Radiation-condensation instability in a self-gravitating dusty astrophysical plasma

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Properties of radiation-condensation instability in a partially-ionized self-gravitating dusty astrophysical plasmas are studied. For this purpose, new dispersion relations for coupled dusty plasma and condensation modes in both unmagnetized and magnetized plasmas are derived. The dispersion relations are numerically analyzed to investigate the interplay between self-gravitation and impurity losses, as well as to study the effects of the external magnetic field and finite plasma beta on instabilities we found.

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I. INTRODUCTION

It is widely believed (e.g., Ref. [1]) that the thermal [also referred to as the radiation-condensation (RC)] instability [2,3] may be responsible for cloudiness of a number of astrophysical objects, for the formation of different phases of molecular clouds in diffuse interstellar and intergalactic media and planetary nebulae, and for the formation of prominences (regions of high-density and low-temperature plasma) in the solar corona. The importance of the RC instability in the edge region of laboratory tokamak discharges [4,5] and in the inner comma of a comet [6] has also been recognized. The physical origin of the thermal instability is attributed to impurity radiation losses, which critically affect the compressibility of emitting gases. In a thermally unstable state, the rate of nonequilibrium impurity radiation cooling is such that any decrease in the electron temperature leads to an increase in the rate of radiation. The latter causes further cooling and radiation enhancement, so that the system becomes unstable. The resulting impurity radiation induced instability can lead to a density condensation, so that one can have a cold, high-density plasma surrounded by a hot, low-density bulk plasma. The thermal instability in a cooling and expanding medium including self-gravity [7] and conduction in the neutral fluid dynamics has been investigated by Gomez-Pelaez and Moreno-Insertis [8].

Since astrophysical cometary and tokamak edge plasmas contain neutrals, electrons, ions, and charged dust grains (composed of graphite, silicate, and metallic compounds), several authors [9–14] investigated the RC instability in such a multicomponent plasma environment, including even the external magnetic field effect [11]. However, in astrophysical environments, the self-gravitational force [7,8] also plays an important role in controlling the dynamics of charged clouds.

In this paper, we present a detailed investigation of the RC instability in a self-gravitating, partially ionized dusty plasma which is either unmagnetized or embedded in an external magnetic field. For this purpose, we present appropriate plasma particle number density perturbations in the presence of electrostatic and electromagnetic disturbances and impurity radiation losses. Invoking the quasineutrality condition, we then derive dispersion relations, which are analyzed both analytically and numerically. The results are applied to understand the filamentation of dense molecular clouds.

II. THEORY

We consider an astrophysical fluid consisting of electrons, ions, charged dust grains, and neutrals. For simplicity, we ignore the size, charge, and mass distributions of the dust particles. The latter in astrophysical environments can be charged due to several competing processes [15,16], viz., collection of electrons and ions from the ambient plasma, photoemission, secondary emission owing to electron and ion impact and electron field emission, etc. At equilibrium, we have $en_{i0}=en_{e0}-Qn_{d0}$, where n_{j0} is the unperturbed number density of the particle species j (j equals e for electrons, i for ions, and d for dust grains) and Q is the dust charge state, viz., $Q=-eZ_d$ for negatively charged dust and $Q=eZ_d$ for positively charged dust. Here, e is the magnitude of the electron charge and Z_d is the number of charges (electrons or ions) residing on the dust grain surface.

A. Unmagnetized dusty plasmas

We first consider the RC instability in an unmagnetized self-gravitating plasma. We focus on the dynamics of low-frequency dust acoustic (DA) [17] and dust ion acoustic (DIA) waves [18] in the presence of impurity radiation losses. For $|\omega(\omega+i\nu_e)|\ll k^2V_{Te}^2$, where ω is the wave frequency, ν_e is the electron-neutral collision frequency, k is the wave number and V_{Te} is the electron thermal speed, the electron density perturbation n_{e1} is obtained by combining the continuity and momentum equations and Fourier analyzing the resultant equation. We have

$$n_{e1}\approx n_{e0}\left(\frac{e\phi}{T_{e0}}-\frac{T_{e1}}{T_{e0}}\right), \tag{1}$$

where ϕ is the wave potential and $T_{e0}(T_{e1})$ is the unperturbed (perturbed) electron temperature. The perturbed electron temperature is determined from the energy equation [19]

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$$\left(\frac{3}{2}\frac{\partial}{\partial t}-\frac{\chi_e}{n_{e0}}\nabla^2\right)T_{e1}-\frac{T_{e0}}{n_{e0}}\frac{\partial n_{e1}}{\partial t}=-\frac{\mathcal{L}(n_e,T_e)}{n_{e0}}, \quad (2)$$

where $\chi_e=3.2n_{e0}V_{Te}^2/\nu_e$ is the electron thermal conductivity and $\mathcal{L}(n_e,T_e)$ accounts for the heating H and cooling C, so that $\mathcal{L}=H-C$. The heating could be due to some external sources (e.g., hard emission, Ohmic heating, etc.) while cooling comes from impurity radiation losses. In an unperturbed state, radiative cooling is balanced by the external energy input. Taylor expanding $\mathcal{L}$ around the equilibrium state, we have

$$\mathcal{L}(n_e,T_e)=\left(\frac{\partial\mathcal{L}}{\partial n_e}\right)_{n_{e0}}n_{e1}+\left(\frac{\partial\mathcal{L}}{\partial T_e}\right)_{T_{e0}}T_{e1}, \quad (3)$$

so that Eq. (2) can be rewritten as

$$\left(\frac{3}{2}\frac{\partial}{\partial t}-D_e\nabla^2+\Omega_T\right)T_{e1}-\frac{T_{e0}}{n_{e0}}\left(\frac{\partial}{\partial t}-\Omega_n\right)n_{e1}=0, \quad (4)$$

where $D_e=3.2V_{Te}^2/\nu_e$ is the electron thermal diffusivity, $\Omega_T=n_{e0}^{-1}(\partial\mathcal{L}/\partial T_e)_{T_{e0}}$ and $\Omega_n=T_{e0}^{-1}(\partial\mathcal{L}/\partial n_e)_{n_{e0}}$.

Eliminating T_{e1} from Eq. (4) by using Eq. (1), we have

$$\left(\frac{5}{2}\frac{\partial}{\partial t}-D_e\nabla^2+\Omega_T-\Omega_n\right)\frac{n_{e1}}{n_{e0}}=\left(\frac{3}{2}\frac{\partial}{\partial t}-D_e\nabla^2+\Omega_T\right)\frac{e\phi}{T_{e0}}. \quad (5)$$

Assuming that ϕ and n_{e1} are proportional to $\exp(i\mathbf{k}\cdot\mathbf{r}-i\omega t)$, we Fourier transform Eq. (5) to obtain

$$\frac{n_{e1}}{n_{e0}}=\frac{\left(\Omega_\chi+\Omega_T-\frac{3}{2}i\omega\right)}{\left(\Omega_\chi+\Omega_T-\Omega_n-\frac{5}{2}i\omega\right)}\frac{e\phi}{T_{e0}}, \quad (6)$$

which shows that a Boltzmann electron response does not hold when the electron temperature perturbation is taken into consideration. We have denoted $\Omega_\chi=D_ek^2$.

The ion density perturbation n_{i1} can be obtained from the ion susceptibility by ignoring the ion temperature perturbation, since in most dusty plasmas one typically has $T_{i0}\ll T_{e0}$, where T_{i0} is the ion temperature. We have [16]

$$\frac{n_{i1}}{n_{i0}}\approx-\left[\frac{1+[(\omega+i\nu_i)/(\sqrt{2}kV_{Ti})]Z(\zeta_i)}{1+(i\nu_i/\sqrt{2}kV_{Ti})Z(\zeta_i)}\right]\frac{e\phi}{T_{i0}}, \quad (7)$$

where $\zeta_i=(\omega+i\nu_i)/\sqrt{2}kV_{Ti}$, Z is the plasma dispersion function, ν_i is the ion-neutral collision frequency, and V_{Ti} is the ion thermal speed. The dust number density perturbation can be derived by combining the dust continuity, the dust momentum, and Poisson's equation for a self-gravitating dusty plasma. We obtain [20]

$$\frac{n_{d1}}{n_{d0}}\approx\frac{Qk^2\phi}{m_d[\omega(\omega+i\nu_d)+\omega_J^2-3k^2V_{Td}^2]}, \quad (8)$$

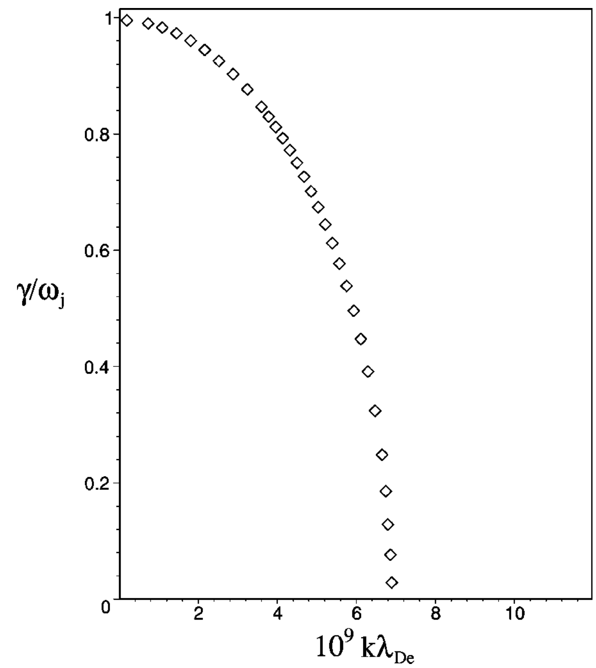

FIG. 1. The normalized growth rate γ/ω_J of the Jeans instability versus $k\lambda_{De}$, where λ_{De} is the electron Debye radius.

where m_d is the dust mass, ν_d is the dust-neutral collision frequency, V_{Td} is the dust thermal speed, $\omega_J=(4\pi G\rho_d)^{1/2}$ is the Jeans frequency [7], G is the gravitational constant, and $\rho_d=n_{d0}m_d$ is the dust mass density.

For the DA wave, the ion density perturbation is obtained from Eq. (7) in the limit $|\zeta_i|\ll 1$, yielding

$$\frac{n_{i1}}{n_{i0}}\approx-\frac{e\phi}{T_{i0}}. \quad (9)$$

Inserting Eqs. (6), (8), and (9) into the quasineutrality condition

$$en_{i1}=en_{e1}-Qn_{d1}, \quad (10)$$

we obtain the dispersion relation

$$1+\frac{\sigma(\Omega_\chi+\Omega_T-3i\omega/2)}{\Omega_\chi+\Omega_T-\Omega_n-5i\omega/2}-\frac{k^2C_D^2}{[\omega(\omega+i\nu_d)+\omega_J^2-3k^2V_{Td}^2]}=0, \quad (11)$$

where $\sigma=n_{e0}T_{i0}/n_{i0}T_{e0}$ and $C_D=Z_d(n_{d0}T_{i0}/n_{i0}m_d)^{1/2}$. Equation (11) represents the DA wave in a self-gravitating dusty plasma including impurity radiative losses.

Several comments are in order. First, in the absence of the electron-temperature perturbation, Eq. (11) yields the dispersion relation [20]

$$\omega(\omega+i\nu_d)+\omega_J^2-k^2\left(3V_{Td}^2+\frac{C_D^2}{1+\sigma}\right)=0, \quad (12)$$

which predicts the Jeans instability of the DA wave both in collisionless [21,22] and collisional [20] limits. In Fig. 1 we present the solution of Eq. (11) by using the plasma parameters that are typical for the photoassociation regions separating H II regions from dense molecular clouds (Ref. [23] and many references therein): $T_{e0}=30$ K, $T_{i0}=10$ K, $n_n=10^3$ cm^{-3}, $n_{i0}=2\times10^{-3}$ cm^{-3}. We took negatively charged grains with $T_d=1$ K, $n_d=5\times10^{-7}$ cm^{-3}, $Z_d=2000$, $m_d=10^{-11}$ gr and we assumed that $\Omega_N\simeq-\Omega_T\simeq\omega_{pd}$, where $\omega_{pd}=(4\pi n_{d0}Z_d^2e^2/m_d)^{1/2}$ is the dust plasma frequency. Second, in the presence of stationary dust grains, viz., $m_d\rightarrow\infty$, we obtain from Eq. (11) a purely growing mode ($\omega=i\gamma$), where the growth rate is

$$\gamma=\frac{2}{5+3\sigma}[\Omega_n-(1+\sigma)(\Omega_T+\Omega_\chi)]. \tag{13}$$

This is the usual radiation condensation mode provided that the isobaric instability criterion, $\Omega_n>\Omega_T+\Omega_\chi$ holds. For $\Omega_T<0$, the latter becomes $\Omega_n+|\Omega_T|>\Omega_\chi$. The growth rate of the RC unstable mode described by Eq. (13) satisfies the inequality $|\omega(\omega+i\nu_i)|\ll k^2V_{Ti}^2$ only for a narrow spectrum of modes with relatively negligible growth rate and large k.

For the DIA waves, we consider the limit $|\zeta_i|\gg1$, which permits us to express Eq. (7) in the form

$$n_{i1}\approx\frac{k^2}{4\pi e}\frac{\omega_{pi}^2}{\omega(\omega+i\nu_i)}\phi, \tag{14}$$

where $\omega_{pi}=(4\pi n_{i0}e^2/m_i)^{1/2}$ is the ion plasma frequency. Inserting now Eqs. (6), (8), and (14) into Eq. (10) we obtain the dispersion relation

$$\frac{k^2v_{Ti}^2}{\omega(\omega+i\nu_i)}-\frac{\sigma(\Omega_\chi+\Omega_T-3i\omega/2)}{\Omega_\chi+\Omega_T-\Omega_n-5i\omega/2}+\frac{k^2C_D^2}{[\omega(\omega+i\nu_d)+\omega_J^2]}=0. \tag{15}$$

One can see that the third term, which corresponds to the response of the dust, is small in comparison with the first term. We have numerically analyzed Eq. (15) using the same plasma parameters we used in the preceding subsection. Figure 2 depicts the normalized growth rate of the RC mode as a function of the normalized wavelength. It is seen that longer wavelength modes grow faster.

B. Magnetized dusty plasmas

We now consider the RC instability in a dusty plasma embedded in an external magnetic field $\hat{\mathbf{z}}B_0$, where $\hat{\mathbf{z}}$ is the unit vector along the z axis and B_0 is the strength of the magnetic field. Here, Eq. (6) holds as long as $|\omega|\ll\nu_e\ll\omega_{ce}$, except that $D_e\nabla^2$ is replaced by $D_\parallel\nabla_\parallel^2+D_\perp\nabla_\perp^2$, where ω_{ce} is the electron gyrofrequency, $D_\parallel\equiv D_e$, $D_\perp=4.7\rho_e^2\nu_e$, $\rho_e=V_{Te}/\omega_{ce}$, and the subscripts $\parallel$ and $\perp$ denote the components parallel and perpendicular to $\hat{\mathbf{z}}$. The dust number density perturbation, given by Eq. (8), remains valid in a magnetoplasma if the wave frequency is much larger than the dust gyrofrequency. However, the external magnetic field can greatly affect the ion dynamics. The ion density perturbation is given by [16]

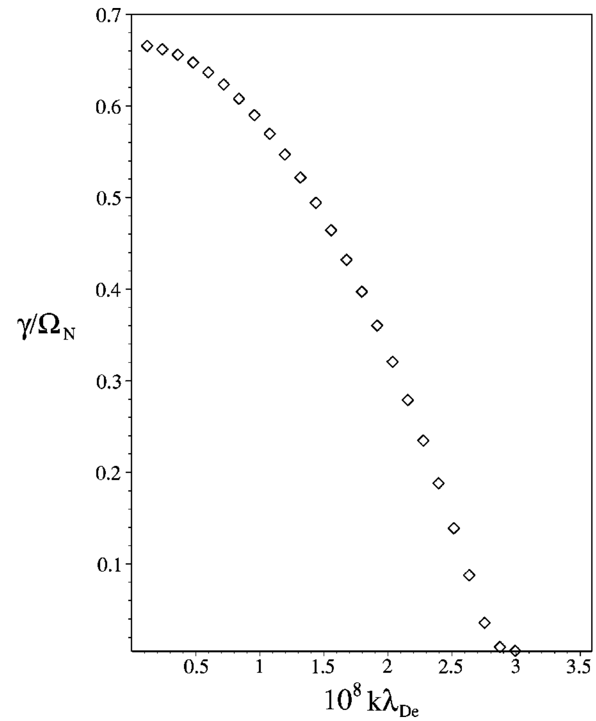

FIG. 2. The normalized growth rate γ/Ω_N of the RC instability is plotted against $k\lambda_{De}$.

$$\frac{n_{i1}}{n_{i0}}\approx-\left[\frac{1+\sum_{n=-\infty}^{\infty}\Gamma_n(b_i)[(\omega+i\nu_i)/(\sqrt{2}k_zV_{Ti})]Z(\xi_i)}{1+\sum_{n=-\infty}^{\infty}\Gamma_n(b_i)(i\nu_i/\sqrt{2}k_zV_{Ti})Z(\xi_i)}\right]\times\frac{e\phi}{T_{i0}}, \tag{16}$$

where $\xi_i=(\omega+i\nu_i-n\omega_{ci})/\sqrt{2}k_zV_{Ti}$, $\Gamma_n=I_n\exp(-b_i)$, I_n is the n-order modified Bessel function, $b_i=k_\perp^2\rho_i^2$, $\rho_i=V_{Ti}/\omega_{ci}$ is the ion gyroradius, $\omega_{ci}=eB_0/m_ic$ is the ion gyrofrequency, m_i is the ion mass, and c is the speed of light in vacuum. We can replace Eq. (16) in the limits $\xi_i\gg1$ and $b_i\ll1$ by [24]

$$n_{i1}\approx\frac{k_\perp^2}{4\pi e}\frac{\omega_{pi}^2(\omega+i\nu_i)}{\omega[(\omega+i\nu_i)^2-\omega_{ci}^2]}\phi+\frac{k_z^2}{4\pi e}\frac{\omega_{pi}^2}{\omega(\omega+i\nu_i)}\phi. \tag{17}$$

Inserting Eqs. (6), (8), and (17) into Eq. (10), we obtain the desired dispersion relation. Focusing on waves with $\omega\ll\omega_{ci}$ we have

$$\frac{\omega+i\nu_i}{\omega}-\frac{k_z^2}{k_\perp^2}\frac{\omega_{ci}^2}{\omega(\omega+i\nu_i)}+\frac{(\Omega_b+\Omega_T-3i\omega/2)}{b_s(\Omega_b+\Omega_T-\Omega_n-5i\omega/2)}-\frac{\omega_{\mathrm{DLH}}^2}{[\omega(\omega+i\nu_d)+\omega_J^2-3k^2V_{Td}^2]}=0, \tag{18}$$

where $\Omega_b=(3.2k_\parallel^2\lambda_m^2+4.7k_\perp^2\rho_e^2)\nu_e$, $\lambda_m=V_{Te}/\nu_e$, $b_s=k_\perp^2\rho_s^2$, $\rho_s=C_s/\omega_{ci}$, $C_s=(T_e n_{i0}/m_i n_{e0})^{1/2}$ and $\omega_{\rm DLH}=kC_D/\sqrt{b_i}$ is the dust-lower-hybrid (DLH) frequency [25].

For the two-dimensional ion motion in a plane perpendicular to $\hat{\mathbf{z}}$, in the presence of static dust grains but in absence of ion-neutral collisions, we obtain from Eq. (18) a purely growing mode $(\omega=i\gamma_b)$, where the growth rate is

$$\gamma_b=\frac{2b_s}{3+5b_s}\left[\Omega_n-\frac{1+b_s}{b_s}(\Omega_T+\Omega_b)\right], \tag{19}$$

if $\Omega_n>(1+b_s)(\Omega_T+\Omega_b)/b_s$.

For the two-dimensional magnetized dust, we have [26]

$$\frac{n_{d1}}{n_{d0}}=\frac{Q(\omega+i\nu_d)k_\perp^2\phi}{m_d\{\omega[(\omega+i\nu_d)^2-\omega_{cd}^2]+(\omega+i\nu_d)(\omega_J^2-3k_\perp^2V_{Td}^2)\}}, \tag{20}$$

so that for this case, instead of Eq. (18), we have

$$\frac{\omega+i\nu_i}{\omega}-\frac{k_z^2}{k_\perp^2}\frac{\omega_{ci}^2}{\omega(\omega+i\nu_i)}+\frac{(\Omega_b+\Omega_T-3i\omega/2)}{b_s(\Omega_b+\Omega_T-\Omega_n-5i\omega/2)}-\frac{C_{de}^2\omega_{ci}^2(\omega+i\nu_d)}{C_s^2\{\omega[(\omega+i\nu_d)^2-\omega_{cd}^2]+(\omega+i\nu_d)(\omega_J^2-3k^2V_{Td}^2)\}}=0, \tag{21}$$

where $C_{de}=Z_d(T_e n_{d0}/m_d n_{e0})^{1/2}$.

We have numerically solved Eq. (21) for $B_0=10$ μG, choosing the same astrophysical plasma parameters as we used in Secs. II A and II B. Figure 3 displays the normalized growth rate of the unstable root, which corresponds to the RC mode, as a function of the normalized wave number $(k\lambda_{De})$ for different propagation angles $\theta=\tan^{-1}(k_\perp/k_\parallel)$. It is seen that the growth rate is larger for waves that are propagating at large angles to the magnetic field direction. This occurs due to the reduction of electron thermal diffusivity at these angles.

C. Electromagnetic effects

Let us finally study the RC instability involving electromagnetic perturbations whose electric and magnetic fields are denoted by $\mathbf{E}=-\nabla\phi-\hat{\mathbf{z}}c^{-1}\partial_t A_z$ and $\mathbf{B}_\perp=\nabla A_z\times\hat{\mathbf{z}}$, where A_z is the parallel (to $\hat{\mathbf{z}}$) component of the vector potential. For $\nu_e\lambda_e^2\nabla_\perp^2\ll|\partial_t|\ll\nu_e\ll\omega_{ce}$ and $|\nu_e\partial_t|\ll V_{Te}^2\partial_z^2$, we replace Eq. (1) with

$$\partial_z n_{e1}=\frac{en_{e0}}{T_e}\partial_z\left(\phi-\frac{T_{e1}}{e}\right)+\frac{en_{e0}}{cT_e}\partial_t A_z, \tag{22}$$

where $\lambda_e=c/\omega_{pe}$ is the electron skin depth and ω_{pe} is the electron plasma frequency. The compressional magnetic field perturbation has been neglected in view of the low-beta approximation.

By using the perpendicular (to $\hat{\mathbf{z}}$) component of electron fluid velocities, viz.,

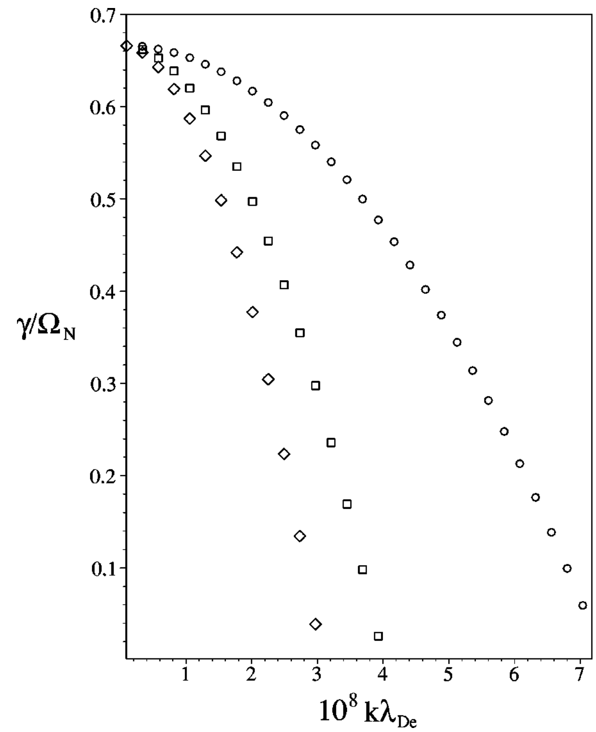

FIG. 3. The normalized growth rate γ/Ω_N of the RC instability versus $k\lambda_{De}$ for different propagation angles. (a) $\theta=\pi/8$, rhombs. (b) $\theta=\pi/4$, squares. (c) $\theta=3\pi/8$, circles.

$$\mathbf{v}_{e\perp}\approx\frac{c}{B_0}\hat{\mathbf{z}}\times\nabla\left(\phi-\frac{T_{e0}}{e}\frac{n_{e1}}{n_{e0}}\right), \tag{23}$$

and Ampère's law

$$v_{ez}\approx\frac{c}{4\pi e n_{e0}}\nabla_\perp^2 A_z, \tag{24}$$

we obtain from the electron continuity equation

$$\partial_t n_{e1}+\frac{c}{4\pi e}\partial_z\nabla_\perp^2 A_z=0. \tag{25}$$

Eliminating A_z from Eq. (22) by using Eq. (25) we have

$$(\partial_t^2+c^2\lambda_{De}^2\partial_z^2\nabla_\perp^2)n_{e1}=\frac{c^2}{4\pi e}\partial_z^2\nabla_\perp^2\left(\phi-\frac{T_{e1}}{e}\right), \tag{26}$$

where $\lambda_{De}=(T_{e0}/4\pi n_{e0}e^2)^{1/2}$ is the electron Debye radius. On the other hand, eliminating T_{e1} from Eq. (26) by using Eq. (4) with an appropriate replacement for the term $D_e\nabla^2$ in an external magnetic field, as mentioned in Sec. II B, we obtain

$$(\partial_t^2+c^2\lambda_{De}^2\partial_z^2\nabla_\perp^2)\left(\frac{3}{2}\partial_t-D_\parallel^2\partial_z^2-D_\perp^2\nabla_\perp^2+\Omega_T\right)n_{e1}+c^2\lambda_{De}^2\partial_z^2\nabla_\perp^2(\partial_t-\Omega_n)n_{e1}=\frac{en_{e0}c^2\lambda_{De}^2}{T_{e0}}\partial_z^2\nabla_\perp^2\left(\frac{3}{2}\partial_t-D_\parallel^2\partial_z^2-D_\perp\nabla_\perp^2+\Omega_T\right)\phi. \tag{27}$$

Fourier transforming Eq. (27) we have

$$\frac{n_{e1}}{n_{e0}}=\frac{\Omega_e^2\left(\Omega_b+\Omega_T-\frac{3}{2}i\omega\right)}{(\Omega_e^2-\omega^2)\left(\Omega_b+\Omega_T-\frac{3}{2}i\omega\right)-\Omega_e^2(\Omega_n+i\omega)}\frac{e\phi}{T_e}, \tag{28}$$

where $\Omega_e=k_z c k_\perp \lambda_{De}$. For $\Omega_e \gg \omega$, Eq. (28) reduces to Eq. (6).

Equations (17) and (19) for the ion and dust number density perturbations, respectively, remain unchanged. The dispersion relation for $\omega_{ci}\gg|\omega|,\nu_i$ turns out to be

$$\begin{aligned}&\frac{\omega+i\nu_i}{\omega}\frac{\omega_{pi}^2}{\omega_{ci}^2}-\frac{k_z^2}{k_\perp^2}\frac{\omega_{pi}^2}{\omega(\omega+i\nu_i)}\\&\quad+\frac{k_z^2c^2\left(\Omega_b+\Omega_T-\frac{3}{2}i\omega\right)}{\left[(\Omega_e^2-\omega^2)\left(\Omega_b+\Omega_T-\frac{3}{2}i\omega\right)-\Omega_e^2(\Omega_n+i\omega)\right]}\\&\quad-\frac{\omega_{pd}^2(\omega+i\nu_d)}{\{\omega[(\omega+i\nu_d)^2-\omega_{cd}^2]+(\omega+i\nu_d)(\omega_J^2-3k^2V_{Td}^2)\}}\\&=0.\end{aligned} \tag{29}$$

Equation (29) has been numerically analyzed and the results show that the electromagnetic effects do not alter significantly the dispersion properties of the RC unstable mode, at least for the astrophysical plasma parameters we considered in Secs. II A and II B.

III. DISCUSSION AND CONCLUSIONS

In this paper, we have investigated the radiation-condensation instability in a self-gravitating dusty plasma, which is either unmagnetized or magnetized. It is found that the combined influence of the dust inertia and impurity radiation losses can produce novel instabilities involving the DA and DLH waves. In the presence of impurity radiation losses, the electron temperature fluctuations break the Boltzmann electron response so that the electron-density perturbation cannot keep in phase with the wave potential. The resulting phase lag causes the RC mode to grow. By choosing the plasma parameters that are relevant for molecular clouds, we have studied the variation of the growth rates associated with the self-gravitating and RC instabilities. Our results show that long wavelength modes grow faster. In the presence of the external magnetic field, nearly perpendicular propagating modes have larger growth rates than those propagating at a small angle to the magnetic field direction. This is due to the angular dependence of the electron thermal diffusivity. Furthermore, for the astrophysical parameters we used, the corrections due to electromagnetic effects are found to be insignificant. In conclusion, we stress that the present RC and self-gravitational instabilities may be the cause of structuring of plasmas in dusty cosmic regions [27]. We are hoping that forthcoming space missions and Hubble Space Telescope observations shall provide more observations that may lend support to our theoretical predictions.

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