Physics of Plasmas

Charge-gradient instability of compressional dust lattice waves in electrorheological plasmas

Cite as: Phys. Plasmas **26**, 083701 (2019); https://doi.org/10.1063/1.5115346
Submitted: 17 June 2019 . Accepted: 16 July 2019 . Published Online: 16 August 2019

V. V. Yaroshenko

Physics of Plasmas ARTICLE scitation.org/journal/php

Charge-gradient instability of compressional dust lattice waves in electrorheological plasmas

Cite as: Phys. Plasmas **26**, 083701 (2019); doi: 10.1063/1.5115346
Submitted: 17 June 2019 · Accepted: 16 July 2019 ·
Published Online: 16 August 2019

View Online Export Citation CrossMark

V. V. Yaroshenko

AFFILIATIONS

Institut für Materialphysik im Weltraum, Deutsches Zentrum für Luft- und Raumfahrt (DLR), 82234 Wessling, Germany

ABSTRACT

It is shown that the longitudinal dust lattice mode in the one-dimensional string of microparticles oriented along the electric field (it highlights the major aspect of the electrorheological plasmas) can be subjected to a specific charge-gradient instability, associated with inhomogeneous dust charge distribution along the string due to the ion flow. This instability leads to the spontaneous excitation of compressional waves at elevated gas pressures with implications for the electrorheological plasmas studied in the Plasmakristall-4 facility in the ground based conditions and in microgravity on board the International Space Station. The obtained results can also be relevant for the stability analysis of the laboratory multilayer plasma crystals observed in the sheath of gas discharges.

I. INTRODUCTION

The most surprising discovery in the last few decades in complex plasma physics was the observation of crystal-like structures which spontaneously form when highly charged microparticles are trapped in a sheath electric field.[1–7] Such strongly coupled plasma agglomerates can support longitudinal and transverse vibrational modes, the so-called dust-lattice (DL) waves. The modes are the result of collective interactions of microparticles in a system with strong electrostatic coupling. The DL waves were extensively studied in many experiments using radio frequency and dc discharge plasmas in laboratory and under microgravity conditions.[4–14]

The first theoretical considerations of the DL modes were based on the convenient approximation of the one-dimensional particle string[8,9] that has been employed in solid state theory for the description of elastic vibrations. Later, the same approach was extended to 2D plasma crystals.[15] Most theoretical studies of the DL modes are usually based on the assumption that all dust particles carry an equal constant charge in the equilibrium state. However, contrary to usual electron-ion plasmas, where particles always have a fixed charge, the charge carried by microparticles is self-consistently related to the variations in surrounding plasma parameters and external conditions and thus may have a rather nonuniform distribution inside the experimental chamber. In addition, it should be borne in mind that in most complex plasma experiments, dust particles are subjected to discharge electric fields and associated ion flows, which introduce an asymmetry in dust-dust interactions in two ways. First, the ion flow around a charged microsphere leads to distributed ion density reduction upstream (where the flow lines diverge) and a more concentrated ion density enhancement downstream (where the flow lines converge). The screening cloud becomes asymmetric, and the associated plasma wake structure is responsible for the attractive forces[16–19] providing the formation of the particle chains along the electric field (the ion flow). The 1D particle strings are observed in many ground-based experiments in the sheath electric fields[4,6] and in microgravity conditions in the PK-4 facility (Fig. 3 in Ref. 21 and Fig. 2 in Ref. 22). In the PK-4 chamber (see Ref. 20 for a detailed description of the setup), the particle chains are formed along the axial discharge electric field in both regimes of dc and ac fields.[21,22] Complex plasma containing the particle chains is often referred to as electrorheological plasma.

On the other hand, the formation of plasma wakes and particle charging are self-consistently interconnected. Dedicated experiment[23] and numerical particle-in-cell (PIC) modeling of the dust charging in the presence of ion flows[24] indicate that charges on the downstream grain aligned with the flow are reduced due to the ion wake (i.e., the downstream microparticle becomes less negatively charged). For the ion Mach number $M_i \sim 1$, the charge reduction in the experiments achieves $\sim$22%, while the numerical simulations predict values up to $\sim$50%.[23] Notable reductions of the charges on the downstream particles have also been numerically found for a system of several grains in the subsonic ion flows ($M_i < 1$).[24]

The appearance of charge gradients seems to be obvious in the dc field but requires some explanation for the ac mode. In the ac regime, the polarity of the electric field is changed slowly ($\sim$milliseconds) on a time scale of the particle charging ($\sim$microseconds). Therefore, charges acquired by microparticles (organized in a string along the ion flow) at the first field polarity Q_i, $i=1, 2, 3,\ldots, N$, will be periodically replaced with Q_{N-i+1}, $i=1, 2, 3, \ldots, N$ at the second polarity. The time-averaged set of charges $\langle Q_i\rangle = (Q_i + Q_{N-i+1})/2$ represents in general a certain function of the particle position that is symmetric with respect to the middle of a string ($\sim N/2$) at the duty cycle 50%. Other duty cycles destroy this symmetry and can produce even more pronounced time-averaged dust charge distribution. Hence, when now studying multiparticle arrangements, it is logical to assume that complex ion dynamics modifies the charges carried by microparticles and existence of nonuniform dust charge distribution is generally a characteristic feature for the large scale dust structures in electrorheological plasmas.

Although some aspects of the dust charge gradients have been discussed for the DL mode-coupling instability,[25] it is a legitimate question to ask whether equilibrium dust charge distribution may affect the propagation of the compressional waves in the large scale particle chains observed experimentally e.g., in the PK-4 facility. In this paper, the step is taken of considering a simple model of weakly nonuniform electrorheological plasmas whose time-averaged dust charge profile can be approximated by linear charge gradients, for which the computations can be done explicitly.

II. AN EQUILIBRIUM STATE

We consider a one-dimensional (1D) particle chain oriented along the external electric field (z axis) providing the nonuniform dust charge distribution along the string due to the ion flows. The chain is composed of spherical particles of the same radius a (mass m_d), carrying negative charges $-Q_j = Q(z_j)$ with the corresponding interparticle distances $\Delta_{j,\ j\pm 1}$, $j=1, 2,\ldots N$. It might be wise to remember that for the electrorheological plasma, Q_j mean the time-averaged values as explained above. The dust potential governing the interaction between charged grains in electrorheological plasmas is still under debate (e.g., Ref. 22); so to focus solely on the dust charge effect, we consider the standard Yukawa potential

$$\varphi_j(z) = \frac{Q_j \exp\left(-|z-z_j|/\lambda_D\right)}{|z-z_j|}, \tag{1}$$

with the screening length λ_D provided by the plasma ions and electrons. The force acting on the nth particle due to the electric field of the jth grain is

$$\mathbf{F}_{n,j} = -Q_n\left(\partial\varphi_j/\partial\mathbf{z}\right)_{z=z_n}. \tag{2}$$

Employing the usual assumption that the interaction exists only between the nearest neighbors (usually the interparticle distance exceeds the screening length: $\Delta_{j,\ j\pm 1}/\lambda_D > 1$) leads to the momentum equation

$$\ddot{\mathbf{z}}_n + \nu_{dn}\dot{\mathbf{z}}_n = m_d^{-1}(\mathbf{F}_{n,n+1} + \mathbf{F}_{n,n-1}). \tag{3}$$

The dust–neutral collision frequency is represented by the standard Epstein theory[26] yielding $\nu_{dn} = \left(8\sqrt{2\pi}/3\right)a^2 n_n m_n v_{Tn}/m_d$, where n_n denotes the neutral gas density and m_n, T_n, and $v_{Tn} = \sqrt{T_n/m_n}$ are the mass, temperature, and thermal velocity of neutrals, respectively.

In the steady state ($\partial/\partial t=0$), the interaction force between the nodal nth particle and the $n+1$ grain has to be balanced by the corresponding force between the nth and ($n-1$)th particles, respectively, thus giving

$$\mathbf{F}^{(0)}_{n,n+1} + \mathbf{F}^{(0)}_{n,n-1} = -\frac{Q_nQ_{n+1}\exp\left(-\Delta_{n,n+1}/\lambda_D\right)}{\Delta^2_{n,n+1}}\left(1+\frac{\Delta_{n,n+1}}{\lambda_D}\right) + \frac{Q_nQ_{n-1}\exp\left(-\Delta_{n,n-1}/\lambda_D\right)}{\Delta^2_{n,n-1}}\left(1+\frac{\Delta_{n,n-1}}{\lambda_D}\right) = 0. \tag{4}$$

Although there have been a number of measurements and computer simulations of the inhomogeneous dust charge distributions,[24,27–31] all of them are related to specific experimental conditions, and the general behavior of $Q(z)$ is not defined. For our purposes, it is then instructive to invoke the simplest model of $Q(z)$, which can be easily discussed throughout and which allows us to ascertain the main physical features. We assume that the charges carried by symmetric particles can be approximated as

$$Q_{n\pm 1} \simeq Q_n \pm Q'_n\Delta_{n,n\pm 1}, \tag{5}$$

where $Q'_n = [\partial Q(z)/\partial z]_{z=z_n}$. Note that in the ac regime, the gradient term relates to the time-averaged dust charge profile, and Q'_n does not change its sign with the polarity switching. The characteristic scale of the dust charge variations $Q(z)/Q'_n \sim L_Q$ is of the order of the chain length, and hence, the approximation (5) follows logically from $L_Q \gg \Delta_{n,\ n\pm 1}$.

Introducing the average interparticle distance $\Delta = (\Delta_{n,n+1} + \Delta_{n,n-1})/2$ and spreading $\delta = (\Delta_{n,n+1} - \Delta_{n,n-1})/2$, we expand the balance equation (4) to the first order in terms of a small positive parameter $\varepsilon = (Q'_n\Delta/Q_n) = \Delta/L_Q \ll 1$ to get

$$\delta \simeq \frac{\varepsilon\Delta(\kappa+1)}{(2\kappa+\kappa^2+2)}, \tag{6}$$

with κ being a local lattice parameter $\kappa = \Delta/\lambda_D$. The equilibrium distances, between the nodal and adjacent particles, $\Delta_{n,\ n\pm 1}$, are then given by

$$\Delta_{n,n\pm 1} = \Delta\left[1 \pm \frac{\varepsilon\Delta(\kappa+1)}{(2\kappa+\kappa^2+2)}\right]. \tag{7}$$

III. DISPERSION RELATION FOR COMPRESSIONAL WAVES

For strongly coupled dust, it is reasonable to assume that the microparticles exhibit only small displacements around their equilibrium positions z_{n0}, i.e., $z_n = z_{n0} + \eta_n$, where ($\eta_n \ll \Delta$). Expanding the right-hand side of Eq. (3) around equilibrium and taking into account relation (7) as well as local variations in the particle charges $Q_j \to Q_j + Q'_j\eta_j$, ($j = n, n\pm 1$), one finds the linearized momentum equation

$$\ddot{\eta}_n + \nu_{dn}\dot{\eta}_n \simeq \Omega^2_{\parallel}(\eta_{n+1} + \eta_{n-1} - 2\eta_n) - \Omega^2_Q(\eta_{n+1} - \eta_{n-1}). \tag{8}$$

The longitudinal DL wave frequency slightly modified by the charge gradient term reads as

$$\Omega^2_{\parallel} = \Omega^2_0(2\kappa + \kappa^2 + 2 + \varepsilon\kappa), \tag{9}$$

with $\Omega_0^2 = Q_n^2 \exp(-\kappa)/(m_d\Delta^3)$. The frequency Ω_Q in the momentum equation (8) is solely specified by the equilibrium charge gradient through

$$\Omega_Q^2 = 2\varepsilon\Omega_0^2(\kappa+1). \quad (10)$$

For the DL perturbations obeying $kL_Q \gg 1$ (i.e., when the DL wavelengths $\sim k^{-1}$ are smaller than the characteristic scale of the dust charge distribution $\sim L_Q$), one can get the local dispersion equation to the first order in ε,

$$\omega^2 + i\nu_{dn}\omega = 4\Omega_{\parallel}^2 \sin^2(k\Delta/2) + 2i\Omega_Q^2 \sin(k\Delta). \quad (11)$$

For the uniform particle string ($\varepsilon = 0$), this gives the standard longitudinal DL wave damped due to the dust-neutral collisions. Such a mode has the wave number cutoff at $k \sim \nu_{dn}/(2\Delta\Omega_{\parallel})$ [see the plot at $\varepsilon = 0$ in Fig. 1(a)]. In the case $\varepsilon \neq 0$, solving the dispersion relation (11) in terms of the real (ω_r) and imaginary (ω_i) parts of the wave frequency, yields

$$\omega_r = \left(\sqrt{2}/4\right)\Big[(32\Omega_Q^4 \sin^2 q + (16\Omega_{\parallel}^2 \sin^2(q/2) - \nu_{dn}^2)^2)^{1/2} + 16\Omega_{\parallel}^2 \sin^2(q/2) - \nu_{dn}^2\Big]^{1/2}, \quad (12)$$

$$\omega_i = (\Omega_Q^2 \sin q)/\omega_r - \nu_{dn}/2, \quad (13)$$

where q denotes the normalized wave number $q = k\Delta$.

To analyze the solutions for typical experimental conditions, we consider the complex plasma parameters relevant for the Plasmakristall-4 (PK-4) facility, where the particle strings are often observed in microgravity conditions on board the International Space Station. The dc discharge is operated in neon/argon gas at various neutral gas pressures, mostly in a range $p_n \sim 20$–60 Pa. In Table I, we give some basic complex plasma parameters estimated at the dc current of 0.5 (mA) in neon gas.[20] The equilibrium dust charge numbers $Z_d = Q_n/e$ are given according to recent modeling of the charging of individual microparticles with radii $a \simeq 1.25\,\mu$m and $a \simeq 1.7\,\mu$m in the appropriate discharge conditions of the PK-4 chamber.[32] We chose the average interparticle distances $\Delta \sim 100\,\mu$m ($a \sim 1.25\,\mu$m) and $\Delta \sim 150\,\mu$m ($a \sim 1.7\,\mu$m) close to the typical experimental values.

The resulting real and imaginary parts of the wave frequency ω as a function of the wave number $q = k\Delta$ at various dust charge gradients are shown in Figs. 1(a) and 1(b) by solid and dashed curves, respectively. The plots are calculated for $a \sim 1.7\mu$ at a gas pressure of $p_n \simeq 40$ Pa. Contrary to the standard DL mode ($\varepsilon = 0$), which has no solution at the range of small wave numbers $q \lesssim \nu_{dn}/(2\Omega_{\parallel})$, the dust charge gradient provides nonzero ω_r at these q values. Here, the behavior of ω_r is specified by a value of ε. For larger wave numbers $q \sim 1$, the real part of the wave frequency manifests the acousticlike dispersion, close to the standard DL solution at $\varepsilon = 0$. The imaginary part of the wave frequency ω_i even more crucially depends on the actual value of the dust charge gradient: small variations of ε can provide the excitation of the compressional waves in the range of small q in the collisional plasma, which are evanescent in a uniform string. As is illustrated in Fig. 1(b), there is a relation between the dust charge gradient and most unstable wave numbers, and larger ε leads to the excitation of the DL waves with larger q (smaller wavelengths). Instability of the compressional mode becomes physically possible at the expense of the external energy which generates an equilibrium charge distribution within the dust chain.

Using the data of Table I, we employ as a first approximation the linear interpolations for the equilibrium plasma density and dust

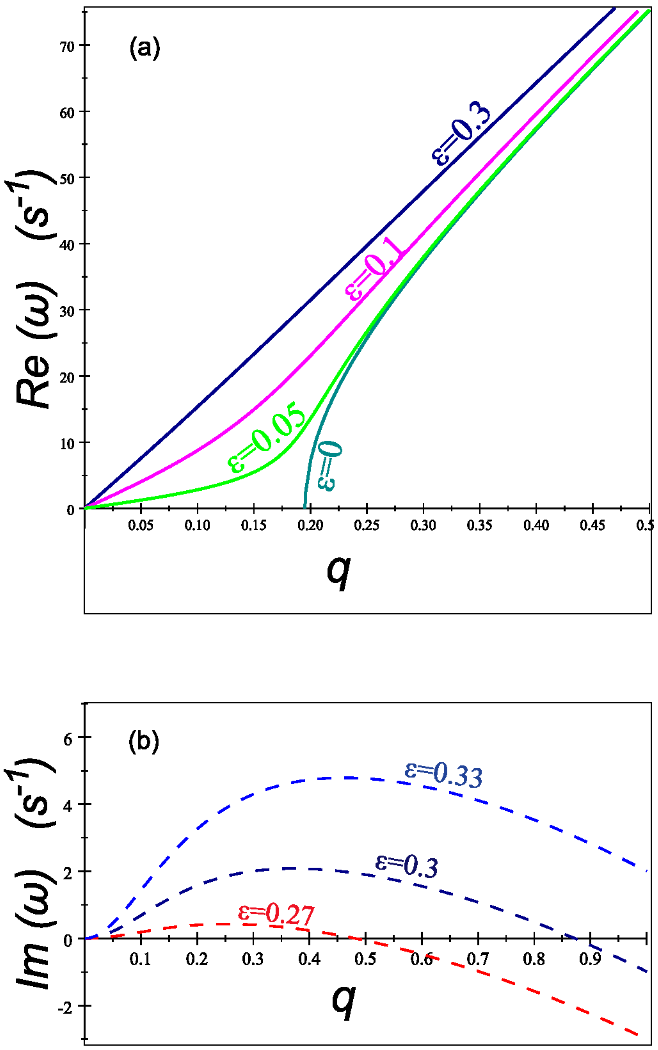

FIG. 1. Real (a) and imaginary (b) parts of the frequency of compressional DL mode vs wave number $q = k\Delta$ for various dust charge gradients ε. Calculations have been done for particles $a = 1.7\,\mu$m, interparticle distance $\Delta = 150\,\mu$m, and neon gas pressure $p_n \simeq 40$ Pa. Other parameters correspond to the respective data of Table I. In the local approximation, the wavenumbers q are limited by ε, and the obtained solutions are valid for $q > \varepsilon$.

TABLE I. Characteristics of complex plasmas in the PK-4 setup.[a]

Pressure, p_n, Pa	20	40	60
Plasma density, n_{0i}, 10^8 cm^{-3}	0.9	1.4	1.5
Electron temperature, T_e, eV	9.8	8.7	8.4
Particle charge number ($a = 1.25\,\mu$m), Z_d, 10^3	2.4	1.6	1.5
Particle charge number ($a = 1.7\,\mu$m), Z_d, 10^3	3.8	2.3	2.1

[a]Numerical data are adopted from Refs. 20 and 32.

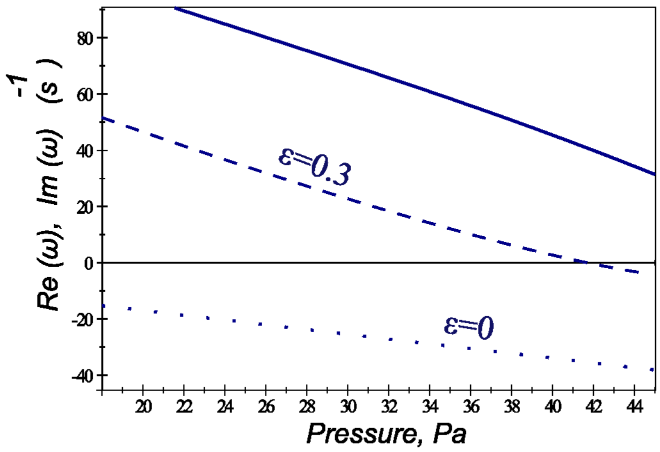

FIG. 2. Real (solid curve) and imaginary (dashed curve) parts of the wave frequency as a function of gas pressure at $q = 0.35$ and $\varepsilon = 0.3$. Other parameters as in Fig. 1. The dotted line displays the damping rate of the standard longitudinal DL wave in the uniform particle chain.

charge in the considered range of the gas pressure p_n, so that we actually deal with the pressure dependent equilibrium quantities $n_i(p_n)$ [or screening length $\lambda_{Di}(p_n)$] and $Z_d(p_n)$. Inserting these into the solutions (12) and (13) yields ω_r and ω_i as functions of the gas pressure, at a fixed wave number. The upper (solid) curve in Fig. 2 illustrates a decrease in the real part of the wave frequency with p_n calculated at $q = 0.35$, and the lower (dashed and dotted) curve represents the respective ω_i describing either the DL mode instability ($\omega_i > 0$ for $\varepsilon = 0.3$, $p_n \lesssim 40$ Pa) or wave damping ($\omega_i < 0$). The example in Fig. 2 yields excitation of the longitudinal DL at elevated gas pressures for rather small dust charge gradients.

Finally, Fig. 3 indicates the instability domains obtained in the frame of local dispersion relation in the parameter space $\{p_n, \varepsilon\}$ for various particle sizes $a \sim 1.25\,\mu$m, $a \sim 1.7\,\mu$m, and $a \sim 3.4\,\mu$m. The

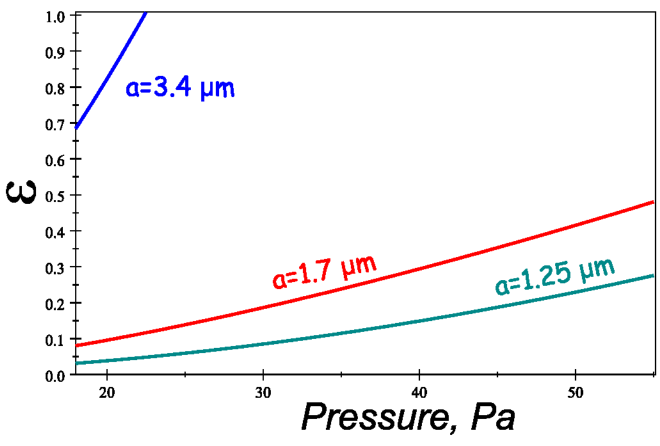

FIG. 3. Instability parameter space in terms of normalized charge gradient ε and neon gas pressure p_n for various microparticle sizes and interparticle distances $\Delta = 100$, $150\,\mu$m and $\Delta = 250\,\mu$m, respectively. The pressure dependent equilibrium dust charges and plasma densities correspond to linear interpolations of Table I. For grains with $a \simeq 3.4\,\mu$m, two times larger values of Z_d than the respective values for $a \sim 1.7\,\mu$m are assumed.

curves express the condition $\omega_i = 0$, and quantities p_n and ε lying above the curves lead to the dust charge gradient instability. In other words, any given p_n implies a threshold value ε providing the excitation of compressional DL waves in a particle chain. It is interesting that the margin of the charge gradient instability is practically independent of a value of the wave number in the long wavelength range ($q \lesssim 1$). For grains with $a \sim 1.25\,\mu$m, the DL mode instability develops already at small charge gradients ($\varepsilon \lesssim 0.1$) and gas pressures $p_n \lesssim 35$ Pa. The chains containing large microparticles ($a \sim 3.4\,\mu$m) appear to be more stable with respect to the considered instability: the excitation of the DL mode requires essentially large dust charge nonuniformity like $\varepsilon \gtrsim 0.5$ and lower neutral gas pressures $p_n \lesssim 15$ Pa.

IV. CONCLUSIONS

The standard DL wave model was suitably modified compared to earlier work to include a weakly nonuniform dust charge distribution along the particle string. The latter is assumed to be caused by discharge electric fields and the associated ion wakes affecting the charging of grains in the multiparticle agglomerations. To deal with the nonuniform string, a simple dust charge gradient model was employed. First, the local dispersion relation was obtained, from which it was concluded that even small dust charge variations significantly modify the compressional DL mode propagation.

For reasonable dust charge gradients, like those predicted by simulations ($\varepsilon \lesssim 0.5$[24]), the difference can be quite appreciable, and a nonuniform particle chain could exhibit DL waves at small wave number-low-frequency domain at high gas pressures. While the compressional mode in the standard uniform chain is strongly damped, the wave propagating in the direction of the charge gradient can exhibit a specific instability, the criterion for which is dependent on the spatial distribution of the microparticle charges along the string and neutral gas pressure (Fig. 3). The changes, therefore, are qualitative rather than quantitative in the sense that the standard DL mode treatment of the uniform chain admits only evanescent wave solutions (Fig. 2). Invoking a more complicated model of the dust-dust interaction than that described by the Yukawa potential does not change the physics of the considered instability but can change the instability criterion quantitatively.

The results presented here can be directly applicable to stability analysis of the particle strings (electrorheological plasma) observed in the PK-4 facility and in the sheath of gas discharges. Our treatment indicates that, at given gas pressure, the instability threshold in terms of ε is increased with the microparticle size, interparticle distance, and discharge current (plasma density). Furthermore, using argon gas also admits a higher threshold value of ε for the development of the DL wave instability under similar plasma conditions than would be observed in neon discharge. In view of this, we predict that stable electrorheological plasmas (with respect to the charge gradient instability) require the use of large microparticles, higher discharge current, and argon gas. Future studies of the large scale strings in the PK-4 chamber under various experimental conditions can verify and quantify the predicted effects.

ACKNOWLEDGMENTS

The author thanks S. Khrapak, M. Thoma, M. Schwabe, and M. Pustylnik for discussions of the theoretical results in the light of observations in the PK-4 facility.

REFERENCES

[1] H. Ikezi, Phys. Fluids **29**, 1764 (1986).
[2] J. H. Chu and I. Lin, Phys. Rev. Lett. **72**, 4009 (1994).
[3] J. H. Chu, J. B. Du, and I. Lin, J. Phys. D **27**, 296 (1994).
[4] H. Thomas, G. E. Morfill, V. Demmel, J. Goree, B. Feuerbacher, and D. Möhlmann, Phys. Rev. Lett. **73**, 652 (1994).
[5] Y. Hayashi and K. Tachibana, Jpn. J. Appl. Phys., Part 2 **33**, L804 (1994).
[6] A. Melzer, T. Trottenberg, and A. Piel, Phys. Lett. A **191**, 301 (1994).
[7] J. B. Pieper and J. Goree, Phys. Rev. Lett. **77**, 3137 (1996).
[8] G. E. Morfill, H. M. Thomas, and M. Zuzic, in *Advances in Dusty Plasmas*, edited by P. K. Shukla, D. A. Mendis, and T. Desai (World Scientific Singapore, 1997), pp. 99–142.
[9] F. Melandsø, Phys. Plasmas **3**, 3890 (1996).
[10] A. Homann, A. Melzer, S. Peters, and A. Piel, Phys. Rev. E **56**, 7138 (1997).
[11] S. V. Vladimirov, P. V. Shevchenko, and N. F. Cramer, Phys. Rev. E **56**, R74 (1997).
[12] S. V. Vladimirov, N. F. Cramer, and P. V. Shevchenko, Phys. Rev. E **60**, 7369 (1999).
[13] G. E. Morfill, A. V. Ivlev, and J. R. Jokipii, Phys. Rev. Lett. **83**, 971 (1999).
[14] A. Nefedov, G. E. Morfill, V. E. Fortov, H. M. Thomas, H. Rothermel, T. Hagl, A. V. Ivlev, M. Zuzic, B. A. Klumov, A. M. Lipaev, V. I. Molotkov, O. F. Petrov, Y. P. Gidzenko, S. K. Krikalev, W. Shepherd, A. I. Ivanov, M. Roth, H. Binnenbruck, J. A. Goree, and Y. P. Semenov, New J. Phys. **5**, 33 (2003).
[15] A. Homann, A. Melzer, S. Peters, R. Madani, and A. Piel, Phys. Lett. A **242**, 173 (1998).
[16] V. A. Schweigert, I. V. Schweigert, A. Melzer, A. Homann, and A. Piel, Phys. Rev. E **54**, 4155 (1996).
[17] A. Melzer, V. A. Schweigert, and A. Piel, Phys. Rev. Lett. **83**, 3194 (1999).
[18] F. Melandso and J. Goree, Phys. Rev. E **52**, 5312 (1995).
[19] V. V. Yaroshenko, H. M. Thomas, and G. E. Morfill, New J. Phys. **8**, 54 (2006).
[20] M. Y. Pustylnik, M. A. Fink, V. Nosenko, T. Antonova, T. Hagl, H. M. Thomas, A. V. Zobnin, A. M. Lipaev, A. D. Usachev, V. I. Molotkov, O. F. Petrov, V. E. Fortov, C. Rau, C. Deysenroth, S. Albrecht, M. Kretschmer, M. H. Thoma, G. E. Morfill, R. Seurig, A. Stettner, V. A. Alyamovskaya, A. Orr, E. Kufner, E. G. Lavrenko, G. I. Padalka, E. O. Serova, A. M. Samokutyayev, and S. Christoforetti, Rev. Sci. Instrum. **87**, 093505 (2016).
[21] A. V. Ivlev, G. E. Morfill, H. M. Thomas, C. Räth, G. Joyce, P. Huber, R. Kompaneets, V. E. Fortov, A. M. Lipaev, V. I. Molotkov, T. Reiter, M. Turin, and P. Vinogradov, Phys. Rev. Lett. **100**, 095003 (2008).
[22] A. V. Ivlev, M. H. Thoma, C. Räth, G. Joyce, and G. E. Morfill, Phys. Rev. Lett. **106**, 155001 (2011).
[23] J. Carstensen, F. Greiner, D. Block, J. Schablinski, W. J. Miloch, and A. Piel, Phys. Plasmas **19**, 033702 (2012).
[24] D. Block and W. J. Miloch, Plasma Phys. Controlled Fusion **57**, 014019 (2015).
[25] V. V. Yaroshenko, A. V. Ivlev, H. M. Thomas, and G. E. Morfill, Phys. Rev. E **71**, 046405 (2005).
[26] P. Epstein, Phys. Rev. **23**, 710 (1924).
[27] V. V. Yaroshenko, T. Antonova, H. M. Thomas, and G. E. Morfill, Phys. Plasmas **16**, 103505 (2009).
[28] A. Samaryan, A. Chernyshev, O. Petrov, A. Nefedov, and V. E. Fortov, J. Exp. Theor. Phys. **92**, 454 (2001).
[29] A. Ivlev, U. Konopka, and G. E. Morfill, Phys. Rev. E **62**, 2739 (2000).
[30] M. R. Akdim and W. J. Goedheer, Phys. Rev. E **65**, 015401(R) (2001).
[31] B. A. Klumov, A. V. Ivlev, and G. E. Morfill, in *Dusty Plasmas in the New Millennium*, edited by R. Bharuthram, M. A. Hellberg, P. K. Shukla, and F. Verheest (AIP, New York, 2002), p. 349.
[32] T. Antonova, S. A. Khrapak, M. Y. Pustylnik, M. Rubin-Zuzic, H. Thomas, A. M. Lipaev, A. D. Usachev, V. I. Molotkov, and M. H. Thoma, "Particle charge in PK-4 dc discharge from ground-based and microgravity experiments," Phys. Plasmas (unpublished).