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Acceleration of dust grains by means of the high energy ion beam

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ABSTRACT

The acceleration of charged dust grains by a high energy ion beam is investigated by obtaining the dispersion relation. The Cherenkov and cyclotron acceleration mechanisms of dust grains are compared with each other. The role of dusty plasma parameters and the magnetic field strength in the acceleration process are discussed. In addition, the stimulated waves by an ion beam in a fully magnetized dust–ion plasma are studied. It is shown that these waves are unstable at different angles with respect to the external magnetic field. It is also indicated that the growth rates increase by either increasing the ion and dust densities or decreasing the magnetic field strength. Finally, the results of our research show that the high energy ion beam can accelerate charged dust grains.

1. Introduction

The presence of massive charged dust grains in the most common state of matter in the universe has created a new branch of plasma physics. There has been a growing interest in the investigation of dusty plasmas (complex plasmas) because of the unique and wonderful properties of these plasmas [1–4]. Dust and dusty plasmas are found in space and astrophysics [5–7] (such as interstellar media, planetary rings, cometary tails, supernova remnants and the lower ionosphere of the earth), in laboratory discharges, in industrial applications [8–10] (such as microelectronic processing, semiconductor technology and carbon nanotubes), in medicine [11,12] (such as microbiology and nanomedicine) and in fusion devices [13–15] (such as tokamaks and stellarators). The dust grains can affect the plasma parameters and stability and consequently the operations of the fusion machines. Moreover, the complex plasma behavior, the radioactivity of dust grains excited by nuclear reactions and the wall–plasma interactions are the main technological problems in future controlled fusion energy production. For these reasons, the investigation of suitable methods for controlling, removing and accelerating of dust grains has attracted a lot of attention in the last decades.

There are a number of theoretical [16–21] and experimental [22–26] works that have focused on the acceleration of dust grains. Shukla et al. [16–18] have investigated the acceleration of dust grains by the ponderomotive force of dust ion-acoustic, Alfven and circularly polarized electromagnetic cyclotron waves. They have shown that the ponderomotive force of dust ion-acoustic, Alfven and electromagnetic cyclotron waves can generate space charge electric fields which can accelerate charged dust grains in dusty plasmas. Furthermore, the acceleration of charged dust grains in tokamak edges have been studied [19]. It has been indicated that the force associated with the space charge electric field at the plasma surface can accelerate of dust particles. On the other hand, it has been shown that the polarization force, arising from interactions of thermal ions with highly charged dust grains, can cause dust grain acceleration [20]. Recently, Ehsan et al. [21] have proposed a new mechanism for the acceleration of dust particles by closed filaments or vortex rings by employing the nonlinear interaction of circularly polarized electromagnetic wave with dusty plasma. The design and construction of a plasma dynamic device to accelerate dust to hypervelocities have been presented by Ticos et al. [22]. They have also investigated the dust acceleration techniques and their applications and have shown that the micron-size dust particles can be accelerated to speeds of the order of km/s over a distance of about 1 m by cold and dense plasma flows [23–25]. Diagnostics of fast dust particles in tokamak edge plasmas have been investigated by Castaldo et al. [26]. It has been shown that fast dust particles can be detected by electrostatic probes on the basis of the process of dust impact ionization.

In the present work, we investigate the interaction between the high energy ion beam and the charged dust grains. We obtain the total dielectric permittivity tensor of the system in the laboratory frame using the Lorentz transformation formulas in the same method as it was used in our previous papers [27–30]. By making use of the system dielectric permittivity and wave equation and obtaining the dispersion relation, the growth rate and the energy loss of the unstable oscillations and their dependency on the dusty plasma parameters and strength of the magnetic field are achieved.

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It is shown that dust grains injected into an ion beam can be captured and accelerated when the Cherenkov or cyclotron resonance conditions are confirmed.

This work is organized into five sections. In Section 2 the problem is formulated and the general equations for calculating the dispersion relation are given. Section 3 describes the dust grain acceleration by a high energy ion beam. The results and discussion are presented in Section 4. Finally, Section 5 contains a summary and conclusions.

2. Formulation

Consider a system of ion–dust beam plasma, like a plasma comprising a small group of ions that a straight monoenergetic beam of dust grains, that direct parallel to the external magnetic field, pumped into it. We assume that the characteristic velocities of system greatly exceed the thermal velocities of the beam and plasma particles. Consequently, we neglect the thermal motion of the particles and the energy converted to heat through the thermal motion. Hence, the energy conservation is hold in the system.

The system of ion–dust beam plasma are considered strongly magnetized, $\omega_{pd}^2 \ll \Omega_d^2$ and $\omega_{pi}^2 \ll \Omega_i^2$, where ω_{pd} and ω_{pi} are the Langmuir frequencies of dust grains and ions, respectively and $\omega_{p\alpha} = \sqrt{4\pi q_\alpha^2 N_{0\alpha}/m_\alpha}$, where $N_{0\alpha}$, m_α and q_α are the density, mass and charge of the α particle species ($\alpha = i, d$) in the laboratory frame, respectively. $\Omega_\alpha = q_\alpha B_0/m_\alpha c$ is the Larmor frequency of the α particle species, B_0 is the external magnetic field strength and c is the speed of light. The ion and dust beam, in their intrinsic frame, have the Maxwellian distribution with nonrelativistic temperature. The quasi-neutrality condition that we use reads as $N_{0i} \simeq Z_d N_{0d}$, where N_{0i} and N_{0d} are the equilibrium ion and dust number densities, respectively, and Z_d is the charge number of a dust particle. Therefore, N_{0d} is less than N_{0i} because Z_d is larger than 1. Thus, we assume that the ion beam density is larger than the dust beam density and ignore induced charge and current densities. Moreover, we assume that the all dust grains have equal masses and charges and the rotation of particles across the magnetic field is ignored, i.e., $p_{0\perp} = 0$, therefore we can use the Lorentz transform of the dielectric permittivity tensor with our solving the kinetic equations [27–31]. On the other hand, the values of drag force strongly depend on plasma parameters as well as dust size, shape, temperature, and electric charge. In this work, we ignore the effect of drag forces and the variation of the forces synchronize with the motion of dust.

The dielectric permittivity tensor of the system of ion–dust beam plasma in the moving frame can be written as follows:

$$\tilde{\varepsilon}' = \begin{pmatrix} \varepsilon'_\perp & ig' & 0 \\ -ig' & \varepsilon'_\perp & 0 \\ 0 & 0 & \varepsilon'_\parallel \end{pmatrix}, \tag{1}$$

with

$$\varepsilon'_\perp = \varepsilon'_{xx} = \varepsilon'_{yy} = 1 - \sum_\alpha \frac{\omega'^2_{p\alpha}}{\omega'^2_\alpha - \Omega^2_\alpha},$$

$$\varepsilon'_\parallel = \varepsilon'_{zz} = 1 - \sum_\alpha \frac{\omega'^2_{p\alpha}}{\omega'^2_\alpha},$$

$$\varepsilon'_{xy} = -\varepsilon'_{yx} = ig' = -i\sum_\alpha \frac{\omega'^2_{p\alpha}\Omega_\alpha}{\omega'_\alpha(\omega'^2_\alpha - \Omega^2_\alpha)},$$

$$\varepsilon'_{xz} = \varepsilon'_{zx} = \varepsilon'_{yz} = \varepsilon'_{zy} = 0, \tag{2}$$

where $\omega'_{p\alpha}$ is the plasma frequency of the particles of the type α in their respective moving frame and ω'_α is the Lorentz transformed frequency ω. Moreover, in this moving frame, we have

$$j'_{i\alpha}(\omega'_\alpha, \vec{k}'_\alpha) = \sigma^\alpha_{ij}(\omega'_\alpha, \vec{k}'_\alpha)E'_{j\alpha}(\omega'_\alpha, \vec{k}'_\alpha), \tag{3}$$

where $\sigma^\alpha_{ij}(\omega'_\alpha, \vec{k}'_\alpha)$ is the conductivity tensor of the α particle in its respective moving frame, $j'_{i\alpha}(\omega'_\alpha, \vec{k}'_\alpha)$ and $E'_{j\alpha}(\omega'_\alpha, \vec{k}'_\alpha)$ are the current density and the electric field components in this frame that are related to $j_{i\alpha}(\omega_\alpha, \vec{k}_\alpha)$ and $E_{j\alpha}(\omega_\alpha, \vec{k}_\alpha)$ in the laboratory frame by the Lorentz transformation, respectively. In the laboratory frame, the total induced current in the plasma is the sum of the currents of its charged particle component as

$$j_i(\omega, \vec{k}) = \sum_\alpha j_{\alpha i}(\omega, \vec{k}) = \sum_\alpha \sigma^\alpha_{ij}(\omega, \vec{k})E_j(\omega, \vec{k}). \tag{4}$$

In the above equations, $\vec{k}'$ is the Lorentz transformed wave vector $\vec{k}$. The physical parameters in these two frames are related together by the Lorentz transformation as follows:

$$\omega'_\alpha = (\omega - \vec{k}\cdot\vec{u}_\alpha)\gamma_\alpha,$$

$$\vec{k}'_\alpha = \vec{k} + \vec{u}_\alpha\gamma_\alpha\left[\frac{\vec{k}\cdot\vec{u}_\alpha}{u^2_\alpha}\left(1 - \frac{1}{\gamma_\alpha}\right) - \frac{\omega}{c^2}\right],$$

$$\gamma_\alpha = \left(1 - \frac{u^2_\alpha}{c^2}\right)^{-1/2},$$

$$j_{\alpha i} = \alpha_{ij}(\vec{u}_\alpha)\vec{j}'_{\alpha j},$$

$$\alpha_{ij}(\vec{u}_\alpha) = \delta_{ij} + \gamma_\alpha\left[\frac{u_{\alpha i}u_{\alpha j}}{u^2_\alpha}\left(1 - \frac{1}{\gamma_\alpha}\right) + \frac{k'_{\alpha j}u_{\alpha i}}{\omega'_\alpha}\right],$$

$$E'_{\alpha i} = \beta_{ij}(\vec{u}_\alpha)E_j$$

$$\beta_{ij}(\vec{u}_\alpha) = \frac{\omega'_\alpha}{\omega}\alpha_{ji}(\vec{u}_\alpha) = \frac{\omega'_\alpha}{\omega}\delta_{ij} + \gamma_\alpha\left[\frac{u_{\alpha i}u_{\alpha j}}{u^2_\alpha}\left(\frac{1}{\gamma_\alpha} - 1\right) + \frac{k_i u_{\alpha j}}{\omega}\right], \tag{5}$$

where $\vec{u}_\alpha$ is the averaged drift velocity of the α species. Using the aforementioned equations, the dielectric tensor of the multicomponent plasma in the laboratory frame can be written as [27–31]

$$\varepsilon_{ij}(\omega, \vec{k}) = \delta_{ij} + \frac{4\pi i}{\omega}\sum_\alpha \alpha_{i\mu}(u_\alpha)\sigma^\alpha_{\mu\nu}(\omega'_\alpha, \vec{k}'_\alpha)\beta_{\nu j}(u_\alpha) = \delta_{ij} + \sum_\alpha \frac{\omega'_\alpha}{\omega}\alpha_{i\mu}(u_\alpha)\left[\varepsilon^\alpha_{\mu\nu}(\omega'_\alpha, \vec{k}'_\alpha) - \delta_{\mu\nu}\right]\beta_{\nu j}(u_\alpha) \tag{6}$$

where $\varepsilon^\alpha_{\mu\nu}(\omega'_\alpha, \vec{k}'_\alpha)$ describes the dielectric permittivity of the α particles in their respective moving frame. Therefore, the components of dielectric permittivity tensor in the laboratory frame will be obtained from Eqs. (2), (5) and (6) as follows:

$$\varepsilon_{xx} = \varepsilon_{yy} = 1 - \sum_\alpha \frac{\omega^2_{p\alpha}\omega'^2_\alpha\gamma^{-1}_\alpha}{\omega^2(\omega'^2_\alpha - \Omega^2_\alpha)},$$

$$\varepsilon_{xy} = -\varepsilon_{yx} = -i\sum_\alpha \frac{\omega^2_{p\alpha}\omega'_\alpha\Omega_\alpha\gamma^{-1}_\alpha}{\omega^2(\omega'^2_\alpha - \Omega^2_\alpha)},$$

$$\varepsilon_{xz} = \varepsilon_{zx} = -\sum_\alpha \frac{\omega^2_{p\alpha}\omega'_\alpha k_\perp u_\alpha}{\omega^2(\omega'^2_\alpha - \Omega^2_\alpha)},$$

$$\varepsilon_{yz} = -\varepsilon_{zy} = i\sum_\alpha \frac{\omega^2_{p\alpha}\Omega_\alpha k_\perp u_\alpha}{\omega^2(\omega'^2_\alpha - \Omega^2_\alpha)},$$

$$\varepsilon_{zz} = 1 - \sum_\alpha\left(\frac{\omega^2_{p\alpha}\gamma^{-1}_\alpha}{\omega'^2_\alpha} + \frac{\omega^2_{p\alpha}\gamma_\alpha k^2_\perp u^2_\alpha}{\omega^2(\omega'^2_\alpha - \Omega^2_\alpha)}\right). \tag{7}$$