Ion-Acoustic Surface Waves on a Dielectric–Dusty Plasma Interface

Konstantin N. Ostrikov and Ming Y. Yu

***Abstract*— The theory of ion-acoustic surface wave propagation on the interface between a dusty plasma and a dielectric is presented. Both the constant and variable dust-charge cases are considered. It is found that massive negatively charged dust grains can significantly affect the propagation and damping of the surface waves. Application of the results to surface-wave generated plasmas is discussed.**

***Index Terms*— Dusty plasma, plasma-wall interaction, surface waves.**

I. Introduction

It is well known that dust grains are ubiquitous in astrophysical plasmas [1], [2] such as the cometary tails, planetary rings, interstellar clouds, and photo-planetary discs. Many theoretical studies on wave propagation and transport in dusty space plasmas have appeared in the last decade [3]–[7]. Recently, there has been a shift of attention to low-temperature plasmas because of the rapidly growing plasma-assisted technologies. The latter, such as plasma-assisted chemical vapor deposition (PACVD) and other thin film processes, are widely used in microelectronics [8]–[10] precision manufacturing. In such processes, the film deposition rate is a key efficiency factor. A crucial problem limiting the deposition rate is the appearance of fine powder in the discharge [10]–[15] when the active species are created at high rates. The dust powder can polymerize in the gaseous state before being deposited onto the substrate as a thin film [10]–[12]. Dust grains (usually regarded as contamination) ranging in size from tenths of microns to several microns have been observed [10], [13], [14]. The obvious need for contamination control in PACVD for modern material processing and manufacturing has stimulated intensive studies on the transport and collective phenomena in dust-containing plasmas.

In this paper we investigate the effect of dust particles on the dispersion and damping characteristics of electrostatic ion-acoustic surface waves (IASW) at the interface between a gas-discharge produced plasma and a dielectric. We shall use both the constant dust-charge approximation, valid for wave frequencies much larger than the rate of charge variation as well as for heavy impurity ions, and the probe model for dust charging. Application of our results to RF produced laboratory plasmas is given.

Manuscript received July 1, 1997; revised August 29, 1997. This work was supported by the Sonderforschungsbereich 191 Niedertemperatur Plasmen and the Science and Technology Center of Ukraine.

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Publisher Item Identifier S 0093-3813(98)00701-2.

II. Impurity Ions or Constant-Charge Dust Grains

We first consider the dispersion properties of IASW in plasmas containing heavy negatively charged impurity ions or constant-charge dust particles. The isotropic plasma with thermal electrons is assumed to be bounded at $x = 0$ by a dielectric with dielectric constant ϵ_0. The dust grains are assumed to have a constant average charge $q_d = Z_d e < 0$, where e is the magnitude of the electron charge. The dust grains are assumed to be much smaller than the electron Debye length r_{De} as well as the interparticle distance. Thus, we can treat the dust grains as heavy (compared to the plasma ions) point masses with constant negative charge. The linearized equations describing the IASW are then

$$d_t n_j + \nabla \cdot (n_{0j} + n_j)\boldsymbol{v}_j = 0 \tag{1}$$

$$d_t \boldsymbol{v}_j + \nu_j \boldsymbol{v}_j = \frac{Z_j e}{m_j}\boldsymbol{E} - \frac{V_{Tj}^2}{n_{0j}}\nabla n_j \tag{2}$$

$$\nabla \cdot \boldsymbol{E} = 4\pi \sum_j Z_j e(n_{0j} + n_j) \tag{3}$$

where $\boldsymbol{E}$ is the electric field of IASW, m_j, n_{0j}, n_j, $\boldsymbol{v}_j$, V_{Tj}, and $q = Z_j e$ are the mass, unperturbed and perturbed densities, fluid velocity, thermal speed, and charge of the species $j = e, i, d$, respectively. In the unperturbed state the quasineutrality condition for a three-component plasma

$$Z_i n_{i0} = n_{e0} + Z_d n_{d0} \tag{4}$$

is fulfilled.

We shall assume electrostatic wave perturbations ($\boldsymbol{E} = -\nabla\varphi$, where φ is the IASW electrostatic potential) propagating in the z direction, so that the dependence of all the perturbation quantities are in the form $f(x, z, t) \sim \tilde{f}(x)\exp[i(kz - \omega t)]$. For hot electrons satisfying $T_e \gg T_i, T_d$, we obtain

$$n_e = \frac{\epsilon_{i,d}}{4\pi e}\nabla^2\varphi \tag{5}$$

$$\boldsymbol{v}_e = i\frac{e}{m_e(\omega + i\nu_e)}\left[\nabla\varphi - r_{De}^2\nabla\left(\epsilon_{i,d}\nabla^2\varphi\right)\right] \tag{6}$$

where $\epsilon_{i,d} = 1 - \omega_{pi}^2/\omega(\omega + i\nu_i) - \omega_d^2/\omega(\omega + i\nu_d)$, ω_i and ω_d are the ion and dust plasma frequencies, ν_i and ν_d are the effective collision frequencies of the ions and dust particles, respectively. Substituting (5) and (6) into (3), one obtains for the IASW electrostatic potential

$$\left[\epsilon + \frac{V_{Te}^2}{\omega(\omega + i\nu_e)}\nabla^2\epsilon_{i,\,d}\right]\nabla^2\varphi = 0 \tag{7}$$

where $\epsilon = \epsilon_{i,\,d} - \omega_{pe}^2/\omega(\omega + i\nu_e)$ and ω_{pe} is the electron plasma frequency.

The differential equation (7) is of fourth order. In the plasma region ($x > 0$), the solution has the form

$$\varphi \sim A_1 \exp(-k_3 x) + A_2 \exp(-l_2 x) \tag{8}$$

where $l_2^2 = k_3^2 - \epsilon\omega(\omega + i\nu_e)/\epsilon_{i,\,d}V_{Te}^2$. In the dielectric region ($x < 0$), it has the form

$$\varphi \sim A_3 \exp(k_3 x) \tag{9}$$

which is to be matched to the solution (8) at the interface $x = 0$.

Three boundary conditions have to be satisfied [16]–[18]. Two of them are obtained directly from (7) by integrating it across the narrow interface region $0_- < x < 0_+$, yielding the conditions

$$\varphi|_{x=0_-} = \varphi|_{x=0_+} \quad \text{and } \epsilon_{i,\,d}\partial_x\varphi|_{x=0_-} = \epsilon_0\partial_x\varphi|_{x=0_+} \tag{10}$$

for the potential and its derivative at the plasma–dielectric interface. The third boundary condition depends on the physical situation considered. For ion-deposition setups such as those for PACVD, it can be in the form of a vanishing normal component of the electron fluid velocity at $x = 0$. This macroscopic model boundary condition has been widely used for bounded fluids and is valid to a good accuracy for well-polished dielectric or metal interfaces [16]–[19]. Accordingly, the boundary condition $v_{ex}(x = 0) = 0$ written in terms of the IASW electrostatic potential is

$$\left[d_x\left(1 - r_{De}^2\epsilon_{i,\,d}\nabla^2\right)\varphi\right]_{x=0} = 0 \tag{11}$$

which leads to $|A_1|/|A_2| = (l_2/k_3)(\omega/\omega_{pe})^2 \ll 1$. Thus, one can set $A_1 = 0$ and obtain from the conditions (10)

$$k_3^2 = \frac{\omega(\omega + i\nu_e)}{V_{Te}^2}\frac{\epsilon\epsilon_{i,\,d}}{\epsilon_{i,\,d}^2 - \epsilon_0^2} \tag{12}$$

which is the dispersion relation for acoustic type of surface waves in dusty plasmas. The corresponding damping coefficient δ ($= \operatorname{Im} k_3$) is given by

$$\frac{\delta}{\operatorname{Re} k_3} = \frac{\nu_i\omega_{pi}^2}{2\omega^3}\frac{\epsilon_0^2 + (\operatorname{Re}\epsilon_{i,\,d})^2}{(\operatorname{Re}\epsilon_{i,\,d})[\epsilon_0^2 - (\operatorname{Re}\epsilon_{i,\,d})^2]} \tag{13}$$

which shows that the damping (δ is negative definite) of IASW in a dusty plasma is determined by the ions.

In the absence of dust, and if the bounding medium is a vacuum ($\epsilon_0 = 1$), the expression (12) reduces to the well-known dispersion relation for IASW in a two-component dusty-free plasma [18], [20], [21]

$$k_3^2 = \frac{\epsilon_i}{1 + \epsilon_i}\left(\frac{\omega}{C_s}\right)^2 \tag{14}$$

where $\epsilon_i = 1 - \omega_{pi}^2/\omega^2$, and $C_s = (T_e/m_i)^{1/2}$ is the ion sound speed in a dusty-free plasma. The damping coefficient (13) can also be reduced to that of a collisional dust-free plasma [18], [20], [21].

Equation (12) is valid for all ion and dust surface waves satisfying $T_e \gg T_i,\ T_e$. In the usual ion-acoustic frequency range $\omega^2 \ll \omega_{pi}^2(1 + \epsilon_0^2)^{-1}$ and $\operatorname{Re}\epsilon_{i,\,d} \gg \epsilon_0$, (12) can easily be reduced to $k_3^2 = \omega^2/C_{sd}^2$, where $C_{sd} \equiv C_s g$ is the modified ion sound speed in the dusty plasma, and $g \equiv (1 + Z_{d0}^2 n_{d0}/n_{e0})^{1/2}$. Terms proportional to $m_i Z_d^2/m_d \ll 1$ have been neglected [3]. Thus, although the heavily charged dusts can significantly affect the charge balance, their motion does not contribute to the wave dynamics. Since for most dusty plasmas $g \geq 1$ [3], the dust grains can cause a significant increase of the ion sound speed and the IASW phase velocity, as has also been found for bulk ion-acoustic waves in infinite plasmas [5]. Clearly, the presence of dust also leads to a g-fold increase of the skin depth $\lambda_{\text{skin}} \approx 2\pi/l_2$, and a corresponding increase of the IASW field localization region. Physically, these increases are caused by a reduction of the electron density in the plasma because of the presence of the dust, which compensates the negative charge imbalance without being moved by the ion-acoustic wave field.

III. Variable-Charge Dust Grains

We now consider the effect of dust-charge variation on SW propagation. It is assumed that the dust charge varies according to the microscopic electron and ion currents flowing into the grain caused by the potential difference between the plasma and the grain surface. Thus, the grain charge variation can be on the ion time scale. On the other hand, as mentioned in Section II, the time scale of dust motion is much smaller ($m_i Z_d^2/m_d \ll 1$), and shall, thus, be neglected. The other assumptions are the same as those of the previous section, except that collisional damping (which can easily be introduced in the linear results later if the need arises) by the light particles is ignored. The particle temperatures satisfy $T_e \gg T_i, T_d$, but the ion temperature shall be kept finite in the dust-charging model. The process of dust charging is described by the charge conservation equation [2]–[4]

$$d_t q_d = I_e(q_d) + I_i(q_d) \tag{15}$$

where q_d is the average charge on the dust grain, and I_e and I_i are the microscopic electron and ion currents at the grain surface. We assume $q_d = q_{d0} + q_{d1}$, $I_e = I_{e0} + I_{e1}$, and $I_i = I_{i0} + I_{i1}$, where $q_{d0} = C(\varphi_g - \varphi_0)$ is the stationary dust charge, $C = 4\pi\epsilon_0 r(1 + r/r_{De})$ is the effective grain capacitance [22], and $\varphi_g - \varphi_0$ is the steady-state potential difference between the grain and the adjacent plasma. The steady-state electron and ion currents flowing into the grain are [2], [4]

$$I_{e0} = -\pi r^2 e(8T_e/\pi m_e)^{1/2} n_{e0} \exp[e(\varphi_g - \varphi_0)/T_e] \tag{16}$$

$$I_{i0} = -\pi r^2 e(8T_i/\pi m_i)^{1/2} n_{i0}[1 - e(\varphi_g - \varphi_0)/T_i] \tag{17}$$

which are equal and define the floating potential φ_0. Here, r is the average radius of the dust grains.

In the case considered, the phase velocity of the low-frequency surface waves is significantly smaller than V_{Ti} and V_{Te} [23], so that one can use fluid theory to calculate the perturbation ion grain current. For other phase speeds, more accurate expressions can be obtained from kinetic theory [2], [7]. Accordingly, the perturbation grain charge is governed by

$$d_t q_{d1} + \nu_{\rm ch} q_{d1} = -|I_{e0}| n_{e1}/n_{e0} + |I_{i0}| n_{i1}/n_{i0} \tag{18}$$

where $\nu_{\rm ch} = e|I_{e0}|[1/CT_e + 1/(CT_i - eq_{d0})]$ is the grain charging rate.

Considering again electrostatic wave perturbations, one can obtain from (1)–(6) the perturbed electron and ion densities

$$n_{e1} = (\epsilon_i / 4\pi e)\nabla^2\varphi - Z_{d1} n_{d0} \tag{19}$$

$$n_{i1} = -\left(Z_i e n_{i0}/m_i\omega^2\right)\nabla^2\varphi \tag{20}$$

where $\epsilon_i = 1 - \omega_{pi}^2/\omega^2$. Substituting the expressions (19) and (20) into (7), one can obtain the following dust charge relaxation equation:

$$d_t q_{d1} + \nu'_{\rm ch} q_{d1} = -\frac{|I_{e0}|\epsilon'_i}{4\pi e n_{e0}}\nabla^2\varphi \tag{21}$$

where $\nu'_{\rm ch} = \nu_{\rm ch} + \tilde{\nu}$, and $\tilde{\nu} \equiv |I_{e0}| n_{d0}/e n_{e0}$ is the correction to the dust-charging frequency arising from electron density perturbations. We have also defined $\beta = 1 - \omega_{pi}^2 f/\omega^2(1+f)$, where $f = Z_{d0} n_{d0}/n_{e0}$.

From (21) and (20), one can easily obtain

$$q_{d1} = -\frac{i}{4\pi e n_{e0}}\frac{|I_{e0}|\beta\nabla^2\varphi}{\omega + \nu'_{\rm ch}} \tag{22}$$

$$n_{e1} = \frac{\epsilon'_{i,d}}{4\pi e}\nabla^2\varphi \tag{23}$$

$$\boldsymbol{v}_e = \frac{ie}{m_e\omega}\left[\nabla\varphi - r_{De}^2\nabla\left(\epsilon'_{i,d}\nabla^2\varphi\right)\right] \tag{24}$$

where $\epsilon'_{i,d} = \epsilon_{i,d}[1 - i\tilde{\nu}\beta/\epsilon_i(\omega + \nu'_{\rm ch})]$. From the electron continuity equation, we obtain for the SW electrostatic potential

$$\left[\epsilon' + \left(V_{Te}^2/\omega^2\right)\nabla^2\epsilon'_{i,d}\right]\nabla^2\varphi = 0 \tag{25}$$

where $\epsilon' = \epsilon'_{i,d} - \omega_{pe}^2/\omega^2$.

Equation (25) has for the plasma- and dielectric-region solutions the same forms as in the last section, namely, (8) and (9), respectively. However, here $l_2^2 = k_3^2 - \epsilon'\omega^2/\epsilon'_{i,d}V_{Te}^2$. To match the solutions at the interface $x = 0$, we use boundary conditions analogous to the expressions (10) and (11). Accordingly, we obtain

$$k_3^2 = \frac{\omega^2}{V_{Te}^2}\frac{\epsilon'\epsilon'_{i,d}}{\epsilon'^2_{i,d} - \epsilon_0^2} \tag{26}$$

as the dispersion relation for IASW propagating in a plasma with variable-charge dusts. Equation (26) can easily be evaluated numerically for any given set of experimental parameters.

It is physically instructive to view (26) as the coupling equation for IASW and the dust-charging process. In fact, the equation can be simplified in the usual ion-acoustic wave limit Re $\epsilon'_{i,d} \gg \epsilon_0$. In this case, the dispersion relation becomes

$$\left(\omega^2 - k_3^2 C_{sd}^2\right)(\omega + i\nu'_{\rm ch}) = -i\alpha\nu'_{\rm ch}k_3^2 C_{sd}^2 \tag{27}$$

where $\alpha \equiv \tilde{\nu} f/\nu'_{\rm ch}(1+f)$ is the coupling coefficient relating the IASW ($\omega^2 \approx k_3^2 C_{sd}^2$) and the dust-charging mode ($\omega \approx -i\nu'_{\rm ch}$). Equation (27) is cubic and can be solved analytically. However, for a wide range of dusty plasmas, we have $\omega \gg \alpha\nu'_{\rm ch}$, which then leads to the solutions

$$\omega_{1,2} = \pm k_3 C_{sd}\left(1 - \frac{\alpha}{2}\frac{\nu'^2_{\rm ch}}{\nu'^2_{\rm ch} + k_3^2 C_{sd}^2}\right) - i\frac{\alpha}{2}\frac{\nu'_{\rm ch}k_3^2 C_{sd}^2}{\nu'^2_{\rm ch} + k_3^2 C_{sd}^2} \tag{28}$$

for the IASW, and

$$\omega_3 = -i\left(\nu_{\rm ch} + \tilde{\nu} - \frac{\alpha\nu'_{\rm ch}k_3^2 C_{sd}^2}{\nu'^2_{\rm ch} + k_3^2 C_{sd}^2}\right) \tag{29}$$

for the purely damped dust-charging mode. From (28) and (29), we see that charge variation slightly lowers (compared to the constant dust charge case) the IASW frequency and causes additional damping of the waves.

IV. Discussion

For typical laboratory conditions [13], $n_{d0} \sim 2 \times 10^5$ cm^{-3}, $n_{e0} \sim 10^9$ cm^{-3}, $T_e \sim 2$ eV, $T_i \sim 0.3$ eV, and $\omega/2\pi \sim 13.56$ MHz. We assume that the dust grains are SiO_2 particles with $r \sim 5$ μm, and model them as perfectly conducting spheres. Accordingly, we have $|Z_d|e = 4\pi\epsilon_0(\varphi_g - \varphi_0)r(1 + r/r_{De})$, where ϵ_0 is in SI units.

Assuming $\varphi_g - \varphi_0 \sim 10$ V and noting that $r \ll r_{De}$, we have approximately $|Z_d| = 7 \times 10^4$. Since the charge of the real (not perfectly conducting) dust grain would be somewhat less than that given by the latter expression, for numerical estimates we take $|Z_d| = 5 \times 10^4$ and obtain $g \sim 3.32$. One can thus see that the presence of dust grains in an RF plasma can significantly increase (nearly three times) the phase velocity of the IASW and a corresponding increase of the IASW field localization region. This results in an often-desirable increase of the volume of low-frequency RF-discharge produced and sustained plasmas. For RF plasmas containing smaller (1 μm or less) dust grains, the IASW phase velocity increase becomes less significant. Nevertheless, the specific frequency shift and damping of IASW can be used for diagnostics and control of the dust particles or impurity ions.

In obtaining (28) and (29), we assumed that the IASW frequency is sufficiently large such that $\omega \gg \alpha\nu'_{\rm ch} = \tilde{\nu} f/(1+f)$ is satisfied. Under typical laboratory conditions, $T_e \sim 10T_i$, $n_{i0} \sim 10^{12}$ cm^{-3}, $T_i \sim 1$ eV, $n_{e0} \sim 0.8n_{i0}$ and $r \sim 1$ μm, one obtains $\alpha\nu'_{\rm ch} \sim 4 \times 10^6$ s^{-1} with the frequency of the RF generator at 13.56 MHz. Thus, we have $|\omega_{1,2}| \approx 20\alpha\nu'_{\rm ch}$, which is realized over a wide range of the plasma parameters and gas-discharge operation frequencies. Here the SW skin depth and wavelength significantly exceeds the width of the sheath between the wall and the main body of the plasma. The density inhomogeneity and charge imbalance in the sheath region is then ignorable, and the homogeneous fluid edge plasma model valid [17], [23]. Quantitatively, the inequality is realized for $\omega \ll \omega_{pi}$. For RF frequencies of the order of tens of MHz the latter is satisfied for $n_e \approx$ O(10^9–10^{10}) cm $^{-3}$. However, for low-pressure (thus low density) discharges, the SW skin depth and wavelength can be of the same order as the sheath width. Also, sufficiently heavy dust grains can be localized near the bottom sheath region [11] due to gravity. A nonlocal kinetic treatment of the problem would then be necessary. Finally, we mention that in the numerical estimates we have considered surface-wave produced and sustained RF discharges with the SW frequency coinciding with that of the external generator. For other (e.g., HF) discharges, the effect of the applied RF field on the dust-charging process may also have to be taken into account.

It is worth mentioning that in the frequency range of interest, the IASW have equal normal and tangential (with respect to the interface) electric field components. This peculiarity allows [19] one to produce and drive ion beams in the field of the IASW and realize controlled ion bombardment of a substrate which is separated from the plasma boundary by a vacuum gap of width approximately the IASW skin depth. On the other hand, the presence of dust grains leads to a g-fold increase of the IASW skin depth and a decrease of both the normal and tangential components of the electric field at the plasma–vacuum interface. This implies that the normal component of the driving electric field becomes more homogeneous than that in the dust-free case. Furthermore, a decrease in the IASW electric field amplitude leads to a decrease of the energy of the ions falling onto the substrate. Thus, in the presence of dust, an additional extracting dc electric field may be necessary to restore the ion bombardment conditions for the dust-free case.

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