langle E_{2x}\rangle \ll \langle E_{2y}\rangle.\end{gathered} \tag{28}$$

The velocity $\langle \mathbf{v}_{\rm d2}\rangle_k$ is equal to

$$\begin{gathered}\langle v_{{\rm d}2y}\rangle_k = \frac{q_{\rm d}}{m_{\rm d}}\frac{i}{\omega}\langle E_{2y}\rangle_k,\\ \langle v_{{\rm d}2x}\rangle \ll \langle v_{{\rm d}2y}\rangle.\end{gathered} \tag{29}$$

Here we have

$$\frac{\langle v_{{\rm d}2y}\rangle}{\langle v_{{\rm d}2z}\rangle} = \frac{\langle E_{2y}\rangle}{\langle E_{2z}\rangle} \sim \frac{k_y}{k_z}\frac{1}{\delta}.$$

When $k_y \sim k_z$ and $\delta \ll 1$ this ratio is large. Note that this regime is possible if $n_{\rm d0} m_{\rm d} \gg n_{\rm i0} m_{\rm i}$, because $\delta\omega_{\rm ci} \gg \omega_{\rm cd}$ here.

4. Transport of dust particles due to electron and ion cyclotron waves

In this section we obtain concrete expressions for $\langle C\rangle_k$ and $\langle C_z\rangle_k$ for the electron and ion cyclotron waves playing the role of high-frequency radiation. In the case of warm ions, $\delta_{ik} \gg 1$, we have from (21) (remember that the expression for $\langle C\rangle_k$ is determined in (22))

$$\begin{gathered}\langle C\rangle_k = \frac{k_y c}{2\omega B_0}\left[\sum_j \frac{\omega_{{\rm p}j}^2}{(\sigma\omega' - \omega_{{\rm c}j})^2} - \lambda_T \frac{\omega_{\rm pe}^2}{(\sigma\omega' - \omega_{\rm ce})(\sigma\omega' - \omega_{\rm ci})}\right]\langle \mathbf{E}_1^2\rangle_k,\\ \langle C_z\rangle_k = -\frac{i}{2k_z}\left[\sum_j \frac{q_j}{T_{\parallel j}}\frac{\omega_{{\rm p}j}^2}{\sigma\omega'(\sigma\omega' - \omega_{\rm cj})}\right]\langle \mathbf{E}_1^2\rangle_k,\end{gathered} \tag{30}$$

where

$$\lambda_T = \frac{T_{\perp\rm e} + T_{\perp\rm i}}{T_{\parallel\rm e} + (1+\delta)T_{\parallel\rm i}}.$$

The first term in the square brackets in the expression for $\langle C\rangle_k$ arises in the cold plasma. The transverse electric field connected with that term causes the dust acceleration considered in [39] in the particular case of unmagnetized dust grains in the quasi-stationary regime (see Sec. 3.2.3). Note that two components, the electrons and ions, must be taken into account in the additional term for the warm plasma proportional to λ_T. For one component the thermal term in $\langle C\rangle_k$ is equal to zero. For cold ions, $\delta_{ik} \ll 1$, and for $k_z'\omega \ll k_z\omega'$ the value λ_T must be substituted by λ_{T1}, where

$$\lambda_{T1} = -\frac{T_{\perp\rm e} + T_{\perp\rm i}}{(1+\delta)T_{\parallel\rm i}}\delta_{ik}.$$

Thus, at $\delta_{ik} \sim 1$ we have $\lambda_T \sim |\lambda_{T1}|$. For the case $\delta_{ik} \ll 1$ and $k_z'\omega \gg k_z\omega'$ the contribution of the thermal term to $\langle C\rangle_k$ is small. Note that the longitudinal electric field $\langle E_z\rangle$ in the warm plasma is the same in terms of order of magnitude as that

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in the cold plasma. In the latter case we obtain (see (19) and (21)) the well-known expression for $\langle E_z \rangle$ (e.g. [44]).

Below we consider (30) for electron and ion cyclotron waves.

4.1. Electron cyclotron waves

The dispersion relation (16) for these waves takes the form ($\sigma = -1,\ \omega' \gg \omega_{\rm ci}$)

$$n_z'^2 = 1 + \frac{\omega_{\rm pe}^2}{\omega'(|\omega_{\rm ce}| - \omega')}.$$

In this case the expressions (30) are

$$\langle C \rangle_k = \frac{k_y c}{2\omega B_0} \frac{\omega_{\rm pe}^2}{(|\omega_{\rm ce}| - \omega')^2} \left(1 + \lambda_T \frac{|\omega_{\rm ce}| - \omega'}{\omega'}\right) \langle \mathbf{E}_1^2 \rangle_k,$$

$$\langle C_z \rangle_k = -\frac{i}{2k_z} \frac{e}{T_{\|\rm e}} \frac{\omega_{\rm pe}^2}{\omega'(|\omega_{\rm ce}| - \omega')} \langle \mathbf{E}_1^2 \rangle_k.$$

We see that when $|\omega_{\rm ce}| \gg \omega'$ and/or $\lambda_T \gg 1$ in the case of the anisotropic temperature, the magnitude of $\langle C \rangle_k$ for the warm plasma is considerably larger than that for the cold plasma.

As an example, let us consider the acceleration of magnetized dust grains in the dynamic regime (Sec. 3.1.2) in the case of the warm ions and $|\omega_{\rm ce}| \gg \omega'$. Using (22)–(25), the expressions for ε_1, c_1 (for $k_y \sim k_z$), ε_3 and the values $\langle C \rangle_k$ and $\langle C_z \rangle_k$ from this section, we find

$$\frac{\partial \langle v_{\rm d2}{}_y \rangle}{\partial t} = -\frac{\lambda_T}{8\pi n_{\rm d0} m_{\rm d}} \frac{\omega_{\rm pe}^2}{\omega' |\omega_{\rm ce}|} \frac{\partial \langle \mathbf{E}_1^2 \rangle}{\partial y},$$

$$\frac{\partial \langle v_{\rm d2}{}_z \rangle}{\partial t} = \frac{q_{\rm d}}{2m_{\rm d}} \frac{c}{B_0 \omega'} \frac{n_{\rm e0} T_{\|\rm i}}{(n_{\rm i0} T_{\|\rm e} + n_{\rm e0} T_{\|\rm i})} \frac{\partial \langle \mathbf{E}_1^2 \rangle}{\partial z}.$$

4.2. Ion cyclotron waves

In this case $\omega' \ll |\omega_{\rm ce}|$. The dispersion relation has the form

$$n_z'^2 = 1 + \frac{\omega_{\rm pi}^2}{\omega_{\rm ci}(\omega_{\rm ci} - \sigma\omega')} \left(1 - \delta \frac{\omega_{\rm ci} - \sigma\omega'}{\sigma\omega'}\right),$$

where $\sigma = \pm 1$. For $\langle C \rangle_k$ and $\langle C_z \rangle_k$ we obtain the following expressions:

$$\langle C \rangle_k = \frac{k_y c}{2\omega B_0} \frac{\omega_{\rm pi}^2}{(\omega_{\rm ci} - \sigma\omega')^2} \left[1 + (1+\delta)\lambda_T \frac{\omega_{\rm ci} - \sigma\omega'}{\omega_{\rm ci}}\right] \langle \mathbf{E}_1^2 \rangle_k,$$

$$\langle C_z \rangle_k = \frac{i}{2k_z} \frac{e}{T_{\|\rm i}} \frac{\omega_{\rm pi}^2}{\sigma\omega'(\omega_{\rm ci} - \sigma\omega')} \left[1 + (1+\delta) \frac{T_{\|\rm i}}{T_{\|\rm e}} \frac{\omega_{\rm ci} - \sigma\omega'}{\omega_{\rm ci}}\right] \langle \mathbf{E}_1^2 \rangle_k.$$

In the given case the value $\langle C \rangle_k$ in the warm and cold plasmas has the same order of magnitude, if $(1+\delta)\lambda_T \sim 1$. For anisotropic temperatures and/or large δ, $(1+\delta)\lambda_T \gg 1$, the value $\langle C \rangle_k$ in the warm plasma is larger than that in the cold one.

Using the expressions obtained in this section, we can easily (as above) find the velocities and accelerations of dust grains in the various cases considered in the present paper, which can take place in real situations.

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5. Conclusion

In the present paper a general theory for the generation of nonlinear low-frequency electromagnetic fields induced by high-frequency radiation in warm multicomponent dusty magnetoplasmas has been developed. The nonlinear low-frequency field arises when the electromagnetic radiation is non-uniform and/or non-stationary. As a concrete example of high-frequency waves electromagnetic cyclotron waves travelling along the background magnetic field have been considered. The equations for the dust particle velocities arising due to the action of the nonlinear electric field have been obtained both for magnetized and unmagnetized dust grains. It has been shown that in some cases the induced electric fields across the background magnetic field in the warm plasma can be larger than those in the cold plasma. The longitudinal nonlinear electric field in both plasmas remains of the same order of magnitude.

We have considered various regimes of nonlinear electric field generation: stationary, dynamic and quasi-dynamic. These regimes depend on the character of spatio-temporal variations of the wave amplitudes and parameters of the system. Under laboratory conditions the evolution of the applied external radiation can be controlled. In the cosmic environment the various regimes considered can, in principle, be realized. We have shown that the induced electric field is an electrostatic one in the purely stationary regime $\mathbf{k}^2 c_{\mathrm{Ai}}^2 \gg \omega\omega_{\mathrm{ci}}$ for magnetized and unmagnetized dust grains, and in the quasi-dynamic regime for unmagnetized dust grains when $\delta \sim 1$. In all other cases the induced field has an electromagnetic nature. In the stationary regime the transverse velocity of magnetized and unmagnetized dust particles can be larger, smaller or of the same order as the longitudinal velocity. In the dynamic regime of magnetized dust grains and in the dynamic and quasi-dynamic regimes of unmagnetized dust grains their transverse acceleration can be considerably larger than the longitudinal one (see (24), (25), (27)–(29)). The quasi-dynamic regime of unmagnetized dust grains ($\delta\omega_{\mathrm{ci}} \gg \omega \gg \omega_{\mathrm{cd}}$) is only possible when $n_{\mathrm{d0}} m_{\mathrm{d}} \gg n_{\mathrm{i0}} m_{\mathrm{i}}$. The relations between the transverse and longitudinal velocities and accelerations of dust grains depend on the concrete regimes, high frequencies and plasma parameters. Knowledge of the dusty plasma and wave parameters in specific situations is important for determining the dust grain dynamics. For wave packets it is also necessary to take into account their movement with the group velocity to determine the possible regimes.

The results obtained in the present paper may be useful for understanding the possible mechanisms of dust grain dynamics in astrophysical objects, in the Earth's mesosphere and in laboratory devices under the action of the unstable electromagnetic waves and/or external radiations on dusty plasmas.

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