Interaction of dust grains in strong collision plasmas: Diffusion pressure of nonequilibrium charge carriers

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The formation of the nonequilibrium charge carriers in the space-charge layer at the interphase boundary in thermal collision plasmas has been studied. It has been shown that the flux of the nonequilibrium charge carriers determines the diffusional pressure on the interphase boundary, which provides for the existence of the force, directed to the area of greater plasma perturbation. The comparison of the electrostatic and diffusional pressures shows the possibility of the balance of the forces and the formation of the ordered structures of dust grains in a complex strong collision plasma. © *2005 American Institute of Physics*. [DOI: 10.1063/1.2062868]

I. INTRODUCTION

The interaction of charged dust grains in complex plasmas differs from the Coulomb interaction in vacuum by various collective effects, which may result in the attraction of the likely charged dust grains.[1] It is clear that the reasons for such interaction should lay beyond the electrical nature. It is assumed, for example,[2–7] that the flux of plasma particles towards one dust grain carries away other dust grains, located around, thus creating an additional force of attraction between the dust grains—the so-called shadow force. For the first time, this mechanism has been considered by Tsytovich,[3] with reference to the low-pressure gas-discharge plasma. Secondly, the transfer of momentum from the directed ion flux to a dust grain is also possible. This force is called the ion drag force.[8–12] In the plasma of combustion products there are also forces similar to the ion drag force, which are caused by the anisotropy of ionic pressure on the grain surface.[13]

In the present paper the formation of nonequilibrium ions and electrons in the space-charge layer at the interphase boundary in the thermal collisional plasmas and their influence on the gas phase pressure on the grain surface are considered.

II. THE STATE OF THE PROBLEM

The subject of the study is low-temperature thermal complex plasma. The thermal complex plasma consists of a buffer gas (usually air) at the atmospheric pressure and easily ionizable additional agents—the atoms of alkali metals. At the temperatures of 1000–3000 K the easily ionizable atoms of additional agents become partially ionized by the collisions of gas particles, forming the low-temperature plasma. Such plasma, being in equilibrium, can be considered within the limits of ideal-gas approximation. The insertion of metal electrodes into the low-temperature plasma or the formation of a condensed phase in the plasma causes the displacement of the system from equilibrium state. However, if perturbations are small, it is still possible to consider the plasma as an ideal or a weakly coupled plasma.

The ionization equilibrium in the weakly coupled thermal plasma is described by the Saha equation:

$$\frac{n_{e0}n_{i0}}{n_{a0}} = 2\frac{g_i}{g_a}\left(\frac{m_eT}{2\pi\hbar^2}\right)^{3/2}\exp\left(\frac{-I_0}{T}\right) \equiv K_S, \tag{1}$$

where K_S is the Saha constant, g_i and g_a are the statistical weights of ions and atoms, I_0 is the ionization potential of the alkali-metal atoms, T is the equilibrium temperature of the plasma, which in this case is an isothermal system, m_e is the electronic mass, and $\hbar$ is the Planck constant.

Here the conditions of the conservation of mass and charge should be observed,

$$n_{i0} + n_{a0} = n_A, \quad n_{e0} = n_{i0} = n_0, \tag{2}$$

where n_A is the number density of additional agents, and n_0 is the unperturbed number density.

In this case the gas particles can be described by the equilibrium distribution functions. In particular, in the area where there is the electrostatic potential ϕ the electron and ion number densities are described by the Boltzmann distribution law,

$$n_{e0} = n_0\exp(e\phi/T), \quad n_{i0} = n_0\exp(-e\phi/T). \tag{3}$$

If in such plasma there is the charged phase boundary (the condensed dust grain or an electrode), the charge conservation equation in the form of Eq. (2) is false. Let us consider the plasma layer contacting a flat metal electrode. For establishing the equilibrium in the contact the Fermi level of metal F_m must be equal to the level of the electrochemical potential of the plasma (a Fermi level of plasma) F_{pl}. Let the electronic work function of a plasma be greater than the electronic work function of a metal $W_{\text{pl}} > W_m$. Then the energy of an electron at the level F_{pl} is less, than at the level F_m and for Fermi level equalization some electrons should transfer from the metal into the plasma. Thus the metal is charged positively and in the plasma formed the space-charge layer enriched with electrons. The total of electrons in the plasma is increased, which cause additional ionization of the alkali-metal atoms. Therefore the condition of electroneutrality [Eq. (2)] and distributions [Eq. (3)] becomes false. In this case it is necessary to modernize the

Saha equation by introducing an effective ionization potential $I_{\text{eff}}=I_0+\psi$, where the parameter ψ characterizes the displacement of the ionization equilibrium as a result of external influence on the plasma,[14]

$$\frac{n_e n_i}{n_a} = K_S \exp \frac{\psi}{T}, \tag{4}$$

and the value $\sqrt{n_e n_i}=n_q$ is the quasiunperturbed number density.

Then the distributions of charge carriers are described by the following expressions:

$$n_e = n_q \exp(e\phi/T), \quad n_i = n_q \exp(-e\phi/T). \tag{5}$$

The ionization degree of the low-temperature plasma is usually so low, that $n_{a0}\cong n_a \cong n_A$; therefore from Eqs. (1) and (4) the quasiunperturbed number density is defined by the following expression:

$$n_q = n_0 \exp(\psi/2T). \tag{6}$$

The parameter ψ linearly corresponds to the bulk plasma potential:[13,14] $\psi=-e\varphi_{\text{pl}}$, which characterizes the volumetric charge of the whole plasma and depends on the potential barriers at the phase boundary. In particular, for a semi-infinite plasma with the potential barrier at the phase boundary $e\phi_s$ the bulk plasma potential is defined by the following expression:[15]

$$\varphi_{\text{pl}} = -2\frac{T}{e} \tanh\left(\frac{e\phi_s}{4T}\right). \tag{7}$$

Accordingly, the quasiunperturbed number density is defined by the following expression:

$$n_q = n_0 \exp[\tanh(e\phi_s/4T)]. \tag{8}$$

The displacement of the ionization equilibrium means the appearance of nonequilibrium charge carriers, which are not described by the Boltzmann distribution law. Accordingly, there are the fluxes of nonequilibrium carriers which interact with the phase boundary.

The electrostatic interaction of dust grains in the example of the charged planes has been considered in Ref. 15, taking into account the variation of the parameter ψ and the bulk plasma potential. The influence on this interaction of nonequilibrium charge carriers is addressed in the present paper.

III. SPACE-CHARGE REGION

The parameter ψ and the bulk plasma potential φ_{pl} characterize the whole volume of the plasma, i.e., describes the volumetric average of the ionization degree change, induced by the interfacial interaction. However, in the space-charge region there is a local displacement of the ionization degree induced by the inhomogeneity of the space distribution of gas particles.

The conservation of the charge carriers is defined by the continuity equations:

$$\frac{\partial n_e}{\partial t} + \nabla \cdot \mathbf{j}_e = G_e, \quad \frac{\partial n_i}{\partial t} + \nabla \cdot \mathbf{j}_i = G_i, \tag{9}$$

where $\mathbf{j}_{e(i)}$ is the flux of electrons (ions), and the source function $G_{e(i)}$ represents the electron (ion) number density variation rate as a result of ionization—recombination processes.

In any volume of the plasma which does not contain the phase boundary, the ionization degree is defined only by the collision ionization of alkali-metal atoms and the volumetric recombination,

$$G_e = G_i = G = \beta n_e n_a - \gamma n_e n_i = \beta n_e n_A - (\beta + \gamma) n_e n_i, \tag{10}$$

where β is the constant of alkali-metal atom ionization rate, and γ is the electron-ion recombination coefficient. Here, we have assumed that $n_i+n_a=n_A$.

Beyond the space-charge region there are equilibrium values of ionization and recombination,

$$G_0 = \beta n_q n_A - (\beta + \gamma) n_q^2 = 0, \tag{11}$$

that is $\beta n_A=(\beta+\gamma)n_q$.

The existence of the volumetric charge in the space-charge layer results in the disbalance between the recombination rate $\gamma n_e n_i=\gamma n_q^2$ and the ionization rate $\beta n_e n_a = \beta n_q \exp(e\phi/T) n_A - \beta n_q^2$, which depends on the potential. Therefore, $G\neq 0$ causes the displacement of the ionization degree in the space-charge layer and forms the nonequilibrium charge carriers, thus the positive potential induces the increase of the ionization degree, and the negative potential—the decrease. It results in the disbalance between diffusive and drift fluxes:

$$\begin{aligned} j_e &= \mu_e n_e^* \nabla \phi - D_e \nabla n_e^* \neq 0, \\ j_i &= -\mu_i n_i^* \nabla \phi - D_i \nabla n_i^* \neq 0, \end{aligned} \tag{12}$$

where $\mu_{e(i)}$ is the electron (ion) mobility, $D_{e(i)}$ is the electron (ion) diffusivity, and "*" means the nonequilibrium character of the charge-carrier number density.

The nonequilibrium number densities can be presented in the form of a deviation from the equilibrium values $n_e^* = n_e + \delta n$ and $n_i^* = n_i + \delta n$, thus the deviation should be identical for the ions and the electrons as the value of the volumetric charge $n_i^* - n_e^* = n_i - n_e = \rho/e$ should be preserved.

In Fig. 1 the sample space distributions of the electron and the ion number densities in the space-charge layer for positive and negative fields are presented. If the field is positive [Fig. 1(a)], the excess electrons increment the negative drift flux and increment the positive diffusion flux because the gradient of the electron number density is incremented. Thus, the drift and diffusion fluxes of the nonequilibrium electrons as well as the equilibrium fluxes are directed to different sides, $j_{eE}^* = \mu_e \delta n \nabla \phi < 0$ and $j_{eD}^* = -D_e \nabla \delta n > 0$.

The drift flux of nonequilibrium ions coincides in terms of the direction with the equilibrium drift flux, i.e., it is positive, $j_{iE}^* = -\mu_i \delta n \nabla \phi > 0$. The equilibrium diffusion flux of ions is negative, but the excess ions diminish the gradient of the number density and, accordingly, diminish the diffusion flux that is equivalent to the occurrence of the positive flux,

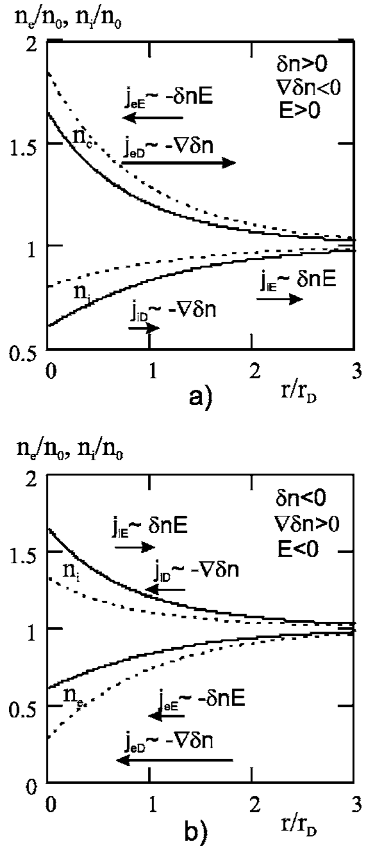

FIG. 1. The space distribution of equilibrium (full curves) and nonequilibrium (dotted curves) number densities in the space-charge layer: (a) the positive field and (b) the negative field.

$$j_{iD}^{*} = (-D_i \nabla n_i^{*}) - (-D_i \nabla n_i) = -D_i \nabla \delta n > 0, \tag{13}$$

because for equilibrium fluxes the balance is observed. Hence, the drift and the diffusion fluxes of nonequilibrium ions are directed to one side—on the field. It should be noted that the gradient of nonequilibrium ions is negative, as the nonequilibrium ionization is decreasing together with diminution of the volumetric charge.

If the field is negative [Fig. 1(b)], the ion number density is increasing in the space-charge layer, but the decrease of the ionization degree results in the decrease of the number density and the gradient of the number density, which cause the decrease of the drift and diffusion fluxes. As a result the fluxes of nonequilibrium ions are opposite to the equilibrium fluxes $j_{iE}^{*}=-\mu_i \delta n \nabla \phi > 0$ and $j_{iD}^{*}=-D_i \nabla \delta n < 0$.

If the field is negative, the positive electron drift flux decreases, because the electron number density is decreasing, which is equivalent to the occurrence of a negative flux of the nonequilibrium electrons $j_{eE}^{*}=\mu_e \delta n \nabla \phi < 0$. The direction of the diffusion flux of the nonequilibrium electrons coincides with the direction of the equilibrium flux, because the decrease in ionization degree increases the gradient of the electron number density $j_{eD}^{*}=-D_e \nabla \delta n < 0$.

Thus, in both cases there are the noncompensating fluxes of the nonequilibrium charge carriers codirectional with the electric field. Accordingly, there is an ambipolar diffusion flux, caused by the flux of nonequilibrium ions at the positive field and the flux of nonequilibrium electrons—at the negative field. The equality of electronic and ionic fluxes is ensured by the field of the ambipolar diffusion,

$$\frac{eE_a}{T} = -\frac{e(D_e - D_i)}{T(\mu_e + \mu_i)\delta n} \nabla \delta n, \tag{14}$$

whence taking into consideration the Einstein relation $\mu = eD/T$, we obtain the flux of nonequilibrium carriers,

$$j^{*} = -2\frac{D_e D_i}{D_e + D_i} \nabla \delta n, \tag{15}$$

where $2D_e D_i/(D_e + D_i) = D$ is the ambipolar diffusion coefficient.

The ambipolar diffusion flux ensures the transition of nonequilibrium charge carriers. At the positive field the nonequilibrium electrons and ions move to the outside of the space-charge layer. If the dust grain is negative the field is also negative, i.e., directed towards the grain. In this case the ionization degree in the space-charge layer decreases, and the stream of nonequilibrium carriers is directed towards the grain.

The conservation of the nonequilibrium carrier number density is defined by the continuity equation:

$$\frac{\partial \delta n}{\partial t} + div j^{*} = G^{*}, \tag{16}$$

where $div j^{*} = -D\Delta\delta n$, and taking into account Eq. (11), $G^{*} = \beta n_A n_e^{*}(1 - n_i^{*}/n_q)$.

In the stationary case Eq. (16) is reduced to the following form:

$$\lambda_R^2 \Delta\left(\frac{\delta n}{n_q}\right) = \left[\frac{1}{1 + (er_D E/T)^2}\left(\frac{\delta n}{n_q}\right)^2 + \frac{\delta n}{n_q} - \frac{\exp(e\phi/T) - 1}{1 + (er_D E/T)^2}\right][1 + (er_D E/T)^2], \tag{17}$$

where $\lambda_R = \sqrt{D/\gamma n_0}$ is the recombination length (or diffusion length of nonequilibrium carriers).[16]

Here an approximation was taken into account that $2\cosh(e\phi/T) - 1 = 1 + (er_D E/T)^2$. At small values of the potential $e\phi \ll T$ the ratio $\delta n/n_q \ll 1$, therefore in Eq. (17) it is possible to neglect the quadratic term. However, it can also be made at more potentials, as in this case $\delta n/n_q \sim 1$, but $er_D E/T \gg 1$.

The solution of Eq. (17) depends on the relation between the Maxwell relaxation time $\tau_M = r_D^2/D$ and the lifetime of the nonequilibrium carriers $\tau_R = \lambda_R^2/D$. In the weak collision plasma $\tau_M \ll \tau_R$. But in the thermal plasmas at the atmospheric pressure, which is strongly collisional, $\tau_R \sim 10^{-8} - 10^{-9}$s and $\tau_M \sim 10^{-6} - 10^{-7}$s, i.e., $\tau_M \gg \tau_R$, accordingly $r_D \gg \lambda_R$. It means that in such plasmas the ionization balance establishes much more quickly than the diffusion-drift balance. Therefore, it is possible to consider the potential and the field as constant values on the recombination length. Then the solution of Eq. (17) is the function

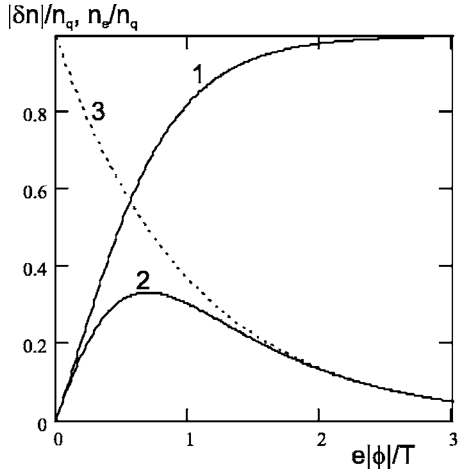

FIG. 2. The dependencies of the nonequilibrium carrier number density on the potential: (1) at the increase of ionization ($\phi>0, \delta n>0$), (2) at the decrease of ionization ($\phi<0, \delta n<0$), and (3) the changes of the electron number density at $\phi<0$.

$$\delta n = n_q \frac{\exp(e\phi/T)-1}{1+(er_D E/T)^2}\left\{1+\frac{\lambda_R}{r_D}\times\exp\left[\frac{-r}{\lambda_R}\sqrt{1+(er_D E/T)^2}\right]\right\}. \quad (18)$$

Accordingly, the flux of the nonequilibrium charge carriers,

$$j^* = \frac{Dn_q}{r_D}\frac{\exp(e\phi/T)-1}{\sqrt{1+(er_D E/T)^2}}\exp\left[\frac{-r}{\lambda_R}\sqrt{1+(er_D E/T)^2}\right]. \quad (19)$$

The shift of the ionization degree in the space-charge layer occurs in different ways at positive and negative potentials of the surface. The dependencies of the nonequilibrium carrier number density on the potential for coordinate $r=0$ are shown in Fig. 2. Here we can see that at the positive potential δn tends to n_q, which compensates the decrease of the ion number density in the space-charge layer. At the negative potential the electron number density decreases in the space-charge layer ($\delta n<0$); the ionization degree decreases too. However, the rate of decrease of ionization (the increase of $-\delta n$) cannot be more than the electron number density, that is $-\delta n \leqslant n_e = n_q \exp(-e|\phi|/T)$. Therefore, the large negative potential causes the decrease of the nonequilibrium charge carriers.

IV. PRESSURE ON THE INTERPHASE SURFACE

Let us consider the transfer of momentum to the dust grain by the gas particles. We shall exclude the buffer gas as its pressure is spherically symmetrical in relation to the grain. Therefore, further, for the purpose of simplicity, the gas phase and the gas particles will be understood only as electrons, ions, and atoms of additional agents. Then, the transfer of total momentum from the gas phase to the dust grain occurs as follows. There are sporadic collisions of gas particles with the grain surface which ensure average fluxes $j_{a0}=(1/4)n_{as}v_{Ta}$ and $j_{i0}=(1/4)n_{is}v_{Ti}$ (in approximation of the flat surface), which correspond to the momentum densities, transferred by the gas particles $p_{a0}=(1/4)\lambda_a m_a n_{as} v_{Ta}$ and $p_{i0}=(1/4)\lambda_i m_i n_{is} v_{Ti}$, and the electronmomentum can be neglected. Here $\lambda_{a(i)}$ is the free length of atoms (ions), and $v_{Ta(i)}=\sqrt{8T/\pi m_{a(i)}}$ is the thermal velocity of atoms (ions).

This interaction is increased or decreased by the flux of nonequilibrium ions [Eq. (19)], that depend on the direction of the field. Hence, taking into account that the free lengths $\lambda_i \cong \lambda_a$, the masses $m_i \cong m_a$, the thermal velocities $v_{Ti} \cong v_{Ta}$, and the number densities $n_i+n_a=n_A$, the momentum transferred by the gas particles to a unit area of the grain surface in the direction of the center of the grain is equal to

$$p = (1/4)\lambda_a m_a v T_a n_A - \lambda_R m_i j^*. \quad (20)$$

Here it is considered that the nonequilibrium ions migrate not along the free length λ_i but along the recombination length λ_R. As the anisotropy of interaction is defined by the nonequilibrium ions, the pressure of gas on the grain surface can be defined as $P=p/\tau_R$, where τ_R is the lifetime of nonequilibrium ions,

$$P = P_0 - \frac{m_i n_q \lambda_R D}{r_D \tau_R}\frac{\exp(e\phi_s/T)-1}{\sqrt{1+(er_D E_s/T)^2}}, \quad (21)$$

where ϕ_s and E_s are the surface values of the potential and the field.

If the dust grain is charged positively, the ionization degree in the space-charge layer increases, the field is positive, the nonequilibrium ionflux causes the decrease of pressure [Eq. (21)] on the grain surface, and greater decrease of the pressure corresponds to the greater ionization degree ($n_q > n_0$). If the dust grain is charged negatively, the ionization degree in the space-charge layer decreases, the field is directed toward the grain, the pressure on the grain surface is increased by the nonequilibrium ionflux from the volume of the plasma into the space-charge layer, and the greater magnification of pressure corresponds to the greater ionization degree.

In any case the surface pressure decreases as the shift of the ionization degree from the equilibrium value increases. Hence, if there is a space inhomogeneity of the ionization degree near the dust grain, then there is anisotropic pressure on the grain surface, ensuring the existence of the average force,

$$\mathbf{F} = \int_S P d\mathbf{s} = -\frac{m_i \lambda_R D}{r_D \tau_R}\frac{\exp(e\phi_s/T)-1}{\sqrt{1+(er_D E_s/T)^2}}\int_S n_q d\mathbf{s}. \quad (22)$$

Equation (22) is obtained with the assumption that the neighboring grains are far enough from the given grain and the space-charge layers are not overlapped. In this case, in Eq. (22) all terms except n_q are radially symmetrical in the space-charge layer and can be put outside the integral. Then, the influences of the neighboring grains are expressed only in the asymmetry of the quasiunperturbed number density, which is determined by the asymmetry of the bulk plasma potential:

$$n_q = n_0 \exp(\psi/2T) = n_0 \exp(-e\varphi_{\rm pl}/2T). \quad (23)$$

Let us compare the pressure of nonequilibrium carriers [Eq. (21)] with the electrostatic pressure between the metal

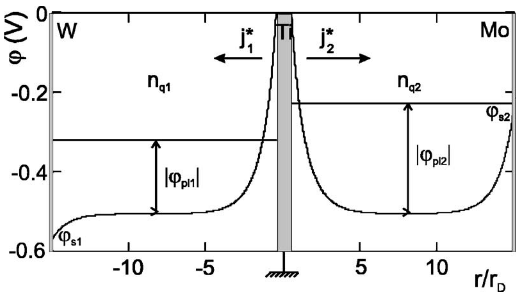

FIG. 3. The spatial potential distribution between electrodes.

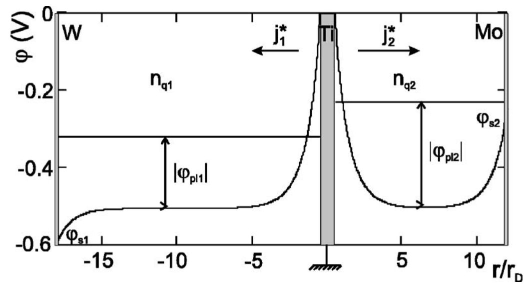

FIG. 4. The distribution of the full potential between the electrodes after reaching the balance of forces.

planes in the plasma. For the purpose of definiteness we shall consider the titanium plane with the work function $W_0=3.9$ eV, placed in the plasma containing an additional of agent of potassium with the number density $n_A=2\times 10^{15}\ \mathrm{cm}^{-3}$ at the temperature $T=0.17$ eV. In this case the potential barrier at the phase boundary is $e\phi_{s0}=0.51$ eV, the unperturbed number density of charge carriers is $n_0=2\times 10^{12}\ \mathrm{cm}^{-3}$, and the screening length is $r_D=1.5\ \mu$.

Let us place the Wolfram plane ($W_1=4.5$ eV) to the left of the titanium plane and the molybdenum one ($W_2=4.15$ eV) to the right. The equilibrium values of the potential barriers are $e\phi_{s1}=-0.09$ eV and $e\phi_{s2}=0.26$ eV. If the central titanium plane is earthed, we obtain the space distribution of the full potential between the planes, as it is shown in Fig. 3. Thus, the values of the surface potentials with respect to the earthed plane are $\varphi_{s1}=-0.6$ V and $\varphi_{s2}=-0.25$ V.

The calculation of the electrostatic pressure on the plane has been described in Ref. 15. In this case, the pressure on the central plane is defined by the following expression:

$$P_E=\frac{1}{8\pi}(E_{\mathrm{right}}^2-E_{\mathrm{left}}^2), \tag{24}$$

where the electric field at the left, E_{left}, of the plane and at the right, E_{right}, of the plane are equal to

$$|E_{\mathrm{left(right)}}|=\frac{2T}{er_D}\sqrt{\sinh^2\left(\frac{e\phi_{s0}}{2T}\right)+\delta_{1(2)}}. \tag{25}$$

The constant δ defines the pattern of the potential distribution,

$$\delta_{1(2)}=-4\tanh\left(\frac{e\phi_{s0}}{4T}\right)\tanh\left[\frac{e\phi_{s1(s2)}}{4T}\right]\exp\left[-\frac{d_{1(2)}}{r_D}\right], \tag{26}$$

where d is the thickness of the plasma layer.

As a result, the electrostatic pressure on the central plane is

$$P_E=\frac{2}{\pi}\left(\frac{T}{er_D}\right)^2\tanh\left(\frac{e\phi_{s0}}{4T}\right)\left[\tanh\left(\frac{e\phi_{s1}}{4T}\right)e^{-d_1/r_D}-\tanh\left(\frac{e\phi_{s2}}{4T}\right)e^{-d_2/r_D}\right]. \tag{27}$$

At the distance between the planes $d_1=d_2=15r_D=22\mu$, the electrostatic pressure is $P_E=-6.1\times 10^{-11}\ \mathrm{N/cm}^2$. The electrostatic forces influencing the central plane are attractive—on the left—and repulsive—on the right.

On the left and right of the central earthed plane there are different values of the bulk plasma potentials $\varphi_{\mathrm{pl1}}=-0.19$ V and $\varphi_{\mathrm{pl2}}=-0.28$ V, which ensure the different values of the quasiunperturbed number densities $n_{q1}=3.8\times 10^{12}\ \mathrm{cm}^{-3}$ and $n_{q2}=5.0\times 10^{12}\ \mathrm{cm}^{-3}$. The fluxes of the nonequilibrium charge carriers from the plane into the volume of plasma, that fits these values n_q, are $j_1^*=-1.3\times 10^{15}\ \mathrm{cm}^{-2}\ \mathrm{s}^{-1}$ on the left and $j_2^*=1.7\times 10^{15}\ \mathrm{cm}^{-2}\ \mathrm{s}^{-1}$ on the right. The recombination length is $\lambda_R=0.1\mu$, and the lifetime is $\tau_R=4\times 10^{-9}$ s. This provides for the existence of the diffusion pressure $P_D=9\times 10^{-10}\ \mathrm{N/cm}^2$, which is opposite to the electrostatic pressure.

Let the central plane be not anchored. Then the total pressures P_E+P_D force it to move to the right, to the molybdenum plane. As we consider the perpetual planes, the bulk plasma potential and the quasiunperturbed number density do not depend on the distance between the planes. Therefore the pressure P_D does not depend on the distance between the planes either. However, the electrostatic pressure P_E depends on this distance. Therefore, when the central plane moves under the diffusion pressure P_D to the right for $3r_D=4.5\mu$, the electrostatic pressure increases up to $P_E=9\times 10^{-10}\ \mathrm{N/cm}^2$ and completely counterbalances the pressure P_D. Thus, the spatial potential distribution between the planes becomes balanced as in Fig. 4. The existence of the diffusion pressure of the nonequilibrium charge carriers provides for the necessity of attraction of the likely charged planes in order to establish the balance in the plane disposition.

The same can be applied also to the charged dust grains. The electrostatic forces and the diffusion forces, caused by the pressure of the nonequilibrium charge carriers, as a result of inhomogeneous ionization of the plasma, can provide for the balance of the space disposition of the charged dust grains.

V. CONCLUSION

It has been demonstrated that the nonequilibrium ionization in the space-charge region at the interphase boundary in the plasma plays an essential role in the interaction of the dust grains. The flux of the nonequilibrium charge carriers changes the pressure of the gas phase on the condensed phase and provides the diffusion pressure of the nonequilib-

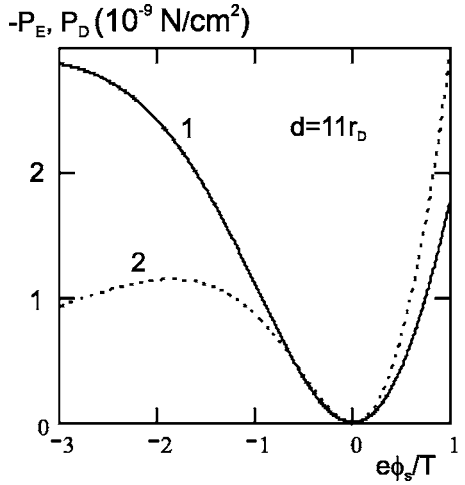

FIG. 5. The dependencies of the electrostatic pressure P_E (1) and the diffusion pressure of the nonequilibrium charge carriers P_D (2) on the surface potential.

rium carriers. The balance between the electrostatic interaction and the diffusion pressure can result in the formation of the equilibrium spatial structures of the dust grains in the thermal plasma at the atmospheric pressure.

The various patterns of behavior of the nonequilibrium carrier number density at the change of the sign of the potential determine the different diffusion pressures.

In Fig. 5 we can see that the diffusion pressure changes differently at positive and negative potentials. For example, at a distance between planes $d=11r_D$ and at the positive surface potential, the electrostatic pressure is less than the diffusion pressure at any value of the potential, and the planes attract. At the negative surface potential such tendency exists only if $-e\phi_s<0.5T$. However, the increase of the negative potential causes the decrease of the diffusion pressure, and the planes repulse. It means that the equilibrium in the space location of the positive and negative dust grains occurs at different distances. The equilibrium distance for the positively charged grains is less, than for negatively charged grains. Hence, the average distance between the grains in the ordered structures, consisting of the positive grains, should be less, than in the structures, consisting of the negative grains.

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