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Instabilities in strongly coupled plasmas

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Abstract

The conventional Vlasov treatment of beam-plasma instabilities is inappropriate when the plasma is strongly coupled. In the strongly coupled liquid state, the strong correlations between the dust grains fundamentally affect the conditions for instability. In the crystalline state, the inherent anisotropy couples the longitudinal and transverse polarizations, and results in unstable excitations in both polarizations. We summarize analyses of resonant and non-resonant, as well as resistive instabilities. We consider both ion-dust streaming and dust beam-plasma instabilities. Strong coupling, in general, leads to an enhancement of the growth rates. In the crystalline phase, a resonant transverse instability can be excited.

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1. Introduction

The conventional Vlasov treatment of beam-plasma instabilities is inappropriate when the plasma is strongly coupled. A realizable example is the case of strongly coupled dusty plasmas under laboratory conditions, permeated by streaming ions: in this case, beam-plasma instabilities may be excited [1, 2]. Another possible example involves the resonant interaction between a dust beam and a strongly coupled dusty plasma under laboratory or microgravity conditions. In the strongly coupled liquid state, the correlations between the dust grains fundamentally affects the dispersion of very low-frequency dust-acoustic waves [3–9] and the conditions for their instability. In the crystalline state, the inherent anisotropy couples the longitudinal and transverse polarizations and can result in unstable excitations in both polarizations.

In this paper, we briefly summarize our recent analyses [10] of resonant and non-resonant beam-plasma instabilities, as well as resistive instabilities, in a model plasma composed of a cold, strongly correlated plasma, penetrated by a cold, weakly correlated beam.

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The strongly coupled plasma particles have charge state Z, thermal energy T and average interparticle spacing d, and satisfy the condition $\Gamma = Z^2e^2/dT \gg 1$. Both the beam and the plasma particles are screened by the background lighter charged particles, with a Debye screening length $\lambda_D = k_D^{-1}$, and interact via a Yukawa (screened Coulomb) potential, $\phi(\mathbf{r}) = Z^2e^2\exp(-k_Dr)/r$. Collisions with background neutral molecules are taken into account through collision frequencies ν_j (where j denotes the charged particle species).

A dielectric function formulation is used for the analysis. The main plasma is characterized by the dielectric function $\epsilon(\mathbf{k}, \omega)$, and the beam by its polarizability $\alpha(\mathbf{k}, \omega - \mathbf{k}\cdot\mathbf{V}_0)$, where $\mathbf{V}_0$ is the velocity of the beam. The dispersion relation is given by

$$\mathrm{Det}\left[\epsilon(\mathbf{k}, \omega) - \frac{k^2c^2}{\omega^2}\left(\mathbf{I} - \frac{\mathbf{kk}}{k^2}\right) + \alpha(\mathbf{k}, \omega - \mathbf{k}\cdot\mathbf{V}_0)\right] = 0. \tag{1}$$

The unstable frequency is in the vicinity of $\mathbf{k}\cdot\mathbf{V}_0$ (the Cerenkov condition) and is given by $\omega = \mathbf{k}\cdot\mathbf{V}_0 + \delta$. When $\omega = \omega(\mathbf{k})$, where $\omega(\mathbf{k})$ is a mode of the plasma, the instability is resonant: otherwise, it is non-resonant.

2. Instability calculations in the liquid state

When the coupling constant is high, but below a certain critical value, the plasma is in a strongly correlated liquid state. The essential feature of the liquid state from the point of view of the excitation of instabilities is that, in contrast to the crystalline phase, it is isotropic. Thus the collective mode structure of the strongly coupled liquid phase of the Yukawa plasma comprises both a longitudinal mode and a transverse mode, similar to the one-component plasma (OCP) [12–15]. In our previous work [3, 4], we investigated the details of this mode structure, on the basis of the quasi-localized charge approximation (QLCA) [11, 12] (for a review, see [16]). The QLCA calculations for the longitudinal and transverse dielectric function of the plasma give [3, 4]

$$\epsilon_{L/T}(\mathbf{k}, \omega) = 1 - \frac{\Omega_{L0}^2(\mathbf{k})}{\omega^2 - D_{L/T}(\mathbf{k})}. \tag{2}$$

Here $\Omega_{L0}(\mathbf{k})$ is the longitudinal mode frequency in the weakly coupled phase of the plasma,

$$\Omega_{L0}^2(\mathbf{k}) = \Omega_0^2\frac{k^2}{k^2 + k_D^2}$$

where Ω_0 is the plasma frequency, and the local field functions $D_L(\mathbf{k})$ and $D_T(\mathbf{k})$ are functionals of the equilibrium pair correlation function (see [16]). The dispersion relations for the longitudinal (plasmon) and transverse (shear) modes are obtained from $\epsilon_L(\mathbf{k}, \omega) = 0$ and $\epsilon_T^{-1}(\mathbf{k}, \omega) = 0$, respectively. These give

$$\Omega_L(\mathbf{k}) = \sqrt{\Omega_{L0}^2(\mathbf{k}) + D_L(\mathbf{k})} \tag{3a}$$

for the longitudinal mode and

$$\Omega_T(\mathbf{k}) = \sqrt{D_T(\mathbf{k})} + \mu \tag{3b}$$

for the transverse mode which arises due to correlational effects; μ is a small electromagnetic correction.

In the following, we consider two scenarios for instabilities driven by a weakly coupled beam (with plasma frequency ω_0 and velocity $\mathbf{V}_0$) in a strongly coupled plasma (with plasma frequency Ω_0): (1) a 'weak' beam ($\omega_0 \ll \Omega_0$) excites a resonant longitudinal instability, or a non-resonant quasi-transverse instability; (2) a 'strong' beam ($\omega_0 \gg \Omega_0$) excites a Buneman type instability. An example of scenario (1) involves injection of a low density dust beam into a strongly coupled dusty plasma, while an example of scenario (2) involves the streaming of ions relative to dust in strongly coupled dusty plasmas under laboratory conditions. In the weakly coupled component, correlations between the particles are negligible, while in the strongly coupled component their role is crucial. Also, the correlations between the particles of the two components will be ignored.

2.1. Dust beam/dust plasma instabilities

We consider a strongly coupled dusty plasma, penetrated by a 'weak', weakly coupled dust beam, with $\omega_0 \ll \Omega_0$. These conditions imply that the dust density in the beam is much lower than the plasma dust density. The beam may excite both longitudinal and transverse waves in the strongly coupled plasma. The resonance conditions for the excitation of longitudinal and of transverse waves, are, respectively,

$$\omega \simeq \mathbf{k} \cdot \mathbf{V}_0 + \delta \simeq \Omega_L(\mathbf{k}) + \delta \tag{4a}$$

$$\omega \simeq \mathbf{k} \cdot \mathbf{V}_0 + \delta \simeq \Omega_T(\mathbf{k}) + \delta \tag{4b}$$

where $\delta \ll \mathbf{k} \cdot \mathbf{V}_0$.

First, we consider the excitation of longitudinal waves. The dispersion relation is given by equation (3a). For $\Gamma \gg 1$, and in the long wavelength limit, $D_L(k \to 0) \approx -l^2k^2$, where $l^2 = \Omega_0^2 d^2 f(\kappa)$, and $f(\kappa) \simeq 4(a_0 + 0.5a_2\kappa^2 + 6a_4\kappa^4)/45$, with $a_0 = 0.899$, $a_2 = 0.103$, and $a_4 = -0.003$, with $\kappa = k_D d$ [3]. Note that since $D_L(k) < 0$, the effect of strong coupling is to 'soften' the mode dispersion. From equation (1), the dispersion relation for the longitudinal wave instability, in the limit $k^2c^2/\omega^2 \to \infty$, is

$$\epsilon_L(\mathbf{k}, \omega) + \alpha_L(\mathbf{k}, \omega - \mathbf{k} \cdot \mathbf{V}_0) \approx 0 \tag{5}$$

where the longitudinal component of the beam polarizability is given by [10]

$$\alpha_L(\mathbf{k}, \omega - \mathbf{k} \cdot \mathbf{V}_0) = -\frac{\omega_{L0}^2(\mathbf{k})}{(\omega - \mathbf{k} \cdot \mathbf{V}_0)^2} \tag{6}$$

with

$$\omega_{L0}^2(\mathbf{k}) = \omega_0^2 \frac{k^2}{k^2 + k_D^2}.$$

Using the resonance condition (4a) in equation (5), and observing that $|D_L(\mathbf{k})| \ll \Omega_{L0}^2(\mathbf{k})$, the growth rate for the longitudinal instability becomes

$$\gamma = \operatorname{Im} \delta = \frac{\sqrt{3}}{2^{4/3}} \Omega_{L0}(\mathbf{k}) \left(\frac{\omega_0}{\Omega_0}\right)^{2/3} \left[1 + \frac{\kappa^2}{6} f(\kappa)\right]. \tag{7}$$

The growth rate $\gamma/\Omega_{L0}(\mathbf{k})$ is proportional to $(\omega_0/\Omega_0)^{2/3}$ as in the weakly correlated case, but strong coupling results in a small enhancement of the instability given by the term in square brackets in equation (7).

Next, we consider the case where a transverse wave is excited, leading to a quasi-transverse instability. The dispersion relation for transverse shear waves is given by equation (3b). For $\Gamma \gg 1$, and in the long wavelength limit, $D_T(k \to 0) = s^2k^2$. Here s is the shear wave velocity, given by $s^2 = \Omega_0^2 d^2 h(\kappa)$, where $h(\kappa) = b_0 + b_2\kappa^2 + b_4\kappa^4$, with $b_0 = 0.034$, $b_2 = -0.009$ and $b_4 = 0.001$ [4]. From (1), in the limit $c^2 \gg s^2$, the dispersion relation for the excitation of waves satisfying the resonance condition (4b) becomes

$$1 - \frac{\Omega_{L0}^2(\mathbf{k})}{D_T(\mathbf{k}) - D_L(\mathbf{k})} - \frac{\omega_{L0}^2(\mathbf{k})}{\delta^2} \approx 0. \tag{8}$$

The shear mode prevails in a frequency domain where $\epsilon_L(\mathbf{k},\omega) < 0$. Thus equation (8) describes a non-resonant type longitudinal instability with a quadratic equation for δ. The growth rate is

$$\gamma = \operatorname{Im}\delta = \omega_{L0}(\mathbf{k})\sqrt{\frac{\kappa^2[h(\kappa)+f(\kappa)]}{1-\kappa^2[h(\kappa)+f(\kappa)]}}. \tag{9}$$

For this case, the growth rate $\gamma/\Omega_{L0}(\mathbf{k})$ is proportional to ω_0/Ω_0, that is, the ratio of the beam to plasma frequencies. Insofar as the growth rate is concerned, the fact that a transverse wave is also excited is not relevant. The resonance of the instability with the transverse wave becomes important in two ways. First, as displayed below, the polarization of the perturbation electric field associated with the instability picks up a small transverse component: in this sense the instability becomes quasi-transverse,

$$\frac{E_T}{E_L} \approx -\frac{\omega_{L0}^2}{\delta^2}\frac{\omega}{kc}\frac{V_{0T}}{c}$$

where E_T and E_L are the transverse and longitudinal components, respectively, of the perturbed electric field, $V_{0T} = V_0 - (\mathbf{k}\cdot\mathbf{V}_0)/k$, and it has been assumed that $|\delta\omega| \gg |\mu^2|$ [10]. Thus in order for the mode to be quasi-transverse, $\mathbf{V}_0$ must have a component perpendicular to $\mathbf{k}$. The second aspect that makes the appearance of the transverse mode important is that a coupling between the transverse and longitudinal modes would complete the feedback loop and would render the instability a genuine resonant transverse instability. While such a coupling mechanism is not part of the present model, it can be due to various physical effects, such as nonlinear mode–mode coupling, finite-size beam effects, and an anisotropy of the medium. An example for the last effect is given in section 3.

2.2. Longitudinal ion-dust instability

In several laboratory dusty plasmas, the ions stream relative to the charged dust with speed V_0 larger the ion thermal speed v_i, thus comprising an ion beam. For example, in dust Coulomb crystals, the grains are localized near plasma sheath interface regions, where the ion flow toward the electrode is approximately equal to the ion sound speed, which can be $\gg v_i$ since, typically, $T_e/T_i \gg 1$ [17, 2]. Further, ions can have flow speeds $V_0 > v_i$ in dusty plasma experiments where the ions are accelerated by an electric field [18]. Ion dust streaming instabilities in dusty plasmas have been investigated previously [17–19, 1], and here we consider the effect of strong coupling on such instabilities.

We adopt a viewpoint from the frame in which the weakly correlated ions (with plasma frequency ω_0) are at rest, and the strongly coupled dust grains (with plasma frequency Ω_0) comprise a 'weak' beam ($\Omega_0 \ll \omega_0$) with velocity $\mathbf{V}_0$. From equation (1), the dispersion relation is

$$1 - \frac{\omega_{L0}^2}{\bar\omega^2} - \frac{\Omega_{L0}^2}{(\bar\omega + \mathbf{k}\cdot\mathbf{V}_0)^2 - D_L(\mathbf{k})} = 0 \tag{10}$$

where the Galilean transformation gives $\bar\omega = \omega - \mathbf{k}\cdot\mathbf{V}_0$. Using the resonance condition

$$\bar\omega \simeq -\mathbf{k}\cdot\mathbf{V}_0 + \delta \simeq \Omega_{L0}(\mathbf{k}) + \delta \tag{11}$$

in equation (10) and observing that $D_L(\mathbf{k})\big/\Omega_{L0}^2(\mathbf{k}) \ll 1$, the solution of the resulting cubic equation gives for the real part of the frequency ω_r (in the lab frame) and the growth rate γ

$$\omega_r = \frac{\Omega_{L0}(\mathbf{k})}{2}\left(\frac{\omega_0}{2\Omega_0}\right)^{1/3}\left[1 - \frac{\kappa^2 f(\kappa)}{3}\left(\frac{2\Omega_0}{\omega_0}\right)^{2/3}\right] \tag{12a}$$

$$\gamma = \frac{\sqrt{3}\Omega_{L0}(\mathbf{k})}{2}\left(\frac{\omega_0}{2\Omega_0}\right)^{1/3}\left[1 + \frac{\kappa^2 f(\kappa)}{3}\left(\frac{2\Omega_0}{\omega_0}\right)^{2/3}\right]. \tag{12b}$$

Thus strong coupling results in a small enhancement of the growth rate.

Under laboratory conditions, collisions with neutrals are important, since both ν_i/ω_0 and ν_d/Ω_0 can be >0.1 (here ν_i, ν_d are the ion-neutral, dust-neutral collision frequencies, respectively). Thus, the instability is likely to be a *dissipative instability* [17], with $\delta \ll \nu_i$. Including collisional effects in equation (10) yields the dispersion relation

$$1 - \frac{\omega_{L0}^2}{\bar{\omega}(\bar{\omega} + i\nu_i)} - \frac{\Omega_{L0}^2}{(\bar{\omega} + \mathbf{k}\cdot\mathbf{V}_0)(\bar{\omega} + \mathbf{k}\cdot\mathbf{V}_0 + i\nu_d) - D_L(\mathbf{k})} = 0. \tag{13}$$

The solution of (13), in the limit $\nu_d \ll \omega_r$, gives for the growth rate

$$\gamma = -\frac{\nu_d}{2} + \mathrm{Im}\,\delta = -\frac{\nu_d}{2} + \frac{\Omega_{L0}(\mathbf{k})}{\sqrt{2}}\left(\frac{\omega_{L0}(\mathbf{k})}{\nu_i}\right)^{1/2}\left[1 + \frac{\nu_i}{2\omega_{L0}(\mathbf{k})}\kappa^2 f(\kappa)\right]. \tag{14}$$

As can be seen from (14), the dissipative instability is also enhanced by strong coupling. Although the effect is small, it is more important than in the collisionless case.

3. Instabilities in the crystal state

The main physical difference between the liquid and crystalline (solid) state is that the liquid is isotropic and the crystal is not. For the crystalline phase we can use the same formalism as in section 2.1, but in this case the dynamical matrix $D_{\alpha\beta}(\mathbf{k})$ (calculated via a summation over lattice sites, see, e.g., [20–22]) is, in general, non-diagonal for an arbitrary direction of wave propagation in the crystal. Thus the longitudinal and transverse modes are coupled. Here we consider a scenario similar to the one in section 2.1; the plasma crystal is taken as 2D and it is pervaded by a 'weak', weakly coupled dust beam with $\omega_0 \ll \Omega_0$. When the off-diagonal elements of $D_{\alpha\beta}(\mathbf{k})$, say $C_{\alpha\beta}(\mathbf{k})$, are small, there is still a quasi-longitudinal mode with frequency $\Omega_{L*}(\mathbf{k})$, and a quasi-transverse mode with frequency $\Omega_{T*}(\mathbf{k})$ (e.g., [21, 22]). The dispersion relation for these two modes now become

$$\Omega_{L*}(\mathbf{k}) = \sqrt{\Omega_L^2(\mathbf{k}) - \Lambda(\mathbf{k})} \tag{15a}$$

$$\Omega_{T*}(\mathbf{k}) = \sqrt{\Omega_T^2(\mathbf{k}) + \Lambda(\mathbf{k})} \tag{15b}$$

where $\Lambda(\mathbf{k})$ represents the coupling due to anisotropy,

$$\Lambda(\mathbf{k}) \simeq \frac{C^2(\mathbf{k})}{\Omega_L^2(\mathbf{k}) - \Omega_T^2(\mathbf{k})} \simeq \frac{C^2(\mathbf{k})}{\Omega_{L0}^2(\mathbf{k})}.$$

The resonance conditions are now given by equations of the form (4*a*) and (4*b*) but with the replacement $\Omega_L \to \Omega_{L*}$ and $\Omega_T \to \Omega_{T*}$. Preliminary results show that (i) the longitudinal resonant instability is excited (cf equation (7)), whose growth rate is not substantially affected by the anisotropy and (ii) a new type of resonant transverse instability develops. The latter is due to the fact that the polarizability of the quasi-transverse mode now also has a longitudinal component. The frequency of the instability is in the vicinity of Ω_{T*}, with a growth rate of the order [10]

$$\gamma_T \simeq \Omega_{L0}(\mathbf{k})\left(\frac{\omega_0}{\Omega_0}\right)^{2/3} kd\,\sin^{2/3}(2\theta)q(\kappa) \tag{16}$$

Table 1. Possible experimental parameters. (The subscripts n, e, i, 1 and 2 refer to neutrals, electrons, ions, dust species 1 and dust species 2, respectively).

Plasma	Ar, $n_n \sim 3 \times 10^{14}$ cm^{-3}, $n_i \sim 5 \times 10^8$ cm^{-3}, $T_e \sim 2$ eV, $T_i \sim 0.1$ eV, $T_n \sim 0.03$ eV $\sim T_1 \sim T_2$
Dust	$r_1 = 0.3\,\mu$m, $Z_1 \sim 1000$, $\rho_1 \sim 1.5$ g cm^{-3}, $n_1 \sim 2 \times 10^4$ cm^{-3}
Dust beam	$r_2 = 0.05\,\mu$m, $Z_2 \sim 170$, $\rho_2 \sim 10$ g cm^{-3}, $n_2 \sim 6 \times 10^2$ cm^{-3}
Electric field	$E \sim 0.3$ V m$^{-1} \Rightarrow V_{02} \sim 10$ cm s^{-1}

where $q(\kappa)$ is determined by the lattice structure. Thus the transverse growth rate is smaller than the longitudinal one, primarily because of the $kd(<1)$ factor, and it is strongly angle dependent. Nevertheless, the fact that the instability prevails at a much lower frequency than the longitudinal one may render it observable without being masked by the latter.

We propose the following set of experimental parameters to investigate these instabilities. Consider a dusty plasma containing negatively charged dust with two different radii, r_1 and r_2 (with $r_1 \gg r_2$), and with different mass densities, ρ_1 and ρ_2. Suppose there is a horizontal electric field $\mathbf{E}$ in the system (e.g., an axial electric field in a Q-machine [23]). In the presence of this field (where $|\mathbf{E}|$ is much less than that required for an ion-dust streaming instability), the velocity of species 2,

$$\mathbf{V}_{02} = -\frac{eZ_2\mathbf{E}}{m_2\nu_2} \propto \frac{1}{r_2} \tag{17}$$

is much larger than that of species 1 (here Z_2, m_2 and ν_2 are the charge state, mass and collision frequency, respectively, of dust species 2). Dust species 2 would then play the role of a beam (with plasma frequency ω_{02}) penetrating a dusty plasma of species 1 (with plasma frequency Ω_{01}). The critical speed for the transverse instability would be lowest, since s is much less than the phase speed of the dust acoustic wave. Table 1 lists some nominal parameters that satisfy the following conditions: the dusty plasma is strongly coupled $\left(Z_1^2e^2n_1^{1/3}/T_d \gg 1\right)$; the beam speed is larger than the dust acoustic phase speed ($V_{02} > \Omega_{01}\lambda_D$); the beam is weak ($\omega_{02}/\Omega_{01} < 1$); collisional effects are small ($\nu_2/\omega_{02} < 1$, $\nu_1/\Omega_{01} < 1$); 'marginally' cold beam $\left(T_2/m_2V_{02}^2 \sim [\omega_{02}/\Omega_{01}]^{2/3}\right)$.

4. Summary

Strong coupling, in general, enhances the longitudinal beam-plasma instability. The physical reason for this, both in the case of of resonant and Buneman-type instability, may be sought in the reduction of the longitudinal wave frequency from its Vlasov value [3, 4], and the concomitant reduction of $\partial\epsilon_L(\mathbf{k}, \omega)/\partial\omega)$ at $\Omega_L(\mathbf{k})$. For the quasi-transverse instability, the main physical effect is the appearance of the low frequency shear mode [4] which, when excited by the beam, occurs at a frequency where $\epsilon_L(\mathbf{k}, \omega)$ has a large negative value. The enhancement effects are generally small; the most significant enhancement seems to occur in the case of the longitudinal dissipative instability. In the crystalline phase where the longitudinal and transverse modes are coupled owing to the anisotropy of the plasma, a genuine resonant transverse instability may occur. Since the shear phase velocity s is much lower than the longitudinal phase velocity, the transverse and longitudinal instabilities may be excited without interfering with each other [10].

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