BOUNDARY INFLUENCE ON GRAIN SCREENING IN SEMI-INFINITE PLASMA

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The screened potential of a single grain embedded into semi-bounded plasma is calculated. The limiting cases of a collisionless plasma and a plasma described in the drift-diffusion approximation are considered. We assume that plasma boundary does not absorb electrons and ions. The grain interaction with plasma particles is described within the "point sink model". According to thus model the singular sinks are introduced into the equations describing plasma dynamics. This makes possible to recover the screened potentials calculated in the case of unbounded plasma and to perform the appropriate calculations in the case of bounded systems. It turns out that the presence of boundary can considerably modify the screened potentials.

Dusty plasmas is a new and interesting field of plasma physics. The main feature of dusty plasmas is that the grains absorbs encountered electrons and ions and thus can accumulate huge electric charge. This introduces many new effects [1].

Usually, in the theoretical studies of dusty plasmas, the system is regarded as an infinite one. But it is obvious that dusty plasmas have finite sizes in the experimental setups and industrial applications. This fact can be important in view of the existence of long-range asymptotes of grain effective potentials [2–4]. This means that the interaction of grains can be influenced by the boundaries even if the grains are located far from it.

Let us consider semi-bounded plasmas ($z > 0$) with a stationary spherical grain at the point $\vec{r}_0 = (0, 0, z_0)$. The domain $z < 0$ is the external medium with the dielectric constant $\tilde{\varepsilon}$. It is assumed that plasma particles are specularly reflected from the boundary, and the grain absorbs all encountered electrons and ions.

We start from the case of a partially ionized plasma which can be described within the drift-diffusion approximation. Weakly ionized gas at atmospheric pressure gives an example of such a plasma [5, 6].

If the ionization and recombination processes can be neglected, the continuity equations for plasma particles are as follows:

$$\vec{\nabla}\vec{\Gamma}_\sigma = -S\delta(\vec{r} - \vec{r}_0), \tag{1}$$

where

$$\vec{\Gamma}_\sigma = -D_\sigma \left(\frac{e_\sigma}{T_\sigma} n_\sigma \vec{\nabla}\phi + \vec{\nabla} n_\sigma \right). \tag{2}$$

On the right-hand part of Eq. (1), we introduced a point sink which describes the electron and ion absorption by a grain. Here, S is the sink intensity, the subscript σ denotes plasma particle species ($\sigma = e$ – electron, $\sigma = i$ – ion), and the rest of notations is traditional. In Eq. (2), the Einstein relation $D_\sigma/\mu_\sigma = T_\sigma/e_\sigma$ is used.

The effective grain electric potential is governed by the Poisson equation

$$\Delta\phi = -4\pi \sum_{\sigma=e,i} e_\sigma n_\sigma - 4\pi q \delta(\vec{r} - \vec{r}_0), \tag{3}$$

where q is the grain charge. The quantities S and q can be determined by solving the appropriate nonlinear boundary-value problem (see, for example, [3, 6] and references therein). They cannot be found within the linear theory presented here. In what follows, we assume that S and q are known.

Linearizing Eqs. (1), (3), one has

$$\frac{e_\sigma n_0}{T_\sigma}\Delta\phi + \Delta\delta n_\sigma = \frac{S}{D_\sigma}\delta(\vec{r} - \vec{r}_0), \tag{4}$$

$$\Delta\phi = -4\pi \sum_{\sigma=e,i} e_\sigma \delta n_\sigma - 4\pi q \delta(\vec{r} - \vec{r}_0), \tag{5}$$

where $\delta n_\sigma(\vec{r}) = n_\sigma(\vec{r}) - n_0$ is the particle density perturbation, and $e_e = -e_i = -e$, where e is the elementary charge.

In the external medium ($z < 0$), the electric potential satisfies the equation

$$\Delta\tilde{\phi} = 0, \tag{6}$$

which has to be solved under the following boundary conditions:

$$\tilde{\phi}(\vec{r}_\perp, z=-0) = \phi(\vec{r}_\perp, z=+0),$$

$$\tilde{\varepsilon}\frac{\partial\tilde{\phi}(\vec{r})}{\partial z}\bigg|_{z=-0} = \frac{\partial\phi(\vec{r})}{\partial z}\bigg|_{z=+0}, \tag{7}$$

$$\vec{\Gamma}_\sigma(\vec{r}_\perp, z=+0) = 0.$$

The last condition means that there is no plasma flux through the boundary.

Eqs. (4), (5) can be solved using the specular continuation of the electric potential and a density perturbation to the region $z < 0$ [7]:

$$\Phi(\vec{r}) = \Phi(\vec{r}_\perp, z) = \begin{cases} \phi(\vec{r}_\perp, z), & z > 0, \\ \phi(\vec{r}_\perp, -z), & z < 0, \end{cases} \tag{8}$$

$$\delta\bar{n}_\sigma(\vec{r}) = \delta\bar{n}_\sigma(\vec{r}_\perp, z) = \begin{cases} \delta n_\sigma(\vec{r}_\perp, z), & z > 0, \\ \delta n_\sigma(\vec{r}_\perp, -z), & z < 0. \end{cases} \tag{9}$$

Such continued quantities Φ and $\delta\bar{n}_\sigma$ satisfy the equations

$$\frac{e_\sigma n_0}{T_\sigma}\Delta\Phi + \Delta\delta\bar{n}_\sigma = \frac{S}{D_\sigma}(\delta(\vec{r}-\vec{r}_0) + \delta(\vec{r}-\vec{r}_0^{\,+})), \tag{10}$$

$$\Delta\Phi = -4\pi\sum_{\sigma=e,i} e_\sigma\delta\bar{n}_\sigma - 4\pi q(\delta(\vec{r}-\vec{r}_0) + \delta(\vec{r}-\vec{r}_0^{\,+})), \tag{11}$$

where $\vec{r}_0^{\,+} = (0,0,-z_0)$.

In the $\vec{k}$-representation, the solution of Eqs. (10), (11) is the following:

$$\Phi_k = \frac{4\pi(q+\tilde{S})\left(e^{-ik_z z_0} + e^{ik_z z_0}\right)}{k^2 + k_{\rm D}^2} -$$

$$-\frac{2\tilde{\varepsilon}k_\perp\phi_{k_\perp}(+0)}{k^2 + k_{\rm D}^2} - \frac{4\pi\tilde{S}}{k^2}\left(e^{-ik_z z_0} + e^{ik_z z_0}\right), \tag{12}$$

where $k_{\rm D}^2 = k_{{\rm D}i}^2 + k_{{\rm D}e}^2, k_{D\sigma}^2 = 4\pi e_\sigma^2 n_\sigma/T_\sigma$,

$$\tilde{S} = \frac{eS}{k_{\rm D}^2}\left(\frac{1}{D_i} - \frac{1}{D_e}\right). \tag{13}$$

The inverse Fourier transformation gives

$$\phi(\vec{r}) = \frac{q+\tilde{S}}{r_-}e^{-k_{\rm D}r_-} + \frac{q+\tilde{S}}{r_+}e^{-k_{\rm D}r_+} - \frac{\tilde{S}}{r_-} - \frac{\tilde{S}}{r_+} -$$

$$-2\tilde{\varepsilon}\tilde{q}\int_0^\infty dk_\perp \frac{k_\perp^2 J_0(k_\perp r_\perp)e^{-z_+\sqrt{k_\perp^2+k_{\rm D}^2}}}{\sqrt{k_\perp^2+k_{\rm D}^2}(\sqrt{k_\perp^2+k_{\rm D}^2} + \tilde{\varepsilon}k_\perp)} +$$

$$+2\tilde{\varepsilon}\tilde{S}\int_0^\infty dk_\perp \frac{k_\perp J_0(k_\perp r_\perp)e^{-z\sqrt{k_\perp^2+k_{\rm D}^2} - z_0 k_\perp}}{\sqrt{k_\perp^2+k_{\rm D}^2} + \tilde{\varepsilon}k_\perp}, \tag{14}$$

where $J_0(k_\perp r_\perp)$ is the Bessel function of zeroth order, $z_\pm = z \pm z_0$, and $r_\pm = \sqrt{r_\perp^2 + z_\pm^2}$.

In the case of unbounded plasma, Eq. (14) is reduced to

$$\phi_u(\vec{r}) = \frac{q+\tilde{S}}{r}e^{-k_{\rm D}r} - \frac{\tilde{S}}{r}. \tag{15}$$

The potential consists of the Debye and Coulomb parts with the appropriate effective charges, this result corresponds to that obtained in [2–4].

Notice that $-\tilde{S}$ can be treated as an effective charge in the unscreened part of the potential. This quantity can be related to the grain charge as $\tilde{S} = -\alpha q$, where α has to be determined by solving the nonlinear boundary-value problem. As was shown in [3], $\alpha \ll 1$ in the case of large grain sizes and $\alpha \le 0.5$ for small grain sizes.

Thus, the first and third terms in Eq. (14) represent the effective potential in unbounded plasma. The second and fourth terms describe the contribution of the charge image. The last two terms (in what follows, I_1 and I_2) are apparently related to the charge induced at the boundary surface.

In the case $r_\perp k_{\rm D}/\tilde{\varepsilon} \gg 1$,

$$I_1 \approx 2\frac{\tilde{q}\tilde{\varepsilon}}{k_{\rm D}^2}\frac{e^{-z_+k_{\rm D}}}{r_\perp^3},$$

and, at $r_\perp k_{\rm D}/\tilde{\varepsilon} \gg 1$ and $z_0 k_{\rm D}/\tilde{\varepsilon} \gg 1$,

$$I_2 \approx 2\frac{\tilde{S}\tilde{\varepsilon}}{k_{\rm D}}\frac{z_0 e^{-zk_{\rm D}}}{(z_0^2 + r_\perp^2)^{3/2}}.$$

In the case $S = 0$, Eq. (14) reduces to the potential of a stationary charge in semi-bounded plasma [7]

$$\phi(\vec{r}) = \frac{q}{r_-}e^{-k_{\rm D}r_-} + \frac{q}{r_+}e^{-k_{\rm D}r_+} -$$

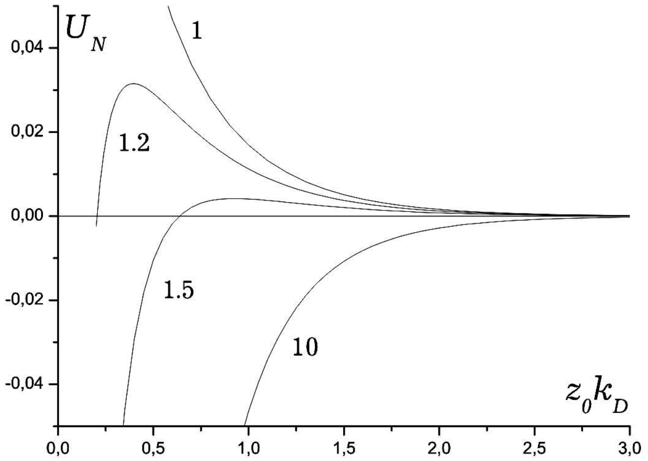

Fig. 1. Dependence of the normalized energy of the grain charge interaction with the induced potential $U_N = U/(q^2 k_\mathrm{D})+1$ on $z_0 k_\mathrm{D}$ and $S = 0$ and various values of $\tilde{\varepsilon}$=1; 1.2; 1.5; 10

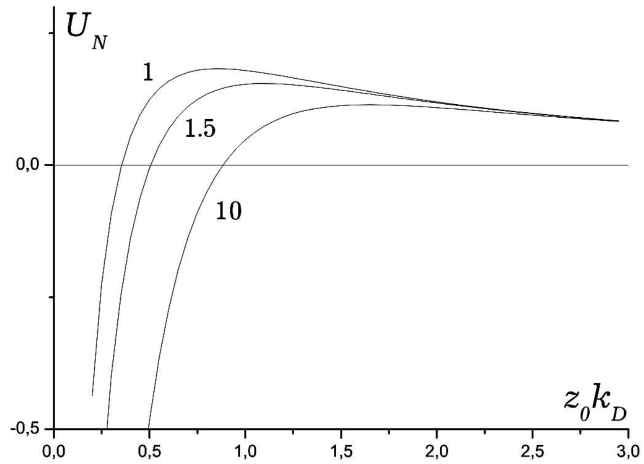

Fig. 2. Dependence of normalized energy of grain charge interaction with induced potential $U_N = U/(q^2 k_\mathrm{D}) + (1-\alpha)$ on $z_0 k_\mathrm{D}$ at $\alpha = 0.5$ and various values of $\tilde{\varepsilon}$=1; 1.5; 10

$$-2\tilde{\varepsilon} q \int_0^\infty dk_\perp \frac{k_\perp^2 J_0(k_\perp r_\perp) e^{-z_+ \sqrt{k_\perp^2 + k_\mathrm{D}^2}}}{\sqrt{k_\perp^2 + k_\mathrm{D}^2}(\sqrt{k_\perp^2 + k_\mathrm{D}^2} + \tilde{\varepsilon} k_\perp)}. \quad (16)$$

The potential energy of the grain charge interaction with the induced potential

$$U(\vec{r}_0) = q \lim_{\vec{r} \to \vec{r}_0} \left(\phi(\vec{r}) - \frac{q}{r_-} \right) \quad (17)$$

is given by

$$\frac{U}{q^2} = -k_\mathrm{D}(1-\alpha) + \frac{1-\alpha}{2z_0} e^{-2k_\mathrm{D} z_0} + \frac{\alpha}{2z_0} -$$

$$-2\tilde{\varepsilon}(1-\alpha) \int_0^\infty dk_\perp \frac{k_\perp^2 e^{-2z_0 \sqrt{k_\perp^2 + k_\mathrm{D}^2}}}{\sqrt{k_\perp^2 + k_\mathrm{D}^2}(\sqrt{k_\perp^2 + k_\mathrm{D}^2} + \tilde{\varepsilon} k_\perp)} -$$

$$-2\tilde{\varepsilon}\alpha \int_0^\infty dk_\perp \frac{k_\perp e^{-z_0(\sqrt{k_\perp^2 + k_\mathrm{D}^2} + k_\perp)}}{\sqrt{k_\perp^2 + k_\mathrm{D}^2} + \tilde{\varepsilon} k_\perp}. \quad (18)$$

At first, we calculate the above-mentioned energy for a charged grain which does not absorb electrons and ions ($S = 0$).

$$\frac{U}{k_\mathrm{D} q^2} + 1 = \frac{e^{-2k_\mathrm{D} z_0}}{2k_\mathrm{D} z_0} - 2\tilde{\varepsilon} \int_0^\infty dx \frac{x^2 e^{-2k_\mathrm{D} z_0 \sqrt{x^2+1}}}{\sqrt{x^2+1}(\sqrt{x^2+1} + \tilde{\varepsilon} x)}. \quad (19)$$

The first term in (19) is positive (the grain is repulsed from the image charge), and the second one is negative (the grain is attracted to the charge induced at the boundary).

The graph of Eq. (19) is shown in Fig. 1. As is seen, the increase of $\tilde{\varepsilon}$ changes the repulsion of the grain from the boundary to the attraction.

Fig. 2 shows the typical behaviour of the normalized energy of the grain charge interaction with the induced potential $U_N = U/(q^2 k_\mathrm{D}) + (1-\alpha)$ as a function of the dimensionless distance from the grain to the boundary $z_0 k_\mathrm{D}$ at various values of $\tilde{\varepsilon}$. As follows from the figure, the potential is nonmonotonic at $z_0 k_\mathrm{D} \sim 1$ for the fixed value of α, which will generate a nonmonotonic density profile for grains near the boundary.

If the collisions between electrons and ions can be neglected, the plasma dynamics is described by the Vlasov equation. In the stationary case in the presence of an immobile grain, such an equation can be written as

$$\left(\vec{v} \frac{\partial}{\partial \vec{r}} - \frac{e_\sigma}{m_\sigma} \nabla \phi(\vec{r}) \frac{\partial}{\partial \vec{v}} \right) f_\sigma(\vec{r}, \vec{v}) = -\delta(\vec{r}) \sigma_\sigma(v) v f_\sigma(\vec{r}, \vec{v}), \quad (20)$$

where $f_\sigma(\vec{r}, \vec{v})$ is the one-particle distribution function.

The singular sink is introduced into the right-hand part of this equation in order to describe the plasma particles absorption by the grain. In the case under consideration, the sink intensity is determined by the charging cross-section $\sigma_\sigma(q, v)$ (see, for example, [8]). Assuming that the perturbations introduced by the sink are small, it is possible to use the linear approximation, i.e. $f_\sigma(\vec{r},\vec{v}) = f_{0\sigma}(v) + \delta f_\sigma(\vec{r},\vec{v})$, where $f_{0\sigma}(v)$ is the Maxwellian distribution.

The boundary condition for $f_\sigma(\vec{r},\vec{v})$ is the following:

$$f_\sigma(\vec{r}_\perp, z=+0, \vec{v}_\perp, \vec{v}_z) = f_\sigma(\vec{r}_\perp, z=+0, \vec{v}_\perp, -\vec{v}_z). \quad (21)$$

It is possible to show that the $\vec{k}$-representation of the specularly extended distribution function perturbation has the form

$$\delta f_{\sigma\vec{k}}(\vec{v}) = -\frac{e_\sigma}{T_\sigma}\Phi_{\vec{k}} f_{0\sigma}(v) +$$

$$+ i\frac{\sigma_\sigma(v) v f_{0\sigma}(v)}{\vec{k}\vec{v} - i0}(e^{-ik_z z_0} + e^{ik_z z_0}), \quad (22)$$

which generates the charge density distribution

$$\delta\rho_{\sigma\vec{k}} = -\frac{e_\sigma^2 n_\sigma}{T_\sigma}\Phi_{\vec{k}} -$$

$$-\frac{2\pi^2}{k} e_\sigma n_\sigma \int dv v^2 \sigma_\sigma(v) f_{0\sigma}(v)(e^{-ik_z z_0} + e^{ik_z z_0}). \quad (23)$$

Substituting Eq. (23) into the Poisson equation for the extended potential, we have

$$\Phi_{\vec{k}} = \frac{4\pi q}{k^2 + k_{\rm D}^2}(e^{-ik_z z_0} + e^{ik_z z_0}) -$$

$$-\frac{8\pi^3(A_i + A_e)}{k(k^2 + k_{\rm D}^2)}(e^{-ik_z z_0} + e^{ik_z z_0}) - \frac{2\tilde{\varepsilon} k_\perp \Phi_{k_\perp}(+0)}{k^2 + k_{\rm D}^2}, \quad (24)$$

where

$$A_\sigma = e_\sigma n_\sigma \int v^2 \sigma_\sigma(v) f_{0\sigma}(v) dv.$$

With regard for the boundary conditions (7),

$$\phi(r) = q\frac{e^{-k_{\rm D} r_-}}{r_-} + q\frac{e^{-k_{\rm D} r_+}}{r_+} - 2\pi(A_i + A_e)\big(F(k_{\rm D} r_-) +$$

$$+ F(k_{\rm D} r_+)\big) - 2\tilde{\varepsilon} q \int_0^\infty dk_\perp \frac{k_\perp^2 J_0(k_\perp r_\perp) e^{-z_+\sqrt{k_\perp^2 + k_{\rm D}^2}}}{\sqrt{k_\perp^2 + k_{\rm D}^2}(\sqrt{k_\perp^2 + k_{\rm D}^2} + \tilde{\varepsilon} k_\perp)} +$$

$$+ 4\pi\tilde{\varepsilon}(A_i + A_e)\int_0^\infty dk_\perp \frac{k_\perp^2 J_0(k_\perp r_\perp) e^{-z\sqrt{k_\perp^2 + k_{\rm D}^2}}}{\sqrt{k_\perp^2 + k_{\rm D}^2} + \tilde{\varepsilon} k_\perp} \times$$

$$\times \int_{-\infty}^{\infty} \frac{dk_z \cos(k_z z_0)}{k(k^2 + k_{\rm D}^2)}, \quad (25)$$

where $F(x) = (e^{-x}{\rm Ei}(x) - e^x {\rm Ei}(-x))/x$, ${\rm Ei}(x)$ is the exponential integral.

As is seen, the structure of potential (25) is similar to that in the case of weakly ionized plasma. Moreover, the difference between Eqs. (14) and (25) vanishes in the case of nonabsorbing grains ($\tilde{S} = 0$ in Eq. (14) and $A_\sigma = 0$ in Eq. (25)). This means that the potential of an external test charge in the semi-bounded plasma does not depend on the plasma dynamics. The difference between Eqs. (14) and (25) is explained by the fact that, in the case of absorbing grains, the system under consideration is an open system, and thus even the stationary characteristics depend on the kinetic coefficients.

Notice that, in the case of unbounded plasma,

$$\phi_u(r) = \frac{q}{r}e^{-k_{\rm D} r} - 2\pi(A_i + A_e)F(k_{\rm D} r). \quad (26)$$

At $k_{\rm D} r \gg 1$,

$$\phi_u(r) \approx -\frac{4\pi(A_i + A_e)}{(r k_{\rm D})^2}. \quad (27)$$

So, the potential has inversely squared asymptote in collisionless plasma.

We use the well-known absorption cross-sections from OML theory [1]

$$\sigma_e(v) = \pi a^2 \begin{cases} 1 - \dfrac{2qe_e}{a m_e v^2}, & v^2 > \dfrac{2qe_e}{a m_e}, \\ 0, & v^2 < \dfrac{2qe_e}{a m_e}, \end{cases}$$

$$\sigma_i(v) = \pi a^2 \left(1 - \frac{2qe_i}{a m_i v^2}\right), \quad (28)$$

where a is the grain radius.

The grain charge q is found from the condition of equality of the electron and ion currents through the grain surface in the stationary state. The equation for dimensionless grain charge $z_g = qe_e/(aT_e)$ [1] is

$$e^{-z_g}\sqrt{t\frac{m_i}{m_e}} = (t + z_g)\frac{n_i}{n_e}, \quad (29)$$

where $t = T_i/T_e$.

In the case of the Maxwellian distribution, using cross-section (28), we get

$$A_i = e_i n_i \frac{a^2}{4}\left(1 + 2\frac{z_g}{t}\right),$$
$$A_e = e_e n_e \frac{a^2}{4}\left(2e^{-z_g}\sqrt{\frac{z_g}{\pi}} + (1-2z_g)(1-\mathrm{erf}(\sqrt{z_g}))\right), \tag{30}$$

where $\mathrm{erf}(x) = \frac{2}{\sqrt{\pi}}\int\limits_0^x e^{-t^2}dt$ is the error function.

For typical values of $t \ll 1$, $A_e \ll A_i$, and Eq. (27) becomes

$$\phi_u(r) \approx \frac{t+2z_g}{4(1+t)}\frac{a^2}{r^2}, \tag{31}$$

which coincides with the result obtained in [9], see also [1].

Energy (17) of the grain interaction with the induced potential is

$$U = -q^2 k_{\mathrm{D}} + \frac{e^{-2k_{\mathrm{D}}z_0}}{2z_0} - 2\pi(A_i + A_e)\big(F(k_{\mathrm{D}}a)+$$

$$+F(2k_{\mathrm{D}}z_0)\big) - 2\tilde{\varepsilon}q^2\int\limits_0^\infty dk_\perp \frac{k_\perp^2 e^{-2z_0\sqrt{k_\perp^2+k_{\mathrm{D}}^2}}}{\sqrt{k_\perp^2+k_{\mathrm{D}}^2}(\sqrt{k_\perp^2+k_{\mathrm{D}}^2}+\tilde{\varepsilon}k_\perp)} +$$

$$+4\pi\tilde{\varepsilon}(A_i+A_e)\int\limits_0^\infty dk_\perp \frac{k_\perp^2 e^{-z_0\sqrt{k_\perp^2+k_{\mathrm{D}}^2}}}{\sqrt{k_\perp^2+k_{\mathrm{D}}^2}+\tilde{\varepsilon}k_\perp}\int\limits_{-\infty}^{\infty}\frac{dk_z\cos(k_z z_0)}{k(k^2+k_{\mathrm{D}}^2)}. \tag{32}$$

In order to plot the graph of (32), one needs to calculate firstly z_g from (29) and A_i, A_e. Figure 3 presents the normalized energy $U_N = U/(q^2k_{\mathrm{D}}) + 1 + 2\pi(A_i+A_e)F(k_{\mathrm{D}}a)$ for a single grain with dimensionless radius $ak_{\mathrm{D}} = 0.1$ in argon with $t = 0.05$.

The main feature of Figs. 2 and 3 in comparison with Fig. 1 is the presence of the long-range repulsion of a grain from the boundary at any values of $\tilde{\varepsilon}$.

We calculated the effective potential of a stationary grain absorbing electrons and ions in the case of semi-bounded plasma. It is shown that the plasma particles absorption by the grain leads to a considerable deviation of the effective potential from that obtained in the case of a nonabsorbing charged grain. The influence of the plasma dynamics on the obtained results is studied as well. The behavior of the charge induced by a grain is quantitatively different for collisionless and collisional plasmas. This shows that the equilibrium grain distribution near the plasma boundary can be inhomogeneous, and the stationary grain density profiles can be influenced by the plasma parameters.

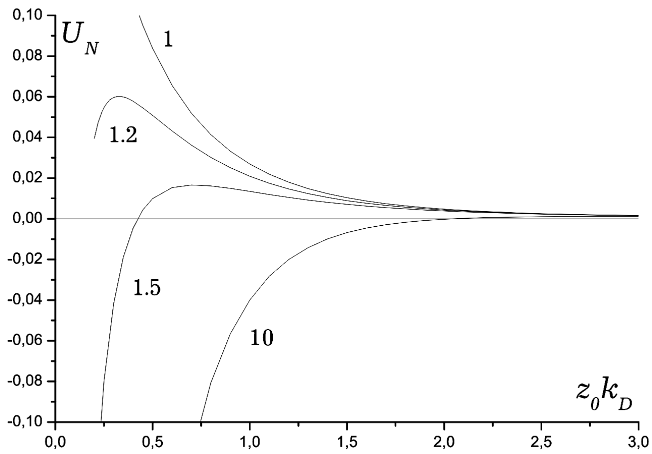

Fig. 3. Dependence of the normalized energy of grain charge interaction with the induced potential on z_0k_{D} at $ak_{\mathrm{D}} = 0.1$, $t = 0.05$ at various values of $\tilde{\varepsilon}$=1; 1.2; 1.5; 10

1. V.E. Fortov, A.V. Ivlev, S.A. Khrapak, A.G. Khrapak, G.E. Morfill, Phys. Repts. **421**, 1 (2005).
2. A.G. Zagorodny, A.I. Momot, Ukr. J. Phys. **51**, 569 (2006).
3. O. Bystrenko, A. Zagorodny, Phys. Rev. **E 67**, 066403 (2003).
4. A.V. Filippov, A.G. Zagorodny, A.F. Pal', A.N. Starostin, JETP Lett. **81**, 146 (2005).
5. V.V. Ivanov, A.F. Pal', T.V. Rakhimova, A.O. Serov, N.V. Suetin, Zh. Eksp. Teor. Fiz. **115**, 2020 (1999).
6. A.F. Pal', A.O. Serov, A.N. Starostin, A.V. Filippov, V.E. Fortov, Zh. Eksp. Teor. Fiz. **119**, 272 (2001).
7. R. Tageback, A.S. Usenko, I.P. Yakimenko, A.G. Zagorodny, J. Plasma Phys. **18**, 113 (1977).
8. A.V. Filippov, A.G. Zagorodny, A.I. Momot, A.F. Pal', A.F. Starostin, Zh. Eksp. Teor. Fiz. **130**, 1 (2006).
9. V.N. Tsytovich, Yu.K. Khodatayev, R. Bingham, Comments Plasma Phys. Controlled Fusion **17**, 249 (1996).

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ВПЛИВ МЕЖІ НА ЕКРАНУВАННЯ ПОРОШИНКИ
У НАПІВНЕСКІНЧЕННІЙ ПЛАЗМІ

А.Г. Загородній, А.І. Момот

Р е з ю м е

Розраховано екрановані потенціали порошинки у напівобмеженій плазмі. Розглянуто граничні випадки плазми без зіткнень і плазми, що описується у дрейфово-дифузійному наближенні. Вважається, що межа плазми не поглинає плазмові частинки. Ми вважаємо порошинку точковою. В рівняння, які описують динаміку плазми, вводяться відповідні точкові стоки, що дає змогу врахувати поглинання плазмових частинок порошинкою. Такий модельний підхід дозволяє відтворити результати для екранованих потенціалів у необмеженій плазмі і легко може бути узагальнений на випадок обмеженої системи. Виявляється, що динаміка плазми суттєво впливає на асимптотичну поведінку екрануючого потенціалу. Детально вивчено модифікацію цих потенціалів за рахунок межі.