29. Oscillations in a Dusty Plasma Medium

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1. ABSTRACT

This chapter discusses novel properties introduced by charged particulates in a plasma medium, and how they influence excitation and propagation of waves. Such a medium, commonly known as a dusty plasma, is generated in the near-Earth environment by dust and other debris of meteoric origin and exhausts and effluents from space platforms. A novel feature of dusty plasmas is that the charge-to-mass ratio can become a dynamical variable, and represents an additional degree of freedom that is not available to a classical plasma. Charged dust particles in a plasma introduce unique potential structures, and significantly alter the short- and long-range forces that can affect the ordering of the dust grains. More interestingly, large amounts of charges on dust grains can allow the average potential energy of the dust component to exceed its average kinetic energy. This can give rise to a strongly coupled plasma component, with liquid-like and solid-like characteristics. These aspects can introduce new types of plasma oscillations or significantly modify existing ones. Selected theoretical and experimental studies of low-frequency electrostatic waves, in weakly and strongly coupled plasmas containing negatively charged dust grains, are used to illustrate the unique oscillations in a dusty plasma medium. The presence of charged dust is shown to modify the properties of ion-acoustic waves and electrostatic ion-cyclotron waves through the quasi-neutrality condition, even though the dust grains do not participate in the wave dynamics. If dust dynamics is included in the analysis, new "dust modes" appear.

2. INTRODUCTION

A plasma is generally considered to be an ensemble of ions and electrons. But, in fact, plasmas often contain large numbers of fine, solid particles, loosely referred to as "dust." For example, the Earth's *D* region (roughly 60 to 100 km altitude) is known to be populated by substantial amounts of dust and other debris of meteoric origin. Low temperatures in this region can lead to condensation of water vapor on these "smoke" particles from meteorites [*Hunten et al.*, *1980*], and can form larger particles, which are observed as noctilucent and polar mesospheric clouds [*Thomas*, *1984*]. In addition, there are large numbers of multi-hydrated molecules and ions, which are capable of attaching free electrons. Any of these species could provide electron-removing mechanisms that might be responsible for the order-of-magnitude changes in the conductivity [*Maynard et al.*, *1981*; *1984*] that have been observed in horizontally stratified layers. It has also been suggested that the strong radar echoes [*Balsley et al.*, *1983*; *Hoppe et al.*, *1990*; *Cho and Kelley*, *1993*] and the electron “bite-outs” [*Ulwick et al.*, *1988*; *Kelley and Ulwick*, *1988*], which are often experienced in the *D* region, can be a direct consequence of highly charged aerosol (or dust) particles in this region forming a layer of dusty plasma [*Havnes et al.*, *1990*].

Typical metallic elements of meteoric origin introduced in the ionosphere are Fe, Al, and Ni [*Castleman*, *1973*]. The solar sources, on the other hand, introduce metallic elements of higher atomic weights, like La, Tu, Os, Yt, and Ta [*Link*, *1973*]. These atomic species are assumed to arise from high-temperature activity on the Sun. Most of the metallic elements introduced in the ionosphere are oxidized easily, forming FeO, AlO, TiO, etc., and are suspected of forming aggregates, which become constituents of the background dust. Dust particles immersed in plasmas and ultraviolet (UV) radiation tend to collect large amounts of electrostatic charges, and respond to electromagnetic forces in addition to all the other forces acting on uncharged grains. The charged dust particles participate in complex interactions with each other and the plasma, leading to completely new types of plasma behavior. Dust particles in plasmas are unusual charge carriers. They are many orders of magnitude heavier than any other plasma component, and they can have many orders of magnitude larger (negative or positive) charges, which can fluctuate in time. They can communicate non-electromagnetic effects (gravity, drag, and radiation pressure) to the plasma that can represent new free-energy sources. Dusty plasmas represent the most general form of space, laboratory, and industrial plasmas.

For a long time, dusty plasmas were mainly of interest to researchers in the astrophysical community. The subject gained popularity in the early 1980s with the Voyager spacecraft observations of peculiar features in the Saturnian ring system (e.g., the radial spokes), which could not be explained by gravitation alone [*Goertz*, *1989*]. This led to the development and successful application of the gravitoelectrodynamic theory of dust dynamics [*Mendis et al.*, *1982*]. In this theory, the finely charged dust particles are about equally influenced by planetary gravity and electromagnetic forces in the rotating planetary magnetosphere. These dynamical studies were complemented in the early 1990s by the study of collective processes in dusty plasmas. This led to the discovery of new modes of oscillations, with wide-ranging consequences for the space environment. It also stimulated laboratory studies that led to the observation of several of these modes, including the very-low-frequency dust acoustic mode. This mode can be made strikingly visual to the naked eye by scattering laser light off the dust [*Thompson et al.*, *1999*]. The role of charged dust in enhanced electromagnetic-wave scattering, and its possible application to the observed enhanced radar backscatter from the high-latitude summer mesopause, where noctilucent clouds are present, have also been investigated [*Tsytovich et al.*, *1989*; *Cho and Kelley*, *1993*].

Perhaps the most fascinating new development in dusty plasmas is the phenomenon of crystallization [*Morfill et al.*, *1998*]. In these so-called “plasma crystals,” micron-sized dust, which is either externally introduced or internally grown in the plasma, acquires large negative charge and forms Coulomb lattices. This was theoretically anticipated by *Ikezi* [*1986*]. Ikezi argued that since a dust grain can hold a large number of electrons, the average potential energy of the dust component can substantially exceed its average kinetic energy. This could result in a sharp deviation in the plasma properties when compared with the familiar Vlasov (or weakly coupled) plasma, and could lead to the formation of a strongly coupled plasma component for ordinary density and room temperatures. Strongly coupled plasma, which otherwise can exist under extraordinary conditions of extremely high densities and very low temperatures (such as a stellar environment), can display liquid-like or solid-like properties. This entirely new material, where phase transition and crystalline structure are so vividly observed by the naked eye, is becoming a valuable tool for studying physical processes in condensed matter, such as melting, annealing, and lattice defects. It also provides a strong motivation for investigating the collective properties in a strongly coupled plasma, an area that has so far remained largely unexplored.

The motivation for studying dusty plasmas is the realization of their occurrence in both the laboratory and space environments. As alluded to earlier, examples include cometary environments, planetary rings, the interstellar medium, and the lower ionosphere. Dust has been found to be a detrimental component of the radio-frequency (RF) plasmas used in the microelectronic processing industry. It may also be present in the limiter regions of fusion plasmas confined in Tokamak devices, as the result of the sputtering of carbon by energetic particles. It is interesting to note that the recent flurry of activity in dusty plasma research has been driven largely by discoveries of the role of dust in quite different settings: the rings of Saturn [*Goertz*, *1989*], and plasma-processing devices [*Selwyn*, *1993*]. The purpose of this article is not to review all possible waves in a dusty plasma medium, but to highlight the novelties of this medium, and to elucidate how wave generation and propagation in this medium are altered from our classical notion. For a detailed review of waves in dusty plasmas, see *Verheest* [*1996*; *2000*]. In Section 3, we first treat waves in weakly coupled dusty plasma. We discuss the case in which the dust charge is time stationary, and then consider the effects of dust-charge fluctuations on the evolution of waves. In Section 4, we discuss waves in a strongly coupled dusty plasma. Finally, in Section 5, we discuss some outstanding issues that need to be resolved in the future.

3. WAVES IN A WEAKLY COUPLED DUSTY PLASMA

In a plasma, the ratio of the potential energy to the kinetic energy of the particles is given by $\Gamma = q^2 \exp\left(-b/\lambda_D\right)/bT$, where $b = \left(3/4\pi\, n\right)^{1/3}$ is the inter-particle distance, *T* is the temperature, λ_D is the Debye length, *q* is the charge, and *n* is the particle density. It is found that if $\Gamma > \Gamma_c$, where Γ_c is a critical value, then there is a phase transition to a solid state, and a Coulomb lattice is formed. Typically, for a Coulomb system, $\Gamma_c \sim 170$ [*Slattery et al*., *1980*], but it has a different value if plasma-shielding effects are considered. *Hamaguchi et al*. [*1997*] discussed phase transition for a screened Coulomb (Yukawa) system. Under normal circumstances for ordinary ion/electron plasmas, Γ is much smaller than Γ_c. Hence, the gaseous state normally prevails, and we are most familiar with this plasma regime. The $\Gamma \ll \Gamma_c$ condition defines the weak-coupling regime. However, in a dusty plasma, a dust grain can acquire a large amount of charge, i.e., $q_d = Ze$, where *Z* can be very large $(\sim 10^3\text{-}10^5)$. Under such circumstances, it is possible for Γ to exceed Γ_c for ordinary dust density and temperature values. This condition $\left(\Gamma \gg \Gamma_c\right)$ defines the strong-coupling regime, where a solid-like behavior, such as Coulomb crystals, is manifested. For intermediate values of Γ, a liquid-like state exists. In the following, we discuss waves in a weakly coupled dusty plasma, first with stationary grain charge in Section 3.1, and we subsequently examine the effects of grain-charge fluctuations on the evolution of waves in Section 3.2.

3.1 Stationary Grain Charge

In this section, we assume that the dust-grain charge does not vary during wave evolution. The presence of charged dust grains acts as a third species, but dust grains can significantly affect the behavior of a plasma in which they are immersed because of their unusual value of charge-to-mass ratio. Both electrons and ions will be collected by the dust grains, but since the electrons move about more swiftly than the ions, the grains tend to acquire a negative charge. Secondary and photoelectron emission from grains in radiative or energetic plasma environments may also contribute to grain charging, and can lead to positively charged grains. As a result, the balance of charge is altered by the presence of the dust, so that the condition for charge neutrality in a plasma with negatively charged grains becomes

$$n_i = n_e + Zn_d \, , \tag{1}$$

where $n_\alpha \left(\alpha = e,i,d\right)$ is the number density of electrons, ions, and dust grains, and $Z = q_d / e$ is the ratio of the charge, q_d , on a dust grain to the electron charge, *e*.

The presence of charged dust can have a strong influence on the characteristics of the usual plasma-wave modes, even at frequencies where the dust grains do not participate in the wave motion. In these cases, the dust grains simply provide an immobile charge-neutralizing background (see Equation (1)). When one considers frequencies well below the typical characteristic frequencies of an electron/ion plasma, new "dust modes" appear in the dispersion relations, derived from either the kinetic or fluid equations for the three-species system consisting of ions, electrons, and charged dust grains. Some of these new modes are very similar to those found in negative-ion plasmas, but with some important differences, unique to dusty plasmas. For example, dusty plasmas in nature tend to be composed of grains with a range of sizes (and shapes!). This means, of course, that one must deal with a range of grain masses and charges.

3.1.1. Low-Frequency Electrostatic Waves in a Dusty Plasma: Theory

3.1.1.1 Dispersion relation

The linear dispersion relation for low-frequency electrostatic waves in a magnetized dusty plasma can be obtained using a multi-fluid analysis [*D'Angelo*, *1990*]. By low frequencies, we mean frequencies on the order of or less than Ω_{ci} , the ion gyrofrequency, and ω_{pi} , the ion plasma frequency. We consider a three-component plasma that is uniform and immersed in a uniform magnetic field, **B**, oriented along the *z* axis of a Cartesian coordinate system. Each species has a mass, m_α ; charge, q_α ; charge state, $Z_\alpha = q_\alpha / e$; density, n_α ; temperature, T_α ; thermal velocity, $v_{t\alpha} = \left(\kappa T_\alpha / m_\alpha\right)^{1/2}$; gyrofrequency, $\Omega_{c\alpha} = eZ_\alpha B / cm_\alpha$; and gyroradius, $\rho_\alpha = v_{t\alpha} / \Omega_{c\alpha}$. All dust grains are assumed to have the same mass and the same negative charge. The three plasma components are described by their continuity and momentum equations:

$$\frac{\partial n_\alpha}{\partial t} + \nabla \bullet \left(n_\alpha \mathbf{v}_\alpha\right) = 0 \, , \tag{2}$$

$$n_\alpha m_\alpha \frac{\partial \mathbf{v}_\alpha}{\partial t} + n_\alpha m_\alpha \left(\mathbf{v}_\alpha \bullet \nabla\right)\mathbf{v}_\alpha + \kappa T_\alpha \nabla n_\alpha + q_\alpha n_\alpha \nabla\varphi - q_\alpha n_\alpha \mathbf{v}_\alpha \times \mathbf{B} = 0\,. \qquad (3)$$

For the low-frequency waves being considered, the electron inertia can be neglected, and we can also take the electron motion to be entirely along **B**. This amounts to assuming that the electrons are in Boltzmann equilibrium, i.e., $\kappa T_e \nabla n_e = e n_e \nabla\varphi$. In addition to the continuity and momentum equations, the charge-neutrality condition (Equation (1)) is also used, both in the equilibrium and in the perturbed state. A standard linear-perturbation analysis is performed around the uniform, non-drifting, equilibrium plasma, with $E_0 = -\left(\nabla\varphi_0\right) = 0$. Assuming that the first-order quantities vary as $\exp\left[i\left(k_x x + k_z z - \omega t\right)\right]$, the following dispersion relation is obtained:

$$\frac{G}{\xi_i^2 - G} + \varepsilon Z_d^2 \mu_{i/d} \frac{H}{\xi_i^2 - \tau_{d/i}\mu_{i/d} H} - \tau_{i/e}\left(1 - \varepsilon Z_d\right) = 0\,, \qquad (4)$$

where

$$G = \left(\frac{\xi_i^2}{\xi_i^2 - 1}\right) k_x^2 \rho_i^2 + k_z^2 \rho_i^2\,, \qquad (4a)$$

and

$$H = \left[\frac{\xi_i^2}{\xi_i^2 - (\xi_i/\xi_d)^2}\right] k_x^2 \rho_i^2 + k_z^2 \rho_i^2\,, \qquad (4b)$$

$\xi_i = \omega/\Omega_{ci}$, $\xi_d = \omega/\Omega_{cd}$, $\mu_{i/d} = m_i/m_d$, $\tau_{d/i} = T_d/T_i$, and $\tau_{i/e} = T_i/T_e$. The parameter $\varepsilon = n_{d0}/n_{i0}$, so that from Equation (1), $n_{e0} = \left(1 - \varepsilon Z_d\right) n_{i0}$. The subscript 0 implies equilibrium (zero-order) quantities. The quantity εZ_d represents the fraction of negative charge per unit volume on the dust. In a plasma without dust – $\varepsilon = 0$ – the dispersion relation (Equation (4)) yields the usual two roots, corresponding to ion-acoustic and electrostatic ion-cyclotron waves. For $\varepsilon \neq 0$, the dispersion relation has four positive solutions in ω/Ω_{ci} , corresponding to electrostatic ion-cyclotron (EIC), ion-acoustic (IA), dust-acoustic (DA), and electrostatic dust-cyclotron (EDC) modes. Numerical solutions of Equation (4) can be obtained for arbitrary values of k_x/k_z , but it is more instructive to obtain the "pure" roots, i.e., those corresponding to propagation either along **B** (acoustic modes), or nearly perpendicular to **B** (cyclotron modes).

Acoustic modes $\left(k_x = 0\right)$: We first obtain the dispersion relations for the acoustic modes, valid in the long-wavelength limits $k\lambda_{De} \ll 1$ and $k\lambda_{Dd} \ll 1$, where $\lambda_{De(d)}$ is the electron (dust) Debye length.

DIA – dust-ion acoustic mode $\left(\omega \gg k_Z v_{td}\right)$: This is the usual ion-acoustic wave, with modifications introduced by the presence of the negatively charged dust [*D'Angelo*, *1990*; *Shukla* *and Silin*, *1992*]. In this case, we can consider the dust to be a static background $(m_d \to \infty)$, yielding the dispersion relation

$$\frac{\omega}{k_z} = \left[\frac{\kappa T_i}{m_i} + \frac{\kappa T_e}{m_i (1 - \varepsilon Z_d)} \right]^{1/2} = C_{S,d} \, , \tag{5}$$

where $C_{S,d}$ is the dust-modified ion acoustic speed. Note that the wave phase velocity, ω/k_z, of the DIA wave increases with increasing relative dust concentration, ε. One can see this by writing the linearized momentum equation for the ions in the form $m_i n_{i0} \partial \mathbf{v}_{i1}/\partial t = -[\kappa T_i + \kappa T_e/(1-\varepsilon Z)](\partial n_{i1} \partial x)$, where the Boltzmann relation has been used to express the wave electric field, E_1, in terms of $\partial n_{e1}/\partial t$. The subscript 1 implies fluctuating (first-order) quantities. The term $m_i n_{i0} \partial v_{i1}/\partial t$ is the force per unit volume on a typical ion fluid element in the presence of the wave perturbation. The right-hand side of the equation is the acoustic restoring force per unit volume on the fluid element, which increases with increasing ε. Physically, as more and more electrons become attached to the immobile dust grains, fewer electrons are available to neutralize the ion space-charge perturbations. An increase in the restoring force then gives rise to an increase in the wave frequency. This increases the wave phase speed, thereby removing the wave from Landau resonance. *Chow and Rosenberg* [*1995*] interpret the term $\kappa T_e/(1-\varepsilon Z_d)$ as an effective electron temperature.

DA – dust-acoustic mode $(\omega \ll k_z v_{ti})$**:** This is a very-low-frequency acoustic mode, in which the dust grains participate directly in the wave dynamics [*Rao et al.*, *1990*]. For this mode, both the electron and ion inertia can be neglected, and the dust provides the mode inertia. The restoring force is provided by the electron and ion pressures, and is described by the linearized dust-momentum equation, with $T_d = 0$: $m_d n_{d0}(\partial v_{d1}/\partial t) = -[\kappa T_e(\partial n_{e1}/\partial x) + \kappa T_i(\partial n_{i1}/\partial x)]$. The dispersion relation is

$$\frac{\omega}{k_z} = \left[\frac{\kappa T_d}{m_d} + \varepsilon Z_d^2 \frac{\kappa T_i}{m_d} \frac{1}{1 + (T_i/T_e)(1-\varepsilon Z_d)} \right]^{1/2} = C_{DA} \, , \tag{6}$$

where C_{DA} is the dust acoustic velocity.

Cyclotron modes $(k_z \ll k_x)$**:** These are modes that propagate nearly perpendicular to the **B** field, but with a finite k_z, so that the assumption that the electrons remain in Boltzmann equilibrium along **B** remains valid.

EDIC – electrostatic dust-ion cyclotron mode $(\omega \sim \Omega_{ci})$**:** This is the dust-modified EIC mode. For $\omega \sim \Omega_{ci}$, the dust grains can be taken as immobile, and the dispersion relation reduces to

$$\omega^2 = \Omega_{ci}^2 + k_x^2 \left[\frac{\kappa T_i}{m_i} + \frac{\kappa T_e}{m_i \left(1 - \varepsilon Z_d\right)} \right]. \tag{7}$$

In Equation (7), we note that the frequency increases with increasing ε.

EDC – electrostatic dust-cyclotron mode $(\omega << \Omega_{ci})$**:** For this mode, the dynamics of the magnetized dust grains must be taken into account. The ions can be taken to be in Boltzmann equilibrium along **B** in response to the very small, but finite, E_z. The dispersion relation is

$$\omega^2 = \Omega_{cd}^2 + k_x^2 \left[\frac{\kappa T_d}{m_d} + \varepsilon Z_d^2 \frac{\kappa T_i}{m_d} \frac{1}{1 + \left(T_i / T_e\right)\left(1 - \varepsilon Z_d\right)} \right]. \tag{8}$$

Summary: The four dispersion relations can be expressed compactly in the form

DIA: $\omega^2 = k_z^2 C_{S,d}^2$ (9a)

DA: $\omega^2 = k_z^2 C_{DA}^2$ (9b)

EDIC: $\omega^2 = \Omega_{ci}^2 + k_x^2 C_{S,d}^2$ (9c)

EDC: $\omega^2 = \Omega_{cd}^2 + k_x^2 C_{DA}^2$ (9d)

where $C_{S,d}$ and C_{DA} are defined in Equations (5) and (6). The modes described by Equations (9a) and (9c) are modes in which the dust dynamics does not play a role, although the effect of the dust can be very important in the excitation of the waves. The DA, Equation (9b), and EDC, Equation (9d), modes are the new low-frequency dust modes.

3.1.1.2 Wave Excitation and Damping

The fluid analysis just presented describes the normal modes of the system. To examine the conditions for wave excitation, a Vlasov analysis is required, except when dealing with resistive or non-resonant instabilities. *Rosenberg* [*1993*] investigated the conditions for the excitation of DIA and DA waves by ion and/or electron drifts in an unmagnetized dusty plasma by using a standard Vlasov analysis.

For the DIA wave, the presence of negatively charged dust reduces the strength of the (collisionless) Landau damping. Thus, even in plasmas with $T_e = T_i$ (where Landau damping gives rise to spatial attenuation of the wave over a distance of less than one wavelength), the waves can, in the presence of negatively charged dust, propagate over several wavelengths. This result could, of course, be anticipated from the fluid analysis, which showed that the DIA phase velocity increased with increasing dust density, thus reducing the importance of Landau damping due to wave particle interactions at $\omega/k \approx v_{ti}$. The Vlasov analysis for the DA mode gave similar results. It was shown that DA waves could be driven unstable by weak ion and electron drifts greater than the DA phase speed. This analytical result was confirmed by *Winske et al*. [*1995*], in a one-dimensional particle simulation of a dusty plasma that included electron and ion drifts.

The effect of charged dust on the collisionless electrostatic ion cyclotron instability (EDIC) was investigated by *Chow and Rosenberg* [*1995*; *1996a*]. The critical electron-drift velocity in the presence of either positively or negatively charged dust was determined. For the case of negatively charged dust, they found that the critical drift decreased as the relative concentration of the dust increased. This showed that the mode is more easily destabilized in a plasma containing negatively charged dust. This result could again be surmised from the fluid analysis, which showed that the EDIC mode frequency increased with relative dust concentration, reducing the (collisionless) cyclotron damping that is most important for frequencies close to Ω_{ci}. To the best of our knowledge, the Vlasov theory of the EDC mode has not been reported. However, we might be able to draw some conclusions concerning the EDC instability from the work of *Chow and Rosenberg* [*1996b*] on the heavy-negative-ion EIC mode, excited by electron drifts along the magnetic field. The mode frequency again increases with increasing ε, whereas the critical electron drift decreases with increasing ε. The maximum growth rate was found to shift to larger perpendicular wavelengths with increasing ε.

In many of the laboratory dusty plasma environments that have been investigated, the plasmas are only weakly ionized. Hence, the effects of collisions between charged particles, including the dust, and the neutral gas atoms must be considered. These plasmas also often contain quasistatic electric fields that may excite current-driven (resistive) instabilities [*Rosenberg*, *1996*; *D'Angelo and Merlino*, *1996*; *Merlino*, *1997*]. Using kinetic theory, and taking into account collisions with the neutrals, *Rosenberg* [*1996*] investigated an ion-dust streaming instability that might occur at the plasma-sheath interface of a processing plasma. A similar situation was considered by *D'Angelo and Merlino* [*1996*], who analyzed a dust acoustic instability in a four-component fluid plasma. This plasma consisted of electrons, ions, negatively charged dust, and neutrals, with an imposed zero-order electric field. Relatively small electric fields, which can generally be found in typical laboratory plasmas, were required to excite the DA instability. Similar results were found by *Merlino* [*1997*], who studied the excitation of the DIA mode in a collisional dusty plasma. For one particular set of parameters, the critical electron-drift speed was decreased by a factor of three for a plasma in which 90% of the negative charge was on dust grains, compared to a plasma with no dust.

Finally, we briefly discuss two of the novel wave-damping mechanisms that may arise in a dusty plasma. The first is the so-called "Tromsø damping" [*Melandsø et al*., *1993*] for dust acoustic waves. This mechanism is related to the fact that the charge on a dust grain may vary in response to oscillations in the electrostatic potential of the wave. A finite phase shift between the potential and the grain-charge oscillations leads to wave damping, particularly for wave periods comparable to the characteristic grain-charging time. The effects of charging dynamics on waves are treated in more detail in the next section. As pointed out by *D'Angelo* [*1994*], Tromsø damping may also be an important damping mechanism for the DIA mode. Another damping mechanism for the DIA mode, which is related to the fact that the dust grains continuously absorb electrons and ions from the plasma, is the "creation damping" of *D'Angelo* [*1994*]. This effect is due to the continuous injection of new ions to replace those that are lost to the dust grains. These newly created ions cause a drain on the wave, since some of the wave energy must be expended in bringing them into concert with the wave motion. This damping mechanism is expected to be the dominant one for some typical laboratory dusty plasmas.

3.1.2. The Effect of Negatively Charged Dust on Electrostatic Ion Cyclotron Waves and Ion Acoustic Waves: Experiments

3.1.2.1 The Dusty Plasma Device (DPD)

The dusty plasma device is an apparatus for introducing dust grains into a plasma. It consists of a single-ended Q machine and a rotating dust dispenser (shown schematically in Figure 1). The plasma is formed in the usual manner, by surface ionization of potassium atoms from an atomic-beam oven, on a hot (~2200 K) 6-cm-diameter tantalum plate that also emits thermionic electrons. The electrons and K^+ ions are confined to a cylindrical column about 1 m in length, by a longitudinal magnetic field with a strength up to 0.35 T. Typically, the electron and ion temperatures are $T_e \approx T_i \approx 0.2$ eV, with plasma densities in the range of 10^8 to $10^{10}\,\mathrm{cm}^{-3}$. As in typical single-ended Q machines, the plasma drifts from the hot plate with a speed between one and two times the ion-acoustic speed.

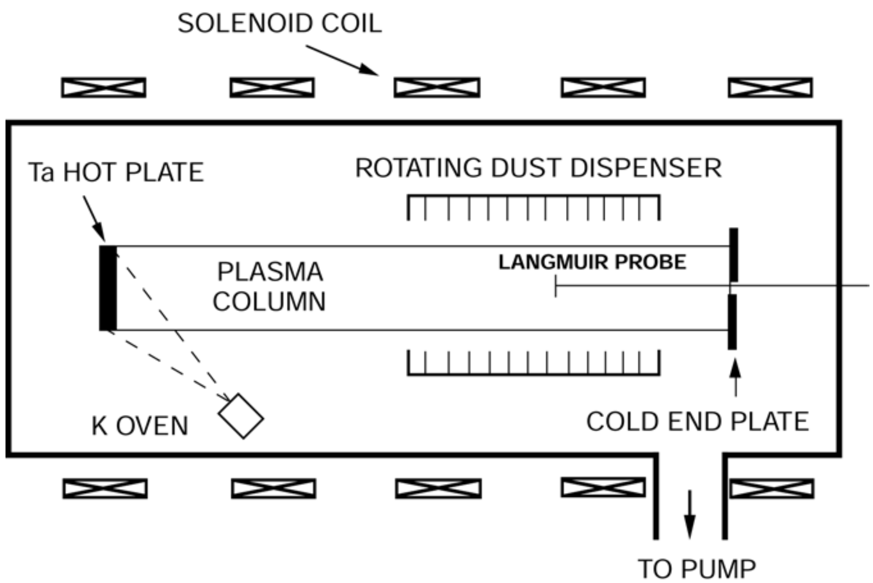

Figure 1. A schematic diagram of the dusty plasma device (DPD) (from *Merlino et al.*, *1998*).

To produce a dusty plasma, kaolin (aluminum silicate) powder is dispersed into a portion of the plasma column. Electron-microscope analysis of samples of the kaolin dust show that the grains are irregular in shape, with sizes ranging from a fraction of a micron to tens of microns. The average grain size is on the order of a few microns. The grains are dispersed into the plasma using the rotating dust dispenser shown in Figure 1. The dispenser consists of a 30-cm-long cylinder surrounding a portion of the plasma column [*Xu et al.*, *1992*]. This cylinder is divided into a number of slots that contain the kaolin powder. A stationary mesh, with an inner diameter slightly smaller than that of the rotating cylinder, also surrounds the plasma column. When the cylinder is rotated, the dust grains are continuously deposited on the outer surface of the stationary mesh. Bristles attached to the rotating slots scrape the outer surface of the mesh, causing it to vibrate, and gently allowing the dust grains to sift through it and fall into the plasma. The fallen dust is then returned to the cylinder through the bottom of the mesh, and re-circulated through the plasma. The amount of dust dispersed into the plasma increases as the rotation rate of the cylinder is increased. The grains attain their equilibrium charge while falling through a very thin layer at the top of the plasma column. The negatively charged dust grains remain in the plasma for a sufficient length of time (~ 0.1 s) to affect the behavior of electrostatic plasma modes, although not long enough to study processes involving dust dynamics.

As pointed out earlier, the quantity $\varepsilon Z_d = 1 - n_e/n_i$ is the fraction of negative charge per unit volume in the plasma on the dust grains. The ratio n_e/n_i can be determined from Langmuir-probe measurements of the reduction in the electron saturation current that occurs when the dust is present, compared to the case with no dust.

3.1.2.2 Current-Driven Electrostatic Dust Ion-Cyclotron Waves (EDIC)

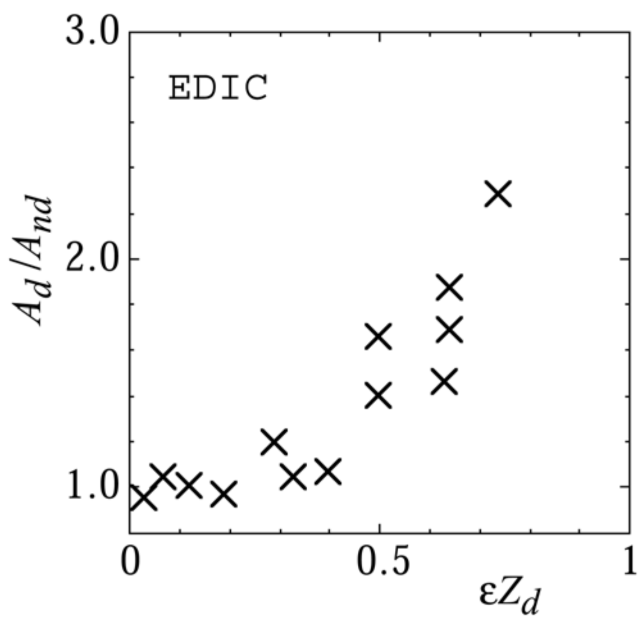

Figure 2. The electrostatic dust ion-cyclotron (EDIC) wave amplitude with dust, divided by the amplitude without dust, as a function of εZ_d (from *Merlino et al.*, *1998*).

The electrostatic ion-cyclotron instability is produced by drawing an electron current along the axis of the plasma column to a 5-mm-diameter disk located near the end of the dust dispenser farthest from the hot plate. A disk bias ~ 0.5 to 1 V above the space potential produces an electron drift sufficient to excite electrostatic waves with a frequency slightly above the ion gyrofrequency. These waves propagate radially outward from the current channel, with a wave vector that is nearly perpendicular to the magnetic field. To study the effect of the dust on the instability [*Barkan et al.*, *1995a*], the wave amplitude, $A_{nd} \equiv (\delta n/n)_{nd}$ (with no dust present), was measured. Without introducing any other changes in the plasma conditions, the dust dispenser was turned on, and the wave amplitude, $A_d \equiv (\delta n/n)_d$ (with dust), was measured. The ratio A_d/A_{nd} could then be used as an indication of the effect of the dust. This procedure was repeated for various dust-dispenser rotation rates. For each value of the rotation rate, the quantity εZ_d was determined from measurements made with a Langmuir probe, located in the dusty plasma. Figure 2 shows the results of these measurements. It appeared that as more and more electrons became attached to the dust grains (larger εZ_d 's), it became increasingly easier to excite EDIC waves, in the sense that, for a given value of the electron-drift speed along the magnetic field, the wave amplitude was higher when the dust was present. By lowering the disk bias to the point that the electron drift was insufficient to excite the waves *with the dust off*, it was possible, *by simply turning the dust on*, to excite the EDIC waves. This result was in line with the prediction of *Chow and Rosenberg* [*1995*; *1996a*], namely that the presence of negatively charged dust reduced the critical electron drift for excitation of the EDIC mode.

3.1.2.3 Ion Acoustic Waves (DIA)

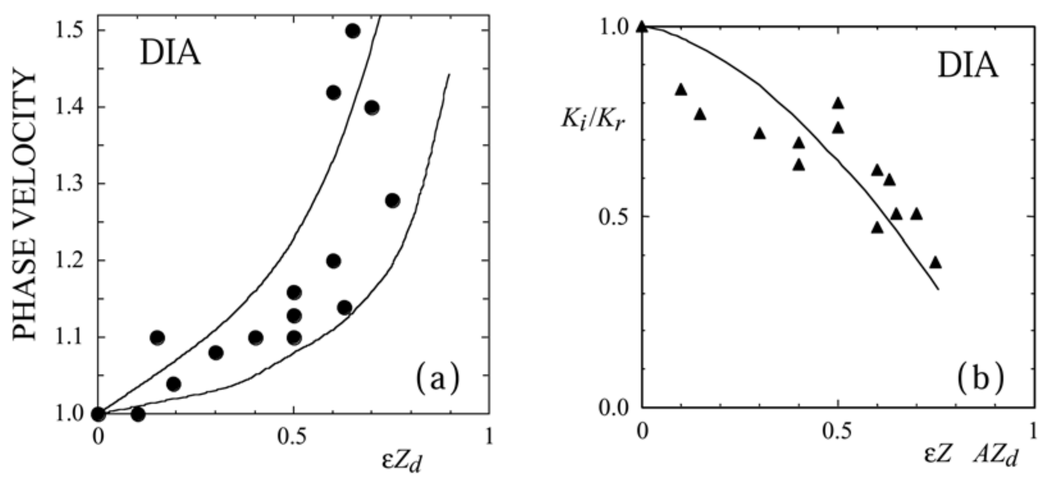

Figure 3. The dispersion properties of grid-launched dust ion-acoustic (DIA) waves. (a) The measured phase velocity as a function of εZ_d (dots). The solid curves were obtained from fluid theory for the case of no plasma drift along the magnetic field (upper curve), and for a drift equal to twice the acoustic speed (lower curve). (b) The measured spatial damping rate (K_i) normalized to the wavenumber (K_r) as a function of εZ_d (triangles). The solid curve was obtained from the solution of the Vlasov equation. In both (a) and (b), the values were normalized to the respective values for $\varepsilon Z_d = 0$ (no dust case) (from *Merlino et al.*, *1998*).

Ion acoustic waves were launched into the dusty plasma by means of a grid that was located approximately 3 cm in front of the dust dispenser (hot-plate side), and oriented perpendicular to the magnetic field [*Barkan et al.*, *1996*]. The grid was biased at several volts negative with respect to the space potential, and a sinusoidal ($f \sim 20$ to 80 kHz) tone burst, of about 4 to 5 V peak-to-peak amplitude, was applied to it. This produced a density perturbation near the grid that traveled down the plasma column as an ion-acoustic wave. By using an axially movable Langmuir probe, the phase velocity $\left(v = \omega/k_r\right)$, wavelength $\left(\lambda = 2\pi/k_r\right)$, and spatial attenuation length $\left(\delta = 2\pi/k_i\right)$ could be measured as a function of the dust parameter εZ_d (k_r and k_i are the real and imaginary parts of the wavenumber.) Figure 3 shows the variation of the phase speed and the spatial-damping parameter, k_i/k_r, with εZ_d. Both quantities were normalized to their respective values in the absence of dust $(\varepsilon Z_d = 0)$. As the fraction of negative charge per unit volume on the dust *increased*, the wave phase velocity *increased* and the wave damping *decreased*. The solid lines in Figure 3a are curves obtained from fluid theory, for the case of no plasma drift along the magnetic field (upper curve), and for a drift of twice the acoustic speed (lower curve). The solid curve in Figure 3b was obtained from the solution of the Vlasov equation. The reduction in the wave damping was a consequence of the reduction in Landau damping that accompanied the increase in phase velocity with increasing εZ_d.

3.1.3. Observations of the Dust Acoustic Wave (DAW)

To observe the low-frequency dust acoustic mode, it was necessary to develop a method for trapping dust grains within a plasma for long times. The initial observations of the DAW were performed using a modified version of the DPD described earlier, in which an anode double layer was formed near the end of the plasma column [*Barkan et al*., *1995a*; *1995b*]. The negatively charged dust grains were trapped in the positive potential region of the anode glow, and DA waves were spontaneously excited, probably due to an ion-dust streaming instability [*Rosenberg*, *1993*].

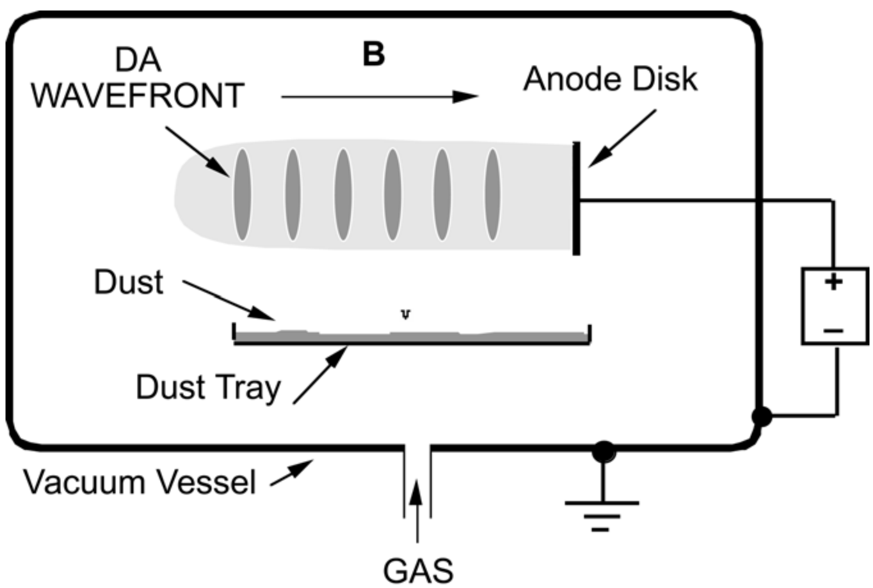

Figure 4a. Experimental observation of the dust acoustic mode: A schematic diagram of the glow-discharge device used to trap negatively charged dust.

Further experiments on the DAW were made in the device shown schematically in Figure 4a [*Thompson et al*., *1997*]. A glow discharge was formed in nitrogen gas ($p \approx 100$ mTorr) by applying a positive potential (200 to 300 V, 1 to 25 mA) to a 3-cm-diameter anode disk, located in the center of a grounded vacuum chamber. A longitudinal magnetic field of about 100 G provided some radial confinement for the electrons, resulting in a cylindrical rod-shaped glow discharge along the magnetic field. Dust grains from a tray located just beneath the anode were attracted into the glow discharge and trapped in this positive potential region. The dust cloud could be observed visually, and its behavior was recorded on VCR tape by light scattered from a high-intensity source, which illuminated the cloud from behind. The trapped grains had a relatively narrow size distribution, with an average size of about 0.7 μm, and a density on the order of $10^5\,\mathrm{cm}^{-3}$.

When the discharge current was sufficiently high ($>1\,\mathrm{mA}$), DA waves appeared spontaneously in the dusty plasma, typically at a frequency of $\approx 20\,\mathrm{Hz}$, with a wavelength of $\approx 6\,\mathrm{mm}$, and propagated at a phase velocity of ≈ 12 cm/s. They were observed as bright bands of enhanced scattered light (from the wave crests), traveling along the axis of the device away from the anode, as shown in Figure 4b.

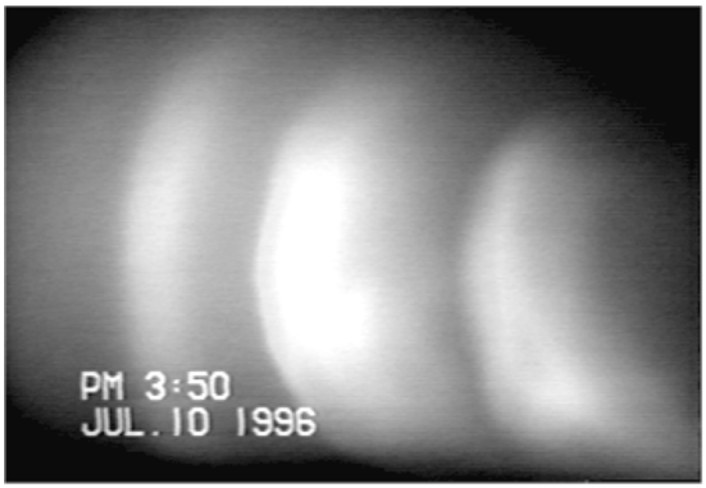

Figure 4b. Experimental observation of the dust acoustic mode: A single frame from the dust acoustic wave video image (from *Merlino et al.*, *1998*).

To investigate the properties of the waves in more detail, a sinusoidal modulation (in addition to the dc bias) was applied to the anode, to generate waves with frequencies in the range of 6 to 30 Hz. At each frequency, a video recording of the waves was made, and the wavelength was measured. Figure 5 shows a plot of the resulting wavenumber (k) as a function of the angular frequency (ω). Over this frequency range, below ω_{pd} $(=\left(4\pi n_d e^2 Z^2 m_d\right)^{1/2}$, the dust plasma frequency), the waves were non-dispersive and had a phase velocity ≈ 12 cm/s. The data in Figure 5 were compared to a theoretical dispersion relation for DA waves, taking into account collisions between the dust grains and neutral gas molecules [*Rosenberg*, *1996*; *D'Angelo and Merlino*, *1996*; *Wang and Bhattacharjee*, *1997*]: $\omega\left(\omega + i\nu_{dn}\right) = k^2 C_{DA}^2$, where ν_{dn} is the dust-neutral collision frequency, and C_{DA} is the dust acoustic speed, defined in Equation (5). Here, $T_e = 2.5$ eV and $T_i = 0.05\,\mathrm{eV}$. ν_{dn} was computed from an expression obtained by *Baines et al.* [*1965*]. $\nu_{dn} \approx 4 m_n N a^2 v_n m_d$, where m_n and v_n are the mass and thermal speed of the neutrals, and N is the neutral density. Using $m_n = 5\times10^{-26}$ kg, $v_n = 283$ m/s, $N = 3\times10^{21}\ \mathrm{m}^{-3}$, a = 0.35μm, and $m_d = 3.6\times10^{-16}$ kg, we find $\nu_{dn} \approx 60\ \mathrm{s}^{-1}$. The solid curve in Figure 5 was then obtained by taking ω to be real and k to be complex, and solving for the wavenumber, k_r, as a function of ω :

$$k_r = \frac{1}{\sqrt{2}C_{DA}}\sqrt{\omega\left(\omega+\sqrt{\omega^2+\nu_{dn}^2}\right)}, \tag{10}$$

with C_{DA} = 12 cm/s. The effect of the collisions produces an offset, which explains why the data points do not extrapolate *linearly* through the origin.

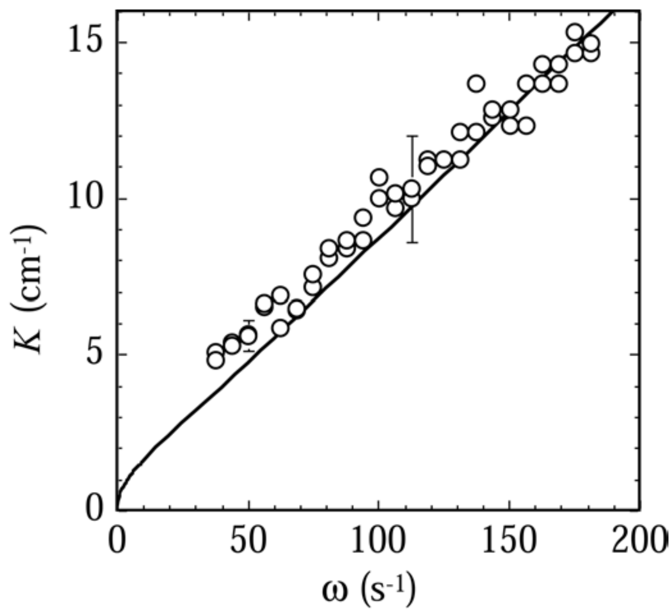

Figure 5. The **m**easured (open circles) dust acoustic-wave dispersion relation (wavenumber k as a function of angular frequency ω). The solid curve was computed from the fluid dispersion relation given in Equation (10) (from *Merlino et al.*, *1998*).

Finally, we note that spontaneous appearance of the DAW may also be explained by the collisional theory of *D'Angelo and Merlino* [*1996*], when account is taken of the equilibrium longitudinal electric field, E_0, in the plasma. This theory predicts the DAW instability for values of $E_0 > 1$ V/cm, quite close to the fields, $E_0 \approx 2$ to 5 V/cm, when measured with an emissive probe.

3.2 Effects of Dust-Charge Fluctuations

In Section 3.1, we described the wave properties in a dusty plasma, in which the effects of dust-charge fluctuation were ignored. Consequently, as far as the waves were concerned, the dust represented a third species in essentially a multi-species plasma. However, in many important cases, such as meteor ablation in the *D* region [*Havnes et al.*, *2001*], discharge of effluents and exhausts from space platforms [*Horanyi et al.*, *1988*; *Bernhardt et al.*, *1995*], etc., the time scales of dust charging and the waves may be comparable. Physically, these situations represent an expansion of a dust cloud through a plasma medium. In such cases, the dust-charging dynamics are likely to influence the wave properties, and vice versa. This led to the development of a self-consistent formalism to account for the dust-charging physics of wave excitation and propagation in a dusty plasma [*Jana et al*., *1993*; *Varma et al*., *1993*].

It is important to realize that, in general, both the charge and mass of the dust grain could be time-dependent. Hence, the Lorentz force on the dust grain is time-dependent, even when the electric and magnetic fields are stationary. For simplicity of discussion, in the following we assume that the mass of the dust grain is time-independent. Thus, the equation of motion for a dust grain is given by

$$\frac{dv}{dt} = \frac{q_d\left(t\right)}{m}\left(E + v \times B\right). \tag{11}$$

To obtain the instantaneous value of dust charge, $q_d\left(t\right)$, we have to solve the dust-charging equation, given by

$$\frac{dq_d}{dt} = \sum_j I_j \ , \tag{12}$$

where the I_j are currents arriving on the surface of the dust grain. The I_j could be of various origins. For example, they could be due to thermal fluxes of electrons and ions, secondary electron emissions, photoelectrons, backscattered electrons, etc. [*Horanyi et al*., *1988*]. The steady state value of the charge is obtained from the $dq_d\left(t\right)/dt = 0$ condition.

To illustrate the important new physics that is introduced by dust-charging dynamics, consider a simple case in which the charging currents are due only to electron and ion fluxes arriving on the dust grain. We express the electron and ion currents as

$$I_e = -\pi a^2 e\left(\frac{8\kappa T_e}{\pi m_e}\right)^{1/2} n_e \exp\left[\frac{e\varphi_f}{\kappa T_e}\right], \tag{13}$$

$$I_i = \pi a^2 e\left(\frac{8\kappa T_i}{\pi m_i}\right)^{1/2} n_i \left[1 - \frac{e\varphi_f}{\kappa T_i}\right], \tag{14}$$

where a is the grain radius and φ_f is the floating potential on the dust grain with respect to the local plasma. Expressions (13) and (14) were derived on the basis of orbit-motion-limited (OML) theory [*Laframboise and Parker*, *1973*; *Allen*, *1992*], which assumes a steady state; collisionless electrons and ions; all ions come in from $r = \infty$, where the potential $\varphi = 0$, and therefore have positive energy; the electron and ion distributions are Maxwellian at $r = \infty$; and the trajectory of any ion or electron can be traced back to $r = \infty$ without encountering grain surfaces or potential barriers. More recently, the validity of the OML theory has been extensively scrutinized [*Allen et* *al*., *2000*]. It was found that as long as the grain size is much smaller than the Debye length and for typical values of T_i and T_e, OML theory remains valid [*Lampe*, *2001*; *Lampe et al*., *2001a*]. For ionospheric applications, these conditions are generally well satisfied.

To consider the effects of dust charging on collective effects, the fluctuating (linearized) currents can be expressed as [*Jana et al*., *1993*]

$$I_{1e} = I_{e0}\left(\frac{n_{1e}}{n_{0e}} + \frac{e\varphi_{1f}}{\kappa T_e}\right), \tag{15}$$

$$I_{1i} = I_{i0}\left(\frac{n_{1i}}{n_{0i}} - \frac{e\varphi_{1f}}{w_0}\right), \tag{16}$$

where $w_0 = \left(\kappa T_i - e\varphi_{f0}\right)$ and $I_{i0} = -I_{e0}$ because of equilibrium constraints. By using Equations (15) and (16) in Equation (12), it can be shown that

$$\frac{dq_{1d}}{dt} + \eta q_{1d} = \left|I_{e0}\right|\left(\frac{n_{1i}}{n_{0i}} - \frac{n_{1e}}{n_{0e}}\right). \tag{17}$$

Here, η is the dust-charge relaxation rate, which is given by

$$\eta = \frac{e\left|I_{e0}\right|}{C}\left(\frac{1}{\kappa T_e} + \frac{1}{\kappa T_i - e\varphi_{f0}}\right), \tag{18}$$

where φ_{f0} is the equilibrium floating potential, and C is the capacitance of the dust grain (assumed to be spherical). Physically, η represents the natural decay rate of the dust-charge fluctuations. These fluctuations arise when the electron/ion current onto the grain surface compensates for the deviation of grain potential from the equilibrium floating potential (Figures 6 and 7).

It is clear from Equation (17) that a fluctuation in the current will lead to a fluctuation in the dust charge via plasma-density fluctuations. Equation (17) is an additional dynamical equation that has to be considered along with the standard equations for deriving the wave-dispersion relation. This has an important bearing on the collective behavior, since the linearized Poisson's equation will now include an additional term proportional to the first-order dust-charge fluctuation:

$$-\nabla^2\phi_1 = 4\pi\left(\sum_\alpha n_{1\alpha} e_\alpha + n_{0d} q_{1d}\right), \tag{19}$$

where $e_d = q_{0d}$ and $q_{1d}, n_{1\alpha}, \varphi_1$, represent fluctuations in dust charge, plasma density, and electrostatic potential, respectively.

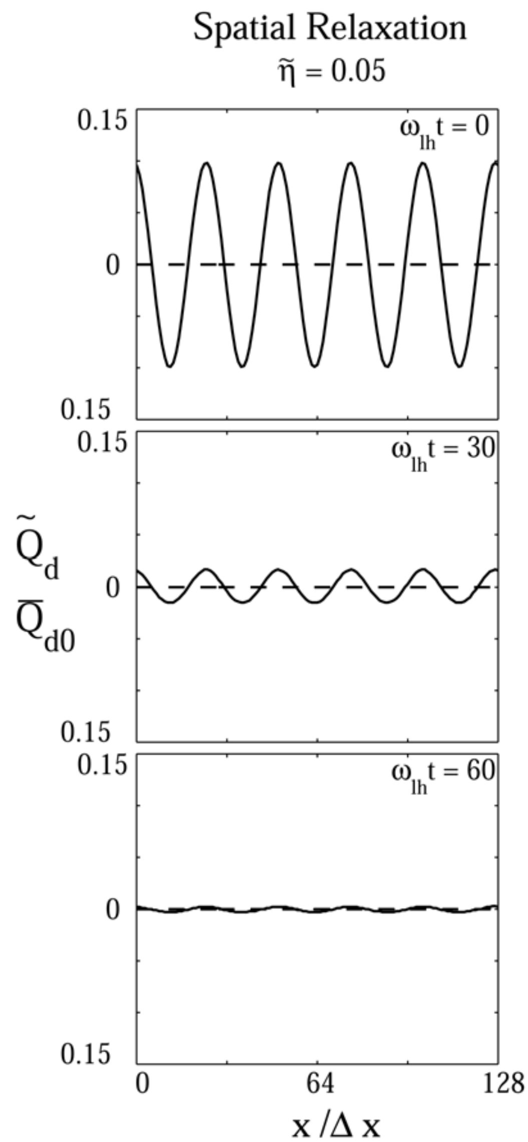

Figure 6. A numerical simulation of the spatial relaxation of an initial spatial dust-charge perturbation $\tilde{Q}_d / Q_{d0}$ at the rate $\eta = 0.05\omega_{1h}$ (from *Scales et al.*, *2001*).

In the simplest form, the local dispersion relation incorporating the effects of dust-charge fluctuations is given by *Jana et al*. [*1993*]:

$$1+\chi_e(\omega,k)\left(1+\frac{i\beta}{\omega+i\eta}\right)+\chi_i(\omega,k)\left(1+\frac{i\beta}{\omega+i\eta}\frac{n_{e0}}{n_{i0}}\right)+\chi_d(\omega,k)=0\,. \tag{20}$$

Here, χ_i, χ_e, and χ_d are ion, electron, and dust susceptibilities. The new physics introduced by the dust-charge fluctuations is represented by the two new parameters, η and β. While η describes the charge relaxation, the second parameter, $\beta = (|I_{e0}|/e)(n_{d0}/n_{e0})$, represents a dissipation rate that is similar to collisional dissipation. It can lead to damping or generation of waves, depending on the circumstances.

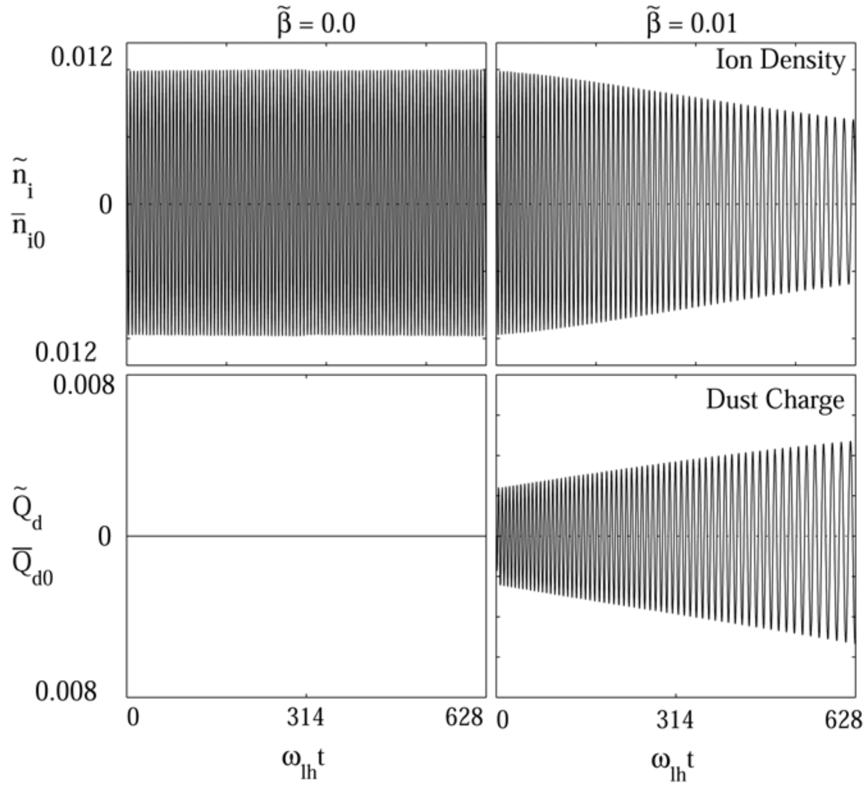

Figure 7. A numerical simulation showing the temporal damping of a lower-hybrid wave due to dust charging. For $\beta = 0.01\omega_{1h}$, the ion-density perturbation for the lower-hybrid wave damps, while the associated dust-charge fluctuations, $\tilde{Q}_d/Q_{d0}$, grow (from *Scales et al., 2001*).

We briefly discuss an application of these ideas to the problem of the expansion of a dust cloud through a plasma environment, since it generically represents a number of important situations in space plasmas. *Scales et al*. [*2001*] have developed a hybrid numerical simulation model that incorporates simple dust-charging physics, and considers the dust charge as a dynamical variable. Both electron and ion components are represented as fluids, while the dust component is treated with the particle-in-cell (PIC) method. Equation (12) is solved at every time step to obtain the instantaneous value of the grain charge, and only electron and ion flux currents, as described in Equations (13) and (14), are used. Poisson's equation is solved to calculate the instantaneous electrostatic potential.

Besides providing the characteristics of an expanding dust cloud, the simulation was also used to validate the fundamental processes of dust-charge relaxation and dissipation. Considering the simple case of lower-hybrid oscillations – for which $\chi_e = \omega_{pe}^2 / \Omega_e^2$ and $\chi_i = \omega_{pi}^2 / \omega^2$, and $\chi_d = 0$, i.e., dust grains are infinitely massive and, hence, they are immobile – it can be shown that there are two fundamental modes [*Scales et al.*, *2001*]. These are $\omega = -i\eta$, the dust-charge fluctuation mode, and $\omega \approx \omega_{lh} - i(\beta/2)(n_{0e}/n_{0i})$, the damped lower-hybrid mode, where ω_{lh} is the lower-hybrid frequency. Simulation was initiated with a sinusoidal dust charge, and its time evolution was monitored. Figure 6 shows the decay of a spatial sinusoidal perturbation in the dust charge at the rate of η. This confirmed the dust-charge relaxation physics. Subsequently, a lower-hybrid wave was launched, and its amplitude was followed in time. Figure 7 shows the ion-density perturbation and dust-charge perturbation in two simulations. The first had no dust charging (i.e., $\beta = 0$), and there was no damping of the wave. In the second case, which was subject to dust charging $(\beta = 0.01\omega_{lh})$, the wave decayed at a rate proportional to β. Also, note that the damping of the wave increased with dust-charge-fluctuation amplitude, as expected. Similar effects of dust charging on ion-acoustic waves have been discussed by *Chae* [*2000*].

Expansion of a dust cloud across a magnetic field was investigated by releasing neutral dust in the middle of the simulation box, and allowing the dust cloud to thermally expand. As the dust particles expanded, they received both electron and ion fluxes on their surface. Since the electrons were faster, the grains got negatively charged. This led to the creation of a narrow layer, less than an ion-gyroradius wide, in which electrons were sharply reduced in density, and there was a corresponding enhancement in the density of the charged dust component. The quasi-neutrality in this layer was maintained by positive ions, electrons, and negatively charged massive dust grains. However, the density gradients of the dust and the electron components had opposite signs. Consequently, while the electron density decreased, the charged dust density increased in the radial direction within the layer. This resulted in an intense transverse ambipolar electric field, localized over a short scale size, on the order of an ion gyroradius [*Ganguli et al.*, *1993*; *Scales et al.*, *1998*]. Because the transverse electric field was localized over such a short scale size, the ions did not experience the electric field in its entirety over one complete gyro orbit and, hence, their $\mathbf{E} \times \mathbf{B}$ drift was not fully developed. The dust grains were so massive that, for all practical purposes, they behaved essentially as an unmagnetized species. However, the electrons experienced a strongly sheared $\mathbf{E} \times \mathbf{B}$ drift. This condition was ideal for exciting the electron-ion-hybrid (EIH) instability around the lower-hybrid frequency, and with wavelengths short compared to the ion gyroradius [*Ganguli et al.*, *1988*; *Romero et al.*, *1992*; *Scales et al.*, *1995*].

Figures 8 and 9 are a transverse (to B) cross-section of the electron and ion densities and dust charge during the development of the EIH instability, at three times during a numerical simulation. (Note that the dust charge, rather than dust density, is shown, since for these time scales, there was no dust motion except for the expansion.) Figure 8 shows a relatively slow-charging-rate case, and Figure 9 shows a relatively high-charging-rate case. The value of $\tilde{\beta} = \beta/\omega_{lh}$ was calculated to be 0.15 in Figure 8, and 0.7 in Figure 9. These values may be used to assess the effects of dust charging on the EIH instability. In both cases, the dust expanded toward the right. It can be seen that as the dust expanded into the background plasma, a distinct layer of enhanced negative charge developed on the front of the dust cloud. Also, there was a slow reduction of negative charge in the core of the cloud as the electron density was reduced (which ultimately reduced the electron density to sustain a dominant electron flux to the grains). In the higher-charging-rate case in

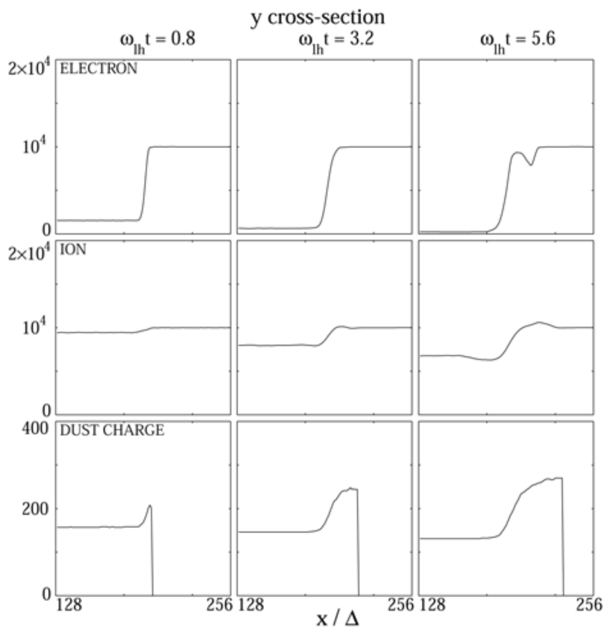

Figure 8. The electron and ion densities and dust charge during a dust-expansion simulation for $\beta = 0.15\omega_{1h}$ (from *Chae*, *2000*).

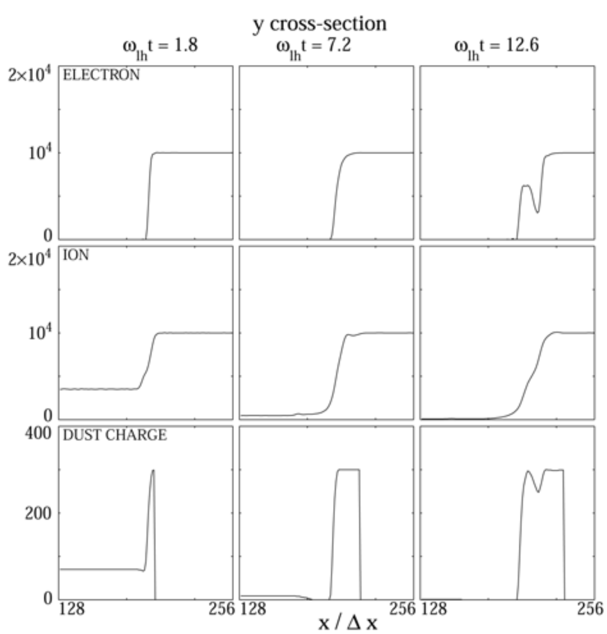

Figure 9. Electron and ion densities and dust charge during a dust expansion simulation for $\beta = 0.7\omega_{1h}$ (from *Chae*, *2000*).

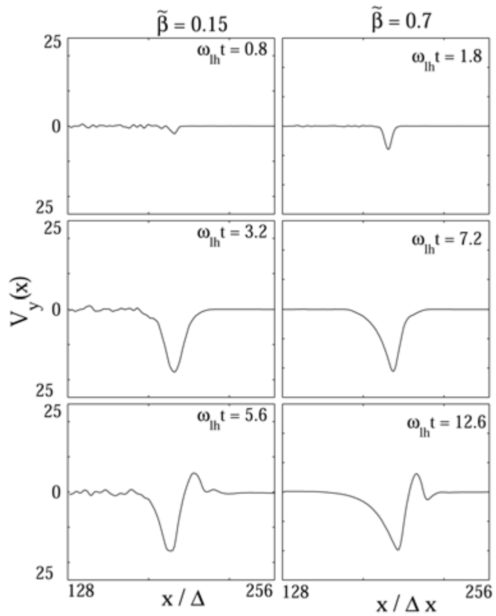

Figure 10. Inhomogeneous electron-flow velocities developing at the boundary layer of an expanding dust cloud for weak and strong charging (from *Chae*, *2000*).

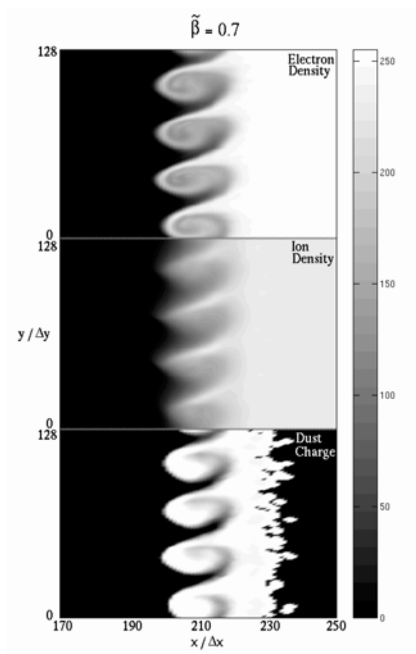

Figure 11. A numerical simulation of the EIH instability in a dusty plasma. The dust expands in the x direction, producing a sheared $\mathbf{E}\times\mathbf{B}$ flow in the $-y$ direction. Note the vortex structures in the dust charge as well as in the ion and electron densities (from *Scales et al.*, *2001*).

Figure 9, a very sharp enhancement in the dust charge was produced at the boundary between the dust cloud and plasma. This enhancement could be localized to less than an ion gyroradius. Note that the structuring, particularly in the electron density, at late times ($\omega_{lh}t = 5.6$ and 12.6) in both cases was due to the development of the EIH instability, which propagated in the *y* direction. Figure 10 shows that highly sheared electron $\mathbf{E}\times\mathbf{B}$ flow was associated with the sharp boundaries in the electron density in Figures 8 and 9. In the higher-charging-rate case, the velocity shear (dv/dx) was larger. This sheared velocity in the dust cloud ultimately drove the EIH instability.

Figure 11 is a snapshot in the morphological evolution of a typical charged-dust layer [*Scales et al*., *2001*]. The EIH instability was spontaneously generated in the layer, and nonlinearly led to coherent (vortex) structures, with typical scale sizes less than an ion gyroradius. Interestingly, *Havnes et al*. [*2001*] discussed radar backscatter from the *D* region from irregularities of a similar scale size, but the origin of such short-scale-size ($\lambda \sim 0.3\,\text{m}$) irregularities is not yet established. Similar small-scale irregularities are a mystery in the meteoric contrails, where a sharp dust layer is also likely [*Kelley et al*., *1998*]. While the dust-expansion mechanism discussed above is an interesting possibility in these near-Earth phenomena, a more detailed comparison of theory with observations is necessary for a positive confirmation.

4. WAVES IN A STRONGLY COUPLED DUSTY PLASMA

As discussed in Section 3, the dust component in a dusty plasma can often be in the strongly coupled regime, where the average Coulomb potential energy between the dust particles can exceed their average thermal energy. This can happen in laboratory dusty plasmas, even at room temperatures, as the result of the large amount of charge that a single dust grain can acquire. An interesting question that naturally arises is what happens to the collective modes of the system – such as those discussed in Sections 3.1 and 3.2 – as the strongly coupled regime is entered. In this section, we discuss the effect of strong coupling on the propagation characteristics of some low-frequency modes of a dusty plasma.

The physical manifestation of strong coupling is to introduce short-range order in the system, as the result of the presence of strong correlations between the particles. This is what distinguishes a solid, with its well-correlated lattice structure, from a gas, where particles behave randomly. The extent of short-range order is controlled by the Coulomb-coupling parameter, Γ, in a dusty plasma. For $\Gamma \ll 1$, the coupling is weak, and the ideal-plasma approximation holds. In such a case, the wave is not affected by the particle correlations, except through weak collisional effects. In the limit when Γ exceeds Γ_c, the short-range order becomes so large that the system freezes to a solid (an ordered crystal structure), and one can then excite lattice vibrations of the dust component in the plasma. We briefly discuss such collective excitations in a dust crystal near the end of this section.

For laboratory dusty plasmas, as well as for dusty plasmas that are commonly found in space, an interesting regime is the one in which $1 < \Gamma \ll \Gamma_c$, where the coupling parameter is less than the crystallization phase, but strong enough to invalidate the weak-coupling assumption of the Vlasov picture. In this intermediate regime, the dust component still retains its fluid character, but develops some short-range order in the system, which keeps decaying and reforming in time. The short-range order gives rise to solid-like properties in the system, e.g., elastic effects, which coexist with the usual fluid characteristics, like viscosity. A phenomenological model incorporating these concepts is the so-called generalized hydrodynamics (GH) model [*Postogna and Tosi*, *1980*; *Berkovsky*, *1992*], which has been successfully used in the past to investigate dense plasmas and liquid metals. We will use this model to study the dynamics of the strongly coupled dust component, while retaining the ordinary hydrodynamics equations for the electron and ion fluids. Such an approach ensures continuity with the multi-fluid approach adopted in the previous sections.

We now replace the linearized momentum equation for the dust component (for an unmagnetized plasma, with $B = 0$) with the following equation given by the GH model:

$$\left(1+\tau_m \frac{\partial}{\partial t}\right)\left[m_d n_{0d} \frac{\partial}{\partial t} v_{d1} + \nabla P_1(r,t) - Z_d e n_{0d} E_1(r,t)\right] = \eta \nabla \bullet \nabla v_{d1} + \left(\varsigma + \frac{\eta}{3}\right)\nabla\left(\nabla \bullet v_{d1}\right), \quad (21)$$

where τ_m is a relaxation time (memory time scale for the decay of the short-scale order), and η and ς are shear- and bulk-viscosity coefficients, respectively. The physical origin of this equation can be traced to the Navier-Stokes equation (to which it reduces, for $\tau_m = 0$), where the Fourier transform of the viscous forces have been generalized from $\left[\eta \mathbf{k}^2 + (\eta/3+\zeta)\mathbf{k}(\mathbf{k}\bullet)\right]$ to $\left[\eta \mathbf{k}^2 + (\eta/3+\zeta)\mathbf{k}(\mathbf{k}\bullet)\right]\big/(1-i\omega\tau_m)$. In real space and time, this amounts to investing the viscosity operator with a non-local character, and leads to memory effects and short-range order. At high frequencies, for example (i.e., for $\omega\tau_m \gg 1$), the dense fluid does not have time to flow, and tends to behave like a solid with elastic properties. At low frequencies, viscous flow is restored. The generalized momentum Equation (21) provides a good physical simulation of this viscoelastic behavior of strongly coupled fluids. The various transport coefficients η, ς, and the relaxation time, τ_m, itself, are functions of the Coulomb parameter, Γ. The exact functional dependence on Γ is model-dependent, and can be deduced either from various first-principle statistical schemes or from fitting to direct molecular-simulation results. The viscoelastic relaxation time is given by

$$\omega_{pd}\tau_m = \frac{\left(\frac{4}{3}\eta + \zeta\right)}{n_{0d}T_{0d}} \frac{\omega_{pd}}{1-\gamma_d \mu_d + \frac{4}{15}u}, \quad (22)$$

where $\omega_{pd} = \left[4\pi\left(Z_d e\right)^2 n_d / m_d\right]^{1/2}$ is the dust plasma frequency, T_{0d} is the dust temperature, γ_d is the adiabatic index, and $\mu_d = \left(1/T_{0d}\right)\left(\partial P/\partial n\right)_T = 1 + u(\Gamma)$ is the compressibility, and $u = E_c / \left(n_{0d}T_{0d}\right)$. E_c is the correlation energy, which is the standard quantity that is calculated from simulations or statistical schemes, and is expressed in terms of an analytically fitted formula. The normalized quantity $u(\Gamma)$ is called the excess internal energy of the system. Typically, for weakly coupled plasmas $(\Gamma < 1)$,

$$u(\Gamma) \approx -\frac{\sqrt{3}}{15}\Gamma^{3/2}. \tag{23}$$

Likewise, in the range of $1 \le \Gamma \le 200$, a widely accepted functional relation is the one provided by [*Slattery et al.*, *1980*]:

$$u(\Gamma) = -0.89\Gamma + 0.95\Gamma^{1/4} + 0.19\Gamma^{-1/4} - 0.81. \tag{24}$$

These relations are based on one-component plasma (OCP) calculations. It should be pointed out that the OCP model ignores Debye shielding effects, which are physically significant for dusty plasmas. For example, the coupling parameter Γ does not include screening effects. A more appropriate model for calculating $u(\Gamma)$ and other transport coefficients would be the Yukawa model, as has been done in the work of *Rosenberg and Kalman* [*1997*]. More recent theoretical [*Kalman et al.*, *2000*] and numerical [*Ohta and Hamaguchi*, *2000*] work on the dispersion of modes in Yukawa fluids for a wide range of wave vectors, coupling, and screening parameters shows how the dispersion curves of longitudinal and shear waves depend on both coupling and screening parameters. However, the shielding contributions do not introduce any fundamental changes in the nature of the effects we have discussed so far. Thus, we will continue to use the simple OCP estimates for our discussions. By using the above expressions, it is straightforward now to carry out a linear stability analysis as before, and to obtain a dispersion relation for low-frequency oscillations of a strongly coupled dusty plasma. We begin by examining the low-frequency dust acoustic modes.

4.1 Dust Acoustic Waves

For longitudinal low-frequency waves $(\omega \ll kv_{te}, kv_{ti})$, with the electrons and ions obeying the Boltzmann law, the dispersion relation for the dust acoustic modes with strong correlation effects is now given by *Kaw and Sen* [*1998*]

$$1 + \frac{1}{k^2\lambda_p^2} - \frac{1}{\omega(\omega + i\nu_{dn}) - \gamma_d \mu_d k^2 \lambda_{Dd}^2 + i\omega k^2 \dfrac{\eta^*}{1 - i\omega\tau_m}} = 0, \tag{25}$$

where all temporal quantities have been normalized by the dust plasma frequency, ω_{pd}, and spatial quantities have been normalized by the inter-particle distance, $b = (4\pi n_{0d}/3)^{-1/3}$. Furthermore, $\eta^* = (4\eta/3 + \varsigma)\big/(m_d n_{0d}\omega_{pd} b^2)$ and $\lambda_p^{-2} = b^2\lambda_p^{-2} = b^2(\lambda_{De}^{-2} + \lambda_{Di}^{-2})$. Equation (25) also includes dust-neutral collision effects (through the normalized collision frequency, ν_{dn}), which can be important in many experimental and space-plasma situations (as discussed in Section 3.1). For $\omega\tau_m \ll 1$, the dispersion relation, Equation (25), simplifies to

$$\omega\left(\omega + i\nu_{dn}\right) = k^2\left(\gamma_d \mu_d \lambda_{Dd}^2 + \frac{\lambda_p^2}{1+k^2\lambda_p^2}\right) - i\omega\eta^* k^2 , \tag{26}$$

and can be readily solved to give

$$\omega_R = \pm\frac{k\lambda_p}{\sqrt{1+k^2\lambda_p^2}}\left[1+\frac{k^2\lambda_{Dd}^2\gamma_d\mu_d - \left(\nu_{dn}+0.5k^2\eta^*\right)^2}{k^2\lambda_p^2}\left(1+k^2\lambda_p^2\right)\right]^{1/2} , \tag{27}$$

$$\omega_I = -\frac{\nu_{dn}+k^2\eta^*}{2} . \tag{28}$$

Note that for $\eta^* = \lambda_{Dd} = v_{dn} = 0$, Equation (27) reduces to the standard dust acoustic result of the weak-coupling regime, namely, $\omega_r = kC_{DA}$, which is Equation (9b). Strong correlations contribute to additional dispersive corrections through the μ_d and η^* terms. As shown in Figure 12, the dispersion curve for the dust-acoustic mode changes significantly with increasing values of Γ. Beyond a certain value of Γ (approximately 3.5), the dispersive corrections change sign. This leads to a turnover in the curve, with the group velocity going to zero, and then to negative values. In addition to dispersive corrections, the mode also suffers additional damping proportional to the η^* contribution, as seen from Equation (28). In the weak-coupling limit, $\Gamma \ll 1$, this term reduces to the usual collisional damping arising from dust-dust collisions, but in the strong coupling limit, it can vary significantly as a function of Γ.

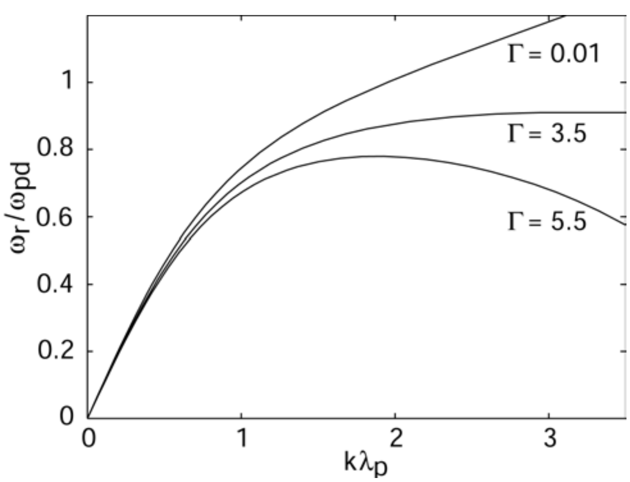

Figure 12. The dispersion curve [Re $\mathrm{Re}(\omega/\omega_{pd})$ as a function of $k\lambda_p$] for the dust-acoustic wave, for different values of Γ (from *Kaw and Sen*, *1998*).

In the opposite limit, of $\omega\tau_m \gg 1$, the dispersion relation, Equation (25), simplifies to

$$\omega^2 = k^2 \left(\gamma_d \mu_d \lambda_{Dd}^2 + \frac{\eta^*}{\tau_m} + \frac{\lambda_p^2}{1+k^2\lambda_p^2} \right). \tag{29}$$

Substituting for η^* and τ_m, this further gives

$$\omega = \pm k \left[\frac{\lambda_p^2}{1+k^2\lambda_p^2} + \lambda_{Dd}^2 \left(1 + \frac{4}{15} u(\Gamma) \right) \right]^{1/2}. \tag{30}$$

In this limit, the dust-acoustic mode does not experience any viscous damping. Dissipation can arise only through Landau damping on the electrons and ions, which are, of course, not included in the GH model. The Γ corrections arise through $u(\Gamma)$, estimates of which can be obtained from the relations of Equations (23) and (24). These corrections, once again, lead to the turnover effect, but the effect is weaker than for the limit.

4.2 Transverse Shear Waves

A solid lattice can support transverse mechanical waves (called shear waves), in addition to longitudinal sound waves. Since strong correlations invest a certain "rigidity," even to the fluid state, can one expect transverse oscillations in a strongly coupled dusty plasma? The answer surprisingly is yes, and the GH model shows this novel behavior quite easily. Taking the curl of the dust equation of motion, Equation (21), and ignoring the negligible electromagnetic contribution arising from the $\nabla \times \mathbf{E}$ term, one immediately gets

$$\left[(1 - i\omega\tau_m) i\omega m_d n_{0d} - \eta k^2 \right] (k \times v_{d1}) = 0 . \tag{31}$$

Since $\mathbf{k} \times \mathbf{v_{d1}} \neq 0$, this leads to a dispersion relation for shear waves that can be written in terms of the dimensionless quantities defined in the previous section:

$$\omega = \frac{-i\eta^* k^2}{1 - i\omega\tau_m}. \tag{32}$$

Its solution is given by

$$\omega_R^2 = \frac{\eta^* k^2}{\tau_m} - \frac{1}{4\tau_m^2}, \tag{33}$$

$$\omega_I = -\frac{1}{2\tau_m}. \tag{34}$$

We thus get a transverse propagating wave for $k^2 > 1/(4\eta^*\tau_m)$, the frequency of which has a linear dependence on **k** in the small-wavelength regime. Substituting for τ_m in Equation (33) and reverting to dimensional variables, we can also express the real frequency (for large **k**) approximately as

$$\omega^2 = k^2\lambda_d^2\omega_{pd}^2 \left\{1-\gamma_d\left[1+\frac{u(\Gamma)}{3}+\frac{\Gamma}{9}\frac{\partial u}{\partial\Gamma}\right]-\frac{4}{15}u(\Gamma)\right\}. \tag{35}$$

In the limit of large Γ, $u(\Gamma) \approx -0.9\Gamma$ is large and negative. For $\Gamma \gg \left|(1-\gamma_d)/(\lambda_d - 4/15)\right|$, the right-hand side of Equation (35) is positive, leading to a propagating wave. In this limit, Equation (35) can be written approximately in the form

$$\omega^2 \approx k^2 \frac{\gamma_d E_c}{m_d n_{0d}}. \tag{36}$$

This is analogous to elastic-wave propagation in solids, with the correlation energy, E_c, playing the role of the elasticity modulus. Shear waves have also been predicted for OCP systems [*Golden and Kalman*, *2000*].

4.3 Dust Lattice Waves

We now briefly discuss wave propagation in a dusty plasma crystal, which is a strongly coupled state, with $\Gamma > \Gamma_c$. Dust crystals have been observed and studied in many laboratory experiments [*Chu and I*, *1994*; *Thomas et al.*, *1994*; *Hayashi and Tachibana*, *1994*]. Typically, they are created in the plasma-sheath region, where the dust particles remain levitated due to a balance between the gravitational force and the electrostatic force of the sheath electric field. Such a state has a high degree of symmetry and order. All of the dust particles are arranged in a hexagonal pattern in a plane, and vertical alignments are perpendicular to the plane. The movement of each dust particle is quite restricted and localized, so that a fluid theory is no longer valid for their description: in fact, they must be treated as discrete entities.

The electrons and ions, on the other hand, continue to be in the weakly coupled regime. They can be considered to be charged fluids that shield the bare Coulomb potential of each dust particle. Because of this shielding, the individual dust particles in a dusty plasma crystal interact only with a few neighboring particles, particularly when the average dust-particle separation, *b*, is larger than the plasma shielding length, λ_p. The collective motion of the lattice sites can then be simply modeled by the excitations of a chain of oscillators (in one dimension) or a network of oscillators (in two dimensions). The one-dimensional model (known as a Bravais lattice model) has been used to study longitudinal oscillations in a dusty plasma [*Melandsø*, *1996*]. Keeping only the nearest-neighbor interaction, the linear equation of motion for each dust particle can be written as

$$\frac{d^2\xi_j}{dt^2} = \frac{\beta(b)}{m_d}\left(\xi_{j-1} - 2\xi_j + \xi_{j+1}\right), \tag{37}$$

where ξ_j is the displacement from equilibrium of the *j*th dust particle, m_d is the particle mass, and

$$\beta(b) = \frac{Q^2}{a^3}\exp\left(-b/\lambda_p\right)\left[1 + \frac{b}{\lambda_p} + \frac{b^2}{2\lambda_p^2}\right]. \tag{38}$$

Assuming a plane-wave solution of the form $\xi_j \propto \exp\left[i\left(\omega t - z_{j,0}k\right)\right]$, where $z_{j,0}$ is the equilibrium location of the *j*th dust particle, provides the dispersion relation

$$\omega^2 = 4\frac{\beta(b)}{m_d}\sin^2\left(\frac{kb}{2}\right). \tag{39}$$

In the long wavelength limit, $kb \ll 1$, the longitudinal dust-lattice wave is seen to have a non-dispersive character, much like the dust-acoustic wave. The phase velocity of the dust-lattice wave (DLW) is given by $v_{DLW} = \left[\beta(b)/m_d\right]^{1/2} b$. A similar dispersion relation has also been obtained for transverse (in-plane) waves, which is valid for horizontal oscillations in a planar crystal [*Nunomura et al*., *2000*]. For out-of-plane (transverse) oscillations, the dispersion relation takes the form [*Vladimirov et al*., *1997*] of

$$\omega^2 = \frac{\gamma}{m_d} - 4\frac{\beta(b)}{m_d}\sin^2\left(\frac{kb}{2}\right), \tag{40}$$

where the constant γ is a measure of the linear restoring force in the vertical direction, and is proportional to the difference between the electrostatic and gravitational force. This transverse wave-dispersion relation has the character of an optical mode (for $kb \ll 1$), but with the frequency decreasing as a function of the wave number. As seen from Equation (38), the phase velocities of the various dust-crystal modes are a function of the plasma screening parameter, $\kappa_1 = b/\gamma_p$, because of the role played by the electron and ion fluids in modifying the inter-dust potential structure.

In summary, we see that low-frequency waves, for which dust dynamics are important, experience significant modifications in the strongly coupled plasma regime. For the longitudinal dust-acoustic waves, for example, there are new dispersive corrections, a lowering of their real frequency and phase velocity, an additional source of damping (due to modified viscous effects), and the existence of parameter regions where $\partial\omega/\partial k < 0$. Strong coupling also gives rise to novel phenomena like the existence of transverse shear modes, even in a fluid medium. These results emerge from the generalized hydrodynamics model for the dust dynamics, where dust correlation effects are physically modeled by Γ-dependent viscoelastic coefficients. However, the basic results are not confined to the GH model but have also been obtained using other methods, such as the quasi-localized-charge approximation (QLCA) [*Rosenberg and Kalman*, *1997*; *Kalman et al*., *2000*], a kinetic approach that uses the so-called “static” and “dynamic” local-field corrections [*Murillo*, *1998*; *Murillo*, *2000*], and the generalized thermodynamic approach [*Wang and Bhattacharjee*, *1997*].

The applicability and physical underpinnings of each of these methods is a subject of much current interest, and provides further motivation for the study of strongly coupled dusty plasmas. In our discussion so far, we have ignored the effect of dust-charge fluctuations on these modes of the strongly coupled regime. As discussed in Section 3.2 for the case of weakly coupled dusty plasmas, charge fluctuations provide an extra degree of freedom in the system, as well as an additional source of free energy that can give rise to interesting new phenomena. These considerations apply to the strongly coupled system, as well. A recent calculation [*Mishra et al.*, *2000*] points out an interesting instability mechanism for the transverse shear mode, arising from finite dust-charging times and the presence of an equilibrium charge gradient. Such effects can facilitate the experimental observation of shear modes.

More recently, there have been a number of experiments and molecular-dynamic simulation studies to understand the collective response of dusty plasmas in the strong coupling regime. One of the earliest experimental measurements on the dust-acoustic mode in such a regime was made by *Pieper and Goree* [*1996*], who excited the plasma with a real external frequency, and measured the complex wave vector. In their experiment, however, they were not able to clearly isolate the strong coupling effects from the collisional effects. This is because the background neutral pressure in their experiment was kept quite high, in order to cool down the dust component through dust-neutral collisions. The damping (as well as the dispersive effects) arising from a large v_{dn} term then tends to mask out the strong coupling contributions (see Equations (27) and (28)). This is one of the major difficulties at present in the unambiguous experimental identification of strong coupling effects. It may be overcome if alternate methods (such as UV ionization) are used for the creation of dusty plasmas.

Numerical simulations, on the other hand, have been more successful in tracking the various strong coupling modifications of the longitudinal dust-acoustic mode. They have also resulted in observation of the transverse shear mode [*Ohta and Hamaguchi*, *2000*]. The transverse shear mode has recently been experimentally seen in a mono-layer dusty plasma [*Nunomura et al.*, *2000*], and also in a three-dimensional fluid-like equilibrium [*Pramanik et al.*, *2000*]. The effects of collisions were also discussed by *Rosenberg and Kalman* [*1997*] and *Kalman et al.* [*2000*]. Both experimental and theoretical investigations of the collective properties of strongly coupled fluid regimes are in their early stages, and a great variety of interesting problems remains to be explored. A particularly attractive area is the experimental exploration of various transport coefficients and their dependence on Γ and κ_1 through measurements of wave-propagation characteristics. This can provide fundamental understanding not only of dusty plasmas, but also of strongly coupled systems, in general.

5. DISCUSSION

In Sections 3 and 4, we discussed how dust introduces new physical phenomena that are markedly different from a classical single- or multi-species plasma. The charge on a dust particle can be very high, yet the charge-to-mass ratio (q/m) of a dust component is much lower than the usual plasma constituent. This itself leads to unusual behavior. While the value of q/m in a multi-species plasma can be different for different components, this value (q/m) remains constant. In a dusty plasma, however, q/m is an additional degree of freedom, which may have a distribution in magnitude and may be time-dependent. In the ionosphere, chemistry is an additional important factor that affects charging and composition of dust particles [*Castleman*, *1973*]. Dust particles charge-exchange with ionospheric ions, and severely alter the local chemistry because of the large charge that can accumulate on a dust particle. In effect, charged particulates can act analogously to catalytic surfaces and promote chemical reactions. Therefore, it is of considerable interest to delineate the physical and chemical processes that affect and transform particulates in a plasma environment.

Also, massive dust particles can no longer be treated as point objects, as in classical plasma treatments. These massive dust grains will have their own unique shielding clouds, which are different from the usual Debye shielding, and which exhibit surface physics and chemistry processes that cannot be ignored. The high dust-charge state can induce strong coupling in the dust component, and thereby introduce short-range ordering in the grains. Research to date indicates that inter-dust-grain forces in a plasma medium are influenced by wakefield force, shadow force, and nonlinear ion flows, and demonstrate the existence of an attractive force between negatively charged dust grains due to wakefield and shadow forces [*Vladimirov and Nambu*, *1995*; *Ishihara and Vladimirov*, *1997*; *Melandsø and Goree*, *1995*; *Lampe et al*., *2000*; *2001a*; *Lampe*, *2001*].

Another important feature that has not been fully addressed is the role of trapped ions [*Goree*, *1992*]. Recent work indicates that the trapped ions can have a profound consequence, and can influence the inter-grain potential and grain-charge magnitude in a plasma medium, provided $T_e \gg T_i$ [*Zobnin et al*., *2000*; *Lampe et al*., *2001b*]. In the future, more research must be conducted to understand and accurately model the effects of trapped ions on the inter-grain potential. The dust grains may lose or gain charges, thereby altering the value of q/m, and affecting various plasma processes such as collective effects, transport, etc. The capacity to lose or gain charges depends greatly on the physio-chemistry of the plasma environment. Hence, charging of dust grains and the influence of the plasma background (e.g., the ionospheric conditions) on this process are very important topics, which must be addressed more thoroughly, both theoretically and experimentally, in the future. Variation of q/m has other important consequences, as well. For example, a distribution of q/m will imply a continuous range of cyclotron frequencies Ω ($= qB/mc$). Thus, the concept of cyclotron resonance and absorption in a dusty plasma is different from that of a classical electron/ion plasma, and its implications on plasma collective effects have yet to be quantified.

Man-made objects (e.g., the Space Station) will increasingly populate the near-Earth space environment in the future. Attitude-control thruster discharges, outgassing, and other water-bearing effluents from these space platforms can lead to the production of ice crystals, and can form a dense "dust" cloud around these objects; these can also charge. Other activities in space, such the as transfer of satellites from low-Earth orbits to higher (geosynchronous) orbits, involve solid-rocket motor (SRM) burns. SRM burns deposit large quantities of aluminum oxide (Al_2O_3) particles in space. Studies show that the flux resulting from just one such SRM burn can exceed the natural meteoroid flux for particles of like size (1 to 10 μm) [*Muller and Kessler*, *1985*]. Another study concludes that 1 to 10 μm-size Al_2O_3 particles can get charged under typical magnetospheric conditions, and can acquire large surface potentials (~ 10 V) [*Horanyi et al*., *1988*]. For a 10 μm-size grain, a surface potential of 10 V corresponds roughly to a charge state of $Z \sim 70000$. These charged grains have a typical residence time of days in the magnetosphere, and can have considerable influence on the plasma environment in the immediate vicinity of these space assets.

One expects that dust will significantly affect collective effects and transport phenomena in space plasmas, if its density exceeds some minimum value. But, to date, no definitive research has been done to establish a general threshold for dust effects. Does the number density of the dust grains need to be a significant fraction of the electron density before effects become observable? This condition would be hard to satisfy in the ionosphere, in most cases. It seems more likely, however, that some much-less-stringent condition applies, e.g., that the total charge density associated with dust grains (which hold very large numbers of electrons when charged) or some other electrodynamic state variable is the controlling factor. However, it is likely that this threshold will be different for different events and, hence, is itself an interesting and important target for future research. Research to date shows that even low dust densities may interfere with space-based systems, by enhancing the scattering cross-section of electromagnetic waves. *Tsytovich et al*. [*1989*] and *Bingham et al*. [*1991*] have shown that even small amounts of charged dust grains can significantly affect the scattering of both electrostatic and electromagnetic waves. They showed that even if the dust-grain mass is assumed to be infinite (and therefore immobile) compared to plasma particles, they are very efficient scattering centers, since only the electron cloud surrounding the grains oscillates in the field of an incoming wave. This is called transition scattering and is, in spirit, similar to the well-known Thomson scattering. For incident waves of wavelength greater than the Debye length $(\lambda > \lambda_D)$, the transition scattering cross-section is enhanced by a factor of Z^2 over Thomson scattering, e.g., $\sigma = Z^2\sigma_T$, where σ_T is electromagnetic wavelengths greater than the Debye length $\left(\lambda > \lambda_D\right)$, the transition scattering the Thomson scattering cross-section. As noted earlier, the dust-charge state can be quite large, typically 10^3 to 10^4 for a 1 to 10 µm-size grain in Earth orbit, leading to huge dust-scattering cross-section enhancements, e.g., $\sigma = \left(10^6\text{-}10^{10}\right)\sigma_T$. Thus, even a small concentration of charged dust may lead to a substantial scattered power, and may potentially interfere with communication and radar capabilities, especially around space assets that are becoming indispensable to our everyday life on Earth. These dust effects will have to be more rigorously assessed and quantified in the future. A general kinetic formalism, necessary to address some of these issues, has recently been developed [*Tsytovich and de Angelis*, *1999*; *2000*; *2001*; *Ricci et al*., *2001*].

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More-specific comments in connection with editing:

In the fifth paragraph of Section 2, you have, “It may also be present in the limiter regions of fusion plasmas....” I’m confused by the term, “limiter regions.” Should this be, “...in limited regions...” or, perhaps, “...in regions at the limits of...?”

In the text below Equation (5), in the linearized momentum equation, you appear to have a vector quantity on the left-hand side of the equation, and only scalar quantities on the right-hand side. Is this correct?

In the last term in parentheses in Equation (31), you have the cross product of two scalars. Shouldn’t this be replaced by $\mathbf{k} \times \mathbf{v}_{d1}$? (as shouldn’t the same expression in the next sentence?)

At the end of Section 5, you have “...$\sigma = Z^2 \sigma_T$, where σ_T is electromagnetic wavelengths greater than the Debye length $(\lambda > \lambda_D)$, the transition scattering the Thomson scattering cross-section.” This doesn’t seem to make sense.

You have only provided the first number page of each reference, rather than the range of page numbers. All other authors have provided the page-number range. It would be nice if you could do so, although I realize time is short.

You may wish to update the following reference(s):
M. Lampe [2001], “Limit of Validity of OML Theory for Small Floating Collector,” *J. Plasma Phys*. (to appear).