Long-range attraction between particles in dusty plasma and partial surface tension of a dusty phase boundary

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Effective potential of a charged dusty particle moving in homogeneous plasma has a negative part that provides attraction between similarly charged dusty particles. A depth of this potential well is great enough to ensure both stability of crystal structure of dusty plasma and sizable value of surface tension of a boundary surface of dusty region. The latter depends on the orientation of the surface relative to the ion flow, namely, it is maximal and positive for the surface normal to the flow and minimal and negative for the surface along the flow. For the most cases of dusty plasma in a gas discharge, a value of the first of them is more than sufficient to ensure stability of lenticular dusty phase void oriented across the counter-ion flow.

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I. INTRODUCTION

A gas-solid phase transition observed in different dusty laboratory plasmas [1–8]. It counts in favor of presence of strong long-range attraction between similarly charged dusty particles. By now some physical mechanisms are proposed to explain formation of a regular arrangement of micron-sized particles embedded in a gas discharged plasma. A part of them is based on the account for electrostatic fields of strata and walls of discharge tubes [9]. In doing so, a generality of solution of the problem of the effective attractive potential is lost and, in particular, the situation of dusty plasma crystal in a thermal dusty plasma of gas burner [5] where there are no walls and external fields drops out of such interpretations.

The most promising approach holds that the ion streaming motion causes an attractive wake potential behind the dust particles. Originally it was developed [10,11] for particular case of supersonic flows which is realized in the sheath of radiofrequency discharges. Later, it was extended with participation of one of the authors of Ref. [10] to the case of subsonic ion flows [12], but physics of shielding of dust charged particles was supposed to be strongly modified with mandatory regard to anisotropy and asymmetry of the ion temperature in the sheath.

Below is shown that the attraction between similarly charged particles can be resulted from a dynamical screening of the Coulomb potential remaining in the frame of a single physical mechanism for both supersonic and subsonic regimes. The distinction of the dynamical screening from the static Debye screening is due to a motion of particles relative to screening charges (counter-ions or electrons). From the physical point of view, this effect can be interpreted as a consequence of loss of spherical symmetry of the Debye screening cloud around the moving charged particle resulted in a space charge with the opposite sign forming in its wake. Not only do this space charge compensates the particle's potential but may also give rise to a local wake potential of the opposite sign. In the supersonic regime the wake potential oscillates in as much of Cherenkov wave generation.

This effect is well known in the usual electron-ion plasma [13–15]. When the charge is moving, the static Debye screening modifies so that the potential in the oncoming flow of electrons grows approaching to the Coulomb potential with an increase of velocity, while the potential in the outward flow decreases up to alternating of its sign.

Similar effect was found [15,16] in a system of gravitating masses where the static screening is absent and accounting for the dynamical screening gives rise to an alternating potential of gravitational interaction.

The aim of the present paper is to propose a general model of dynamical screening of field of dusty particle charge to explain the observable interparticle attraction. It may be supposed that the distance from the particle to the attractive minimum of the wake potential determines a period of the resulted lattice of dusty particles.

Because of an anisotropy of the wake potential, the resulted lattice likewise is to be anisotropic. There are many experimental evidences of such anisotropy down to formatting of cylindrical structures [7,8].

On the other side, an appearence of crystal structures without preferred directions of crystal axes says in favor of an alternative mechanism of the isotropic attractive interaction between dusty particles. A model of this kind for an equilibrium dusty plasma [17] can be mentioned here. It is possible that this model of attraction can give rise to forming of crystal structure, but it is irrelevant to the observed [1–4] dependence of the lattice period on the ion flow velocity and forming of hierarchies along the flow.

In general, different mechanisms of attractive interaction are possible in such complex system as the dusty plasma but here only one of them, namely, the dynamical screening effect will be considered.

Besides, the interparticle attraction is to give rise to a surface tension of a dusty phase interface. Really, this explains a sharp nondiffusion character of a particle density variation observed at the interfaces in the experiments on dusty plasma crystals. The most characteristic phenomenon of such kind are voids in a homogeneous dusty plasma. The anisotropy of the wake potential gives rise to a strong dependence of the surface tension coefficient on the interface ori-

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entation relative to the ion flow. In particular, the interfaces oriented along the ion flow possesses the negative surface tension and thus canot exist. This property explains lenticular and pancake-shaped forms of the observable voids [18–22].

The notion of the surface tension of a dusty phase was mentioned briefly by Tsitovich [23] but his estimation of its value was erroneous. Conceivably that might be the reason why the surface tension was not mentioned in the subsequent papers by Tsitovich (as a co-author) on the theory [24,25] of spherical voids.

The outline of the paper is the following. In Sec. II, the model for effective nonpotential attractive forces of interactions between dust particles will be developed. This model can explain an appearance of the lattice structure both in gas discharge dusty plasma and in thermal dusty plasma of gas burner. In Sec. III, a surface tension of a boundary of dusty phase is estimated and conditions of stability of a lenticular void in the dusty plasma are examined.

II. EFFECTIVE INTERPARTICLE INTERACTION

The system under consideration consists of negatively charged dusty particles of the charge $Q=-Z_p e$ and concentration n_p, and positively charged counter-ions of the charge e (for definiteness sake we will consider singly charged ions) and mass m. For simplicity, neutral molecules of a buffer gas are neglected here and dusty particles are considered as point charges.

While forming of the negative charges of dusty particles the electrons are condensed on particles' surfaces and their concentration in particles' vicinities of the Debye ionic screening radius $\lambda_i=[k_B T_i/(4\pi e^2 n_i)]^{1/2}$ vanishes. As a result, when calculating the dynamical screening of the charged particle by ions, we can restrict our consideration to the simplified model of one-component ion plasma in which the negatively charged dusty particles are immersed.

If the negative test point charge Q is moving in the system of ions and dusty particles with the velocity $\mathbf{u}$ it gives rise to some perturbations of the system state. Due to great difference in masses and concentrations of ions and dusty particles we can restrict our consideration to a perturbation of the ion component only (just as a perturbation of an electron component is taken into account only, as a rule, in the problem of screening of an ion charge in the electron-ion plasma). Such perturbation of ionic subsystem is described by the set of the linear Vlasov equation for a distribution function over coordinates and velocities of ions and Poisson equation for a potential induced by the perturbation of ion density and moving test charge. Interactions between ions and neutral molecules of the buffer gas are neglected here [26].

Such simplified model is well studied in the test-particle approach to the electron-ion plasma theory [13–15], so we can omit intermediate calculations and write down at once the resulted form of the effective potential of the moving test charge,

$$\Phi(\mathbf{r},t)=\frac{Q}{2\pi^2}\int d^3k\frac{\exp\{\mathbf{k}\cdot(\mathbf{r}-\mathbf{u}t)\}}{k^2\varepsilon(k,\mathbf{k}\cdot\mathbf{u})}, \tag{1}$$

where

$$\varepsilon(k,\mathbf{k}\cdot\mathbf{u})=1+\frac{\kappa^2}{k^2}W\left(\frac{\mathbf{k}\cdot\mathbf{u}}{k\tilde{v}}\right) \tag{2}$$

is the dynamical permittivity of the ion subsystem. Here $\kappa=[4\pi e^2 n_i/(k_B T_i)]^{1/2}$ is the ionic Debye wave number, $\tilde{v}=(2k_B T_i/m)^{1/2}$ is the mean heat velocity of the ions and

$$W(t)=1-\sqrt{\pi}te^{-t^2}\ \mathrm{erfi}(t)+i\sqrt{\pi}te^{-t^2}, \tag{3}$$

where $\mathrm{erfi}(t)$ is the imaginary error function. Let us choose an axis z along $\mathbf{u}$ and introduce dimensionless variables $Z=(z-ut)\kappa$, $X=x\kappa$, $\mathbf{K}=\mathbf{k}/\kappa$ and the Mach number $M=u/\tilde{v}$ relative to the ion heat velocity. Then, in the test particle's frame accounting for the cylindrical symmetry, we get

$$\Phi(\mathbf{r})=\frac{Q\kappa}{2\pi^2}\int_{-\pi}^{\pi}d\phi\int_0^{\pi}d\theta\int dKK^2\sin\theta\frac{\exp\{iK\Delta\}}{K^2+W(M\cos\theta)}$$
$$=Q\kappa\varphi(X,Z;M), \tag{4}$$

where

$$\varphi(X,Z;M)=\frac{1}{(X^2+Z^2)^{1/2}}-I(X,Z;M),$$

$$\Delta=X\sin\theta\cos\phi+Z\cos\theta. \tag{5}$$

The triple integral $I(X,Z;M)$ determines the departure of the effective potential from the Coulomb potential to which the first member of $\varphi(X,Z;M)$ corresponds. When the Cauchy integral over K is taken it becomes the double integral

$$I(X,Z;M)=\frac{1}{4\pi^2}\int_{-\pi}^{\pi}d\phi\int_0^{\pi}d\theta$$
$$\times\sin\theta\sqrt{W}[2\,i\sinh(\Delta\sqrt{-W})\mathrm{Ci}(-i\Delta\sqrt{W})$$
$$+\cosh(\Delta\sqrt{W})(\pi+2\,i\,\mathrm{Shi}(\Delta\sqrt{W}))], \tag{6}$$

where $\mathrm{Shi}(t)$ is the hyperbolic sine integral function. It is not difficult to calculate numerically the value of this integral in an arbitrary range of the parameter M variations at any point X, Z. Detailed discussion of the similar numerical calculations can be found elsewhere (see, e.g., Ref. [13,14]).

Corresponding results for $M=3$ are illustrated in Fig. 1. It is seen here that there is well defined negative minimum of the potential along the axis Z at a distance of order of the 5 D radius, which has been formed as a result of deformation of screening Debye cloud in the vicinity of the moving particle.

In gas discharges, where forming of crystal structure of dusty particles was observed, a diffusion ion velocity $\mathbf{u}$ is determined by their mobility and an electric field strength. To define the mobility, we are to take into account a tendency of ions to unite with molecules and atoms into complexes of the types N_4^+, O_4^+, Ne_2^+, He_2^+, and recharge effect also. According to Ref. [28], the diffusion velocity of the complex Ne_2^+ is of order of 50 m c^{-1} in a characteristic for glow-discharge field $E/p=1$ V cm^{-1}torr^{-1} and the heat velocity

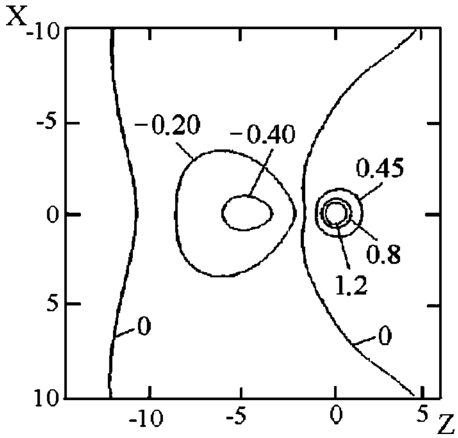

FIG. 1. The reduced potential $\varphi=\Phi/(Q\kappa)$ of the negatively charged dusty particle at $Z=0$ moving relative ions along Z axis with the reduced velocity $M=u/\tilde{v}=3$.

$\tilde{v}=400$ m c^{-1} (at $T=300$ K). Then $M=0.125$. For the ion Ne^{+}, the diffusion velocity decreases near-threefold as a consequence of the resonant recharge, then $M\simeq0.05$. In the case of strong field, the ion drift velocity may not only approach to the heat velocity but surpass it [28]. Because of this, a wide range of values of the parameter $M=u/\tilde{v}$ will be considered below (asymptotic case $M\gg1$ was discussed in Refs. [10,11]).

Observed quasicrystal structure in a thermal plasma, that is, in the dusty flame of gas burner [5,8,29], can be treated also as a consequence of movement of charged particles relative to ions of the flame. To obtain quantitative estimations of the effect, the additional data on a gap of velocities of particle and ion flows are necessary. This gap is determined by conditions of injection of the dusty particles into the flame, and in the steady state limit it is likely to go to a sedimentation velocity of charged particles in the rising flow.

The graphs of $\varphi(X=0,Z<0;M)$ for different M in the wake of charged particle are illustrated in Figs. 2 and 3. As can be seen, the depth of the negative minimum $|\varphi_{min}|$ increases with an increasing of M from 0.01 to 1 (see Fig. 4), and its location Z_{min} shifts to the particle. At low M $|Z_{min}|$ depends on M as $|Z_{min}|=1.6-1.7\ln M$ (see Fig. 5). When further increasing of M, the potential becomes oscillating and

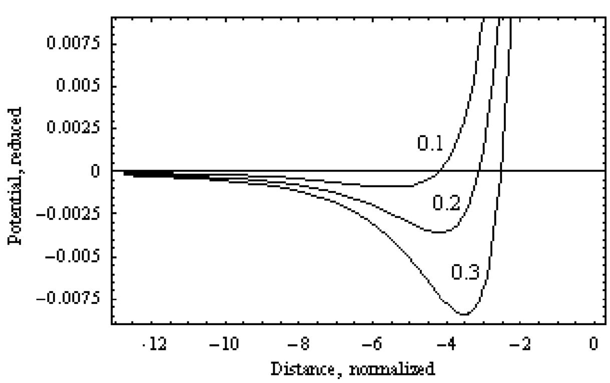

FIG. 2. The same as in Fig. 1 along the axis Z at $M=0.1$, 0.2, 0.3, correspondingly.

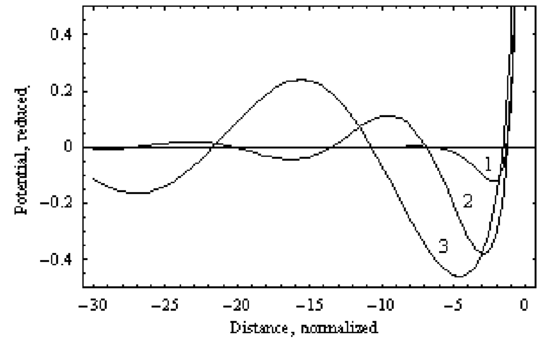

FIG. 3. The same as in Fig. 1 along the axis Z at $M=1$, 2, 3, correspondingly.

its first minimum moves away from the particle. For $M>2$ the graph in Fig. 5 can be well approximated by the expression $|Z_{min}|=1.8\sqrt{M^2-1}$, which is close to the experimental data [1–4] on the period L of a crystal lattice in a dusty plasma if it is identified as $L\simeq|Z_{min}|$. It should be noted here that experimental detection of the minimum of $L(M)$ corresponding to the minimum of the graph in Fig. 5 would be a good evidence in favor of the approach under discussion.

The values $M>1$ correspond to a supersonic movement when Cherenkov radiation of ion-acoustic waves takes a place that explains an appearance of space oscillations of the wake potential. It should be noted that in the range of values of M from 1 to ~3.3 an increasing of the depth of the first minimum takes a place but at further gain of M the depth of the first minimum decreases (see Fig. 4).

In the incoming flow ($Z>0$) the effective potential varies continuously with a gain of M from the Debye potential $\sim\exp\{-Z\}/z$ to Coulomb one $\sim1/z$. This indicates that the forward part of the Debye screening cloud is not pressed by the incoming flow to the particle but blew away from it.

The negative minimum of the wake potential points to presence of a space charge of the opposite sign. Thus, the large charged particles moving relative ion subsystem together attendant space charges constitute a system of dipoles oriented along the direction of the relative motion. This gives rise to formation of hierarchies of particles-dipoles along lines of the ion flow.

Moreover, this interaction between dust particles is asymmetric in such a way that attractive force is communicated only downstream the ion flow. This situation was clearly demonstrated in the experiments [30,31] where upper or

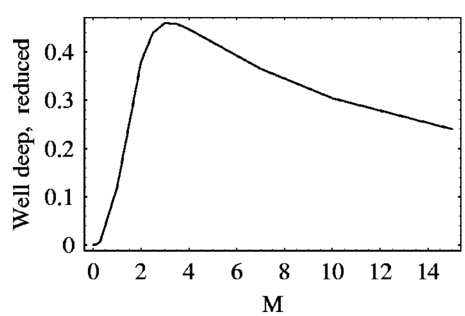

FIG. 4. Dependence of the depth of negative minimum of the reduced potential on M.

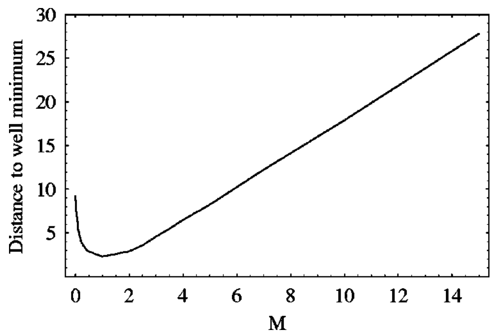

FIG. 5. Dependence of the position of negative minimum of the reduced potential on M.

lower particle was pushed by the laser beam and it was just lower dust particle that fitted its position when the upper particle was shifted and not vice versa. So heavily nonconservative character of the interparticle interactions excludes any possibility to introduce an effective potential of the interaction.

Besides, as it is seen in Fig. 1, the resulted negative potential is long-range in the transverse direction X, also, that provides an attraction between particles from neighbor hierarchies and, in its turn, causes the particles from neighbor hierarchies to shift relative one another along the ion flow for a distance of order of a half dipole length. Because of this interaction is much weaker then along the flow, the situations can take place when the lengthwise attraction appears quite strong for the hierarchies formation but too weak for transverse ordering [8].

Thus, the movement of the charged particle relative ion flow gives rise to anisotropic potential field in which an energy of another similar particle of the charge Q is $U(\mathbf{r}) = Q\Phi(\mathbf{r}) = Q^2\kappa\varphi(\mathbf{r})$. Then, the effective force of interaction between the particles is

$$\mathbf{F}(\mathbf{r}) = -\frac{\partial U(\mathbf{r})}{\partial \mathbf{r}} = -Q^2\kappa\frac{\partial \varphi(\mathbf{r})}{\partial \mathbf{r}}. \qquad (7)$$

For the formation of the crystal structure, it is necessary that the maximum deep of the attraction energy $|U_{min}| = |Q\Phi_{min}| = Q^2\kappa|\varphi_{min}|$ surpass the heat motion energy, that is,

$$\frac{|U_{min}|}{k_B T_p} > 1 \quad \text{or} \quad |\varphi_{min}| > \frac{k_B T_p}{Q^2\kappa}, \qquad (8)$$

where T_p is the kinetic temperature of the particle. According to both experimental data [32] and theoretical estimations [33] T_p may sufficiently surpass the ion temperature. For definiteness sake, we take $T_p = 1000$, $Z_p = 10^4$, $\kappa = 400\ \mathrm{cm}^{-1}$, then the condition of the hierarchy stability becomes $|\varphi_{min}| > 10^{-5}$. According to the approach under discussion (see the graph in Fig. 4), this condition is fulfilled when $M > 0.01$.

It should be noticed here that at the condition (8) the condition of the linear approximation $|e\Phi_{min}| \ll k_B T_i$ may be violated. At the opposite limiting case $|e\Phi_{min}| \gg k_B T_i$ the perturbations of the ion density and Φ_{min} in the ion-dusty plasma should be proportional to $Z^{1/2}$ [34] and $U \propto Z^{3/2}$ in contrast to linear theory. Really, strong decay of the wake potential dependence on Z was found by Winske [35] by numerical dusty plasma simulation at $M = 1.2$ and $Z > 3.2 \times 10^4$. Besides, in the theory of nonlinear energy loss of highly charged ($Z \sim 10^2$) heavy ions, there was found [36] that nonlinear stopping depended on the ion charge state as $Z^{3/2}$ at $M < 1$ and became proportional to Z^2 as M increased. On this basis, the modified condition of the hierarchy stability at the same values of dusty plasma parameters becomes $|\varphi_{min}| > 10^{-3}$ which is fulfilled at $M > 0.1$.

III. SURFACE TENSION OF DUSTY PHASE

One of the characteristic feature of a dispersed phase in dusty plasma is presence of sharp (nondiffusion) boundary surface of a dusty cloud ("dusty drop"). Considerable recent attention has been focused on the problem of formation of voids ("dusty bubbles") of the dispersed phase. It is reasonable to suppose that both phenomena are of the same nature related with presence of a surface tension resulted from the attractive interaction between charged particles.

In general, the surface tension coefficient of an interphase boundary is defined as the integral of a difference between transverse $p_\perp$ and longwise $p_\parallel$ pressures over an interphase layer of thickness l,

$$\gamma = \int_{-l/2}^{l/2} [p_\perp(z_1) - p_\parallel(z_1)] dz_1, \qquad (9)$$

where the axis z_1 directed along the normal to the layer.

On the molecular level, a pressure can be written in the form of a virial equation of state and expressed via a number density $n(z_1)$, virial of a force of intermolecular interaction $\mathbf{F}(\mathbf{r})$ and radial correlation function $g(\mathbf{r})$, where $\mathbf{r} = \mathbf{r}_2 - \mathbf{r}_1$ is the distance between particles. For a bulk equilibrium fluid the function $g(\mathbf{r})$ is calculated as a rule with the use of the Born–Green equation. At an interphase surface the situation is much more complex; this being so, notably rough approximations are used here. One of them is Fowler's step approximation, when the number density $n(z_1)$ is regarded as constant n_0 within the liquid phase and zero outside the liquid, and radial function $g(\mathbf{r})$ within the surface layer is assumed to be the same as in the bulk liquid. Then the surface tension coefficient is determined by the Fowler formula [37,38]

$$\gamma = -\frac{\pi n_0^2}{8}\int_0^\infty r^4 F_r(r) g(r) dr, \quad z > 0. \qquad (10)$$

It is implied here that the z component of the vector $\mathbf{r}$ is normal to the surface. Kirkwood and Buff [39] estimated the surface tension of liquid argon on the base of this formula and obtained quite satisfactory value (error $\sim$25%).

It is supposed in the Fowler formula that the intermolecular interaction is spherically symmetric that does not allow to apply this formula as such to the boundary of dusty phase.

Rederiving the Fowler formula with due regard for the anisotropy of interparticle interaction we obtain

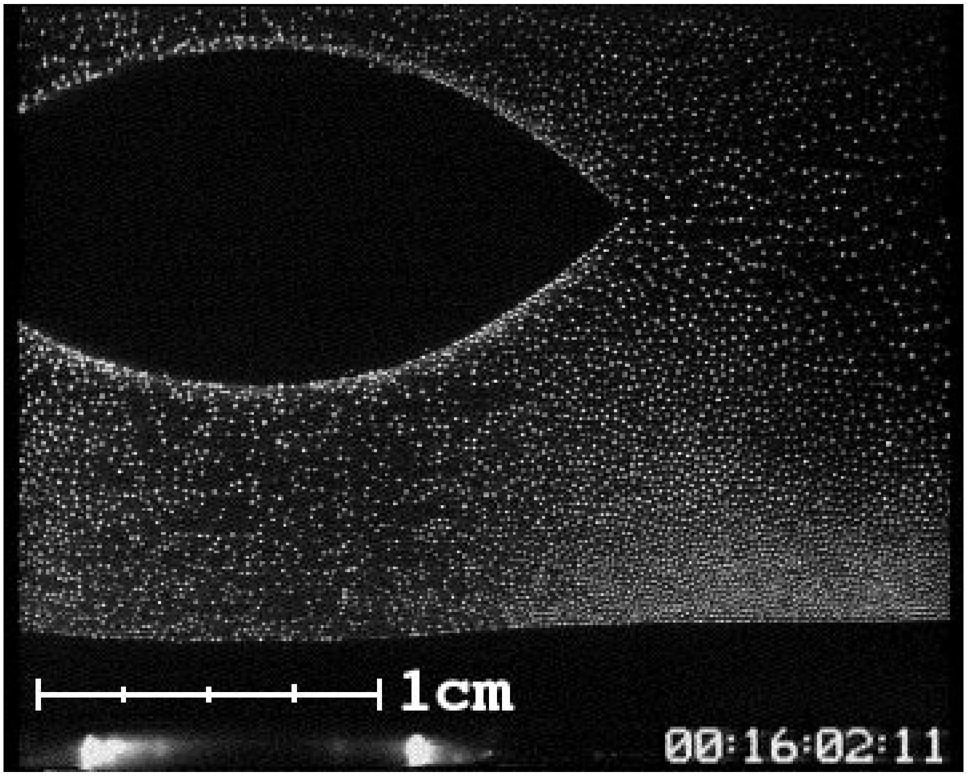

FIG. 6. The typical void in dusty plasma. (Copied from Ref. [40] with Professor H. Thomas' kind permission.)

$$\gamma = -\frac{n_0^2}{2}\int [z^2 F_z(\mathbf{r}) - xzF_x(\mathbf{r})]g(\mathbf{r})d\mathbf{r}, \quad z>0. \tag{11}$$

where F_z and F_x are z and x components of the force (7).

When applying to the boundary of dusty phase, the pressures $p_\perp$ and $p_\parallel$ are to be considered as partial pressures of dusty component. Then, the surface tension of dusty phase should be interpreted as partial property too.

To estimate its value some simplification will be introduced. First, the radial distribution will be used in the form of the Boltzmann factor

$$g(\mathbf{r}) = \exp\left\{-\frac{U(\mathbf{r})}{k_B T}\right\},$$

which is typical for dilute systems. Second, when estimating the surface tension, the interactions within hierarchies will be regarded only as it is much stronger interactions between particles from different hierarchies. Then, for surfaces oriented across and along the ion flow we get

$$\gamma_\perp = -\frac{n_p^2}{2}\int z^2 F_z(\mathbf{r})g(\mathbf{r})d\mathbf{r}, \quad z>0, \tag{12}$$

$$\gamma_\parallel = \frac{n_p^2}{2}\int xzF_x(\mathbf{r})g(\mathbf{r})d\mathbf{r}, \quad z>0. \tag{13}$$

The axis z in both equations directed along the normal to the surface and the forces $F_z(\mathbf{r})$ in the equation for $\gamma_\perp$ and $F_x(\mathbf{r})$ in the equation for $\gamma_\parallel$ are determined by interactions along hierarchies.

A numerical estimation of these expressions for $n_p \sim 10^6\ \mathrm{cm}^{-3}$ and $M=0.01$ leads to $\gamma_\perp \sim 10^5\ \mathrm{dyn\,cm}^{-1}$ and $\gamma_\parallel \sim -10^4\ \mathrm{dyn\,cm}^{-1}$. Both of these coefficients in absolute magnitude surpass sufficiently the typical surface tension coefficient of a liquid at normal conditions $\gamma_{liq} \simeq 10^2\ \mathrm{dyn\,cm}^{-1}$ (as an example, for water $\gamma_{liq} = 70\ \mathrm{dyn\,cm}^{-1}$). Such great difference is quite natural. Really, the negative minimum U_m of the potential U at given conditions is of the order of -10^{-12} erg, while the depth of the potential well of, for example, the Lennard-Jones potential for noble gases is of the order of -10^{-14} erg. In as much as the potential U enters (with sign minus) into the exponent of the radial distribution function, the surface tension coefficient appears to be very sensitive to its value. As a result, we have $|\gamma_{\perp,\parallel}| \gg \gamma_{liq}$ in spite of $\gamma \propto n^2$ and $n_p \ll n_{liq}$.

Hydrodynamic stability of a spherical gas bubble or liquid drop is determined by value of a Weber number $We = \rho v^2 a/\gamma$, which is the ratio of a dynamic head ρv^2 to a surface tension pressure γ/a, where a is the radius of the bubble or drop. For the case of the void in dusty plasma, the dynamic head ρv^2 is determined by the flow of ions and neutral gases with the velocity $v = u = \tilde{v} M \simeq 5 \times 10^4\ M\ \mathrm{cm\,c}^{-1}$. We can put $\gamma \simeq \gamma_\perp$ for the slightly curved surface of a lenticular void with the radius of curvature $a \simeq 2$ cm sufficiently greater its size. Then at characteristic pressure of the order of 0.5 torr we have $We \simeq 10^4 M^2/\gamma_\perp$. The coefficient $\gamma_\perp$ increases with increasing of M, so that the Weber number We attains the greatest value at the minimal from values of M, i.e. at $M=0.01$ when $\gamma_\perp \simeq 10^5$, considered here. Then $We_{max} \simeq 10^{-4}$. When the Weber number is so small the surfaces of the lenticular void are stable, certainly. Their deviations from a spherical form can be due to variations of γ relative to $\gamma_\perp$ as a consequence of inevitable (although small) distortions of the condition of orthogonality to the flow on the curved surface.

The inclusion of nonlinear effects mentioned briefly at the end of Sec. II can increase the low threshold value of M from 0.01 to 0.1 in estimations given above. However, this does not change qualitatively the foregoing picture.

Stability of the lenticular void ensures possibility of its existence. The problem of its appearance remains to be solved. It may be suggested that a phase transition of the first kind takes a place in a homogeneous dusty plasma, and strong anisotropic surface tension can give rise to a peculiar nucleation process.

IV. CONCLUSIONS AND DISCUSSIONS

The wake potential model of attractive interactions between likely charged colloidal or dust particles in plasma have been discussed here in a view to explain the observed crystal structure formed from these particles. The velocity u of the ion flow relative to the dust particles is not necessary to exceed the heat velocity $\tilde{v}$ of counter-ions (or ion–acoustic velocity) as it was supposed in Ref. [10]. Even at small u when $u/\tilde{v} \simeq 10^{-2}$, the effective attractive interaction appears to be sufficiently strong to ensure stability of crystal-like structure.

Another consequence of strong attractive interparticle interactions is the partial surface tension γ of the dusty phase boundary.

This concept was briefly discussed by Tsitovich [23] who estimated γ as a work necessary for construction of a bulk liquid of the unit surface and height h. As a result, he obtained $\gamma = U_m n_p h$, then he took $|U_m| = 100$ eV that correspond to our result at $M=1$ and found $\gamma \simeq 10^{-2}$dyn cm^{-1} neglecting the negative sign of U_m. If he would account for $U_m < 0$ he would get $\gamma < 0$, which is natural as the work for construction of the coupling liquid system from noninteracting gas particles is to be negative.

Here, the surface tension is estimated on the basis of a generalization of the Fowler's formulas taking into account the anisotropy of interparticle interactions. As a result, the strong dependence of γ on orientation of the dusty interface is found. In particular, when $u/\tilde{v} \simeq 10^{-2}$ we get $\gamma_\perp \simeq 10^5$ dyn cm^{-1} and $\gamma_\parallel \simeq -10^4$ dyn cm^{-1} for surfaces oriented across and along the ion flow, correspondingly. So great negative value of $\gamma_\parallel$ exclude a possibility of existence of an interface oriented along the ion flow. Lenticular void fulfills this condition. The same can be said about pancake-shaped voids as they having no surfaces along the ion flow too. Such forms were observed both under microgravity condition (see Fig. 6) and in terrestrial experiments [18–22]. Samsonov and Goree [18] (see Fig. 7) observed also an appearance of a void mode as a penetration of a finger-shaped (in a vertical section) dusty free region through the side boundary of the gas discharge. The process as a whole is similar to a gain of a vapor bubble on the bottom of a pot, and may be interpreted as a heterogeneous nucleation on a wall. A sharp cusplike end of the finger and its fast travel across the gas discharge volume count in favor of negative surface tension of the interface at its end. In the experiment of Thomson *et al.* [41] a void was formed around an object immersed in the dusty plasma and in the experiments of Merlino *et al.* [42,43] a void has been created with the use of electrodes inserted into the dusty plasma of a glow discharge. These experiments may as well be interpreted as examples of heterogeneous nucleation. It quite natural to expect such phenomena at appearance of an embryo of a new phase in a system with a surface tension at interface.

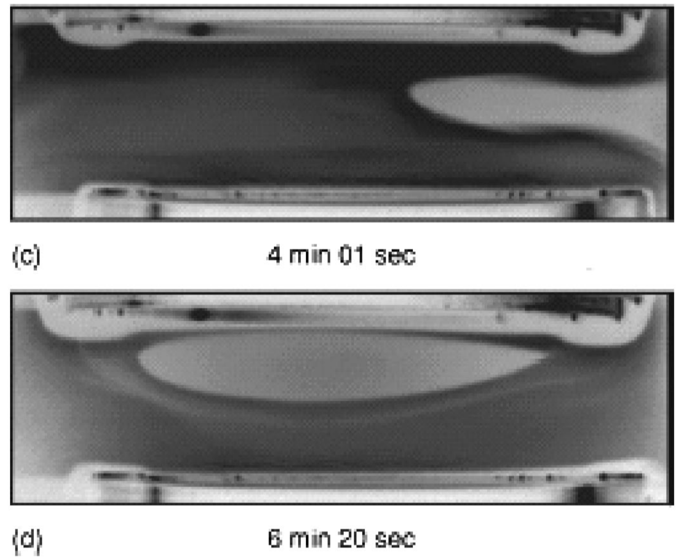

FIG. 7. Formation of the void. (Copied from Ref. [18] with Professor J. Goree kind permission.)

An experimental observation of a spherical void would be the crucial experiment for the model under discussion. An existence of a part of the dusty phase interface along the ion flow would demand a radical reconsideration of the wake potential model to include some kind of isotropic interaction between dusty particles. Up to now I did not find any experimental evidences of the spherical void in literature and internet. Moreover, authors of papers on the theory of spherical voids [24,25], while foundation of their model of spherical void, referred to the experimental works dealing with the lenticular and pancake-shaped voids and even reproduced a photo of the lenticular void from Ref. [18].

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