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Brownian motion of grains and negative friction in dusty plasmas

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Within the approximation of dominant charging collisions the explicit microscopic calculations of the Fokker-Planck kinetic coefficients for highly-charged grains moving in plasma are performed. It is shown that due to ion absorption by grain the friction coefficient can be negative and thus the appropriate threshold value of the grain charge is found. The stationary solutions of the Fokker-Planck equation with the velocity dependent kinetic coefficient are obtained and a considerable deviation of such solutions from the Maxwellian distribution is established.

Key words: *dusty plasma, negative friction, effective temperature, Fokker-Planck equation*

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Description of the Brownian motion in the systems with particle or energy fluxes still remains one of the key problems in the statistical physics of the open systems [1]. Starting from the classical Lord Rayleigh work [2] many studies of the non-equilibrium motion of Brownian particles with additional (inner or external) energy supply have been performed. In particular, such studies are of great importance for physical-chemical [3,4] and biological [5] systems in which non-equilibrium Brownian particle motion is referred to as the motion of active Brownian particles. Recently the dynamical and energetic aspects of motion for the active Brownian particles have been described based on the Langevin equation and the appropriate Fokker-Planck equation [6,7]. The possibility of negative friction (negative values of friction coefficient) of the Brownian particles was regarded as a result of energy pumping.

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For some phenomenological dependences of the friction coefficient as a function of the grain's velocity, one-particle stationary non-Maxwellian distribution function was found.

The traditional formulations of the non-equilibrium Brownian motion are based on some phenomenological expressions for the friction and diffusion coefficients. In particular, it means that deviations from the Einstein relation, as well as the velocity dependence of these coefficients are postulated and high level of uncertainty for the application of such models to the real systems takes place. Moreover, since in the case of the open systems there exist few types of the Fokker-Planck equation which can be related to the nonlinear Langevin equations (see, for example [1] and references cited therein) the form of the Fokker-Planck equation itself could be a matter of the choice. Here we will consider another situation, when the kinetic coefficients can be calculated exactly on the basis of the microscopically derived Fokker-Planck equation for dusty plasmas [8,9]. It will be shown that in the case of strong interaction parameter $\Gamma \equiv e^2 Z_{\rm g} Z_{\rm i}/aT_{\rm i} \gg 1$ (here $Z_{\rm g}$, $Z_{\rm i}$ are the charge numbers for the grains and ions, respectively, a is the grain radius, $T_{\rm i}$ is the ion temperature) the negative friction coefficient appears for some velocity domain. If the charging collisions are dominant, this domain is determined by the inequalities: $v^2 < 10v_{\rm Ti}^2(\Gamma-1)/(3\Gamma-1)$ for $\Gamma \geqslant 1$, but $(\Gamma-1) \ll 1$,$(v_{\rm Ti}^2 \equiv T_{\rm i}/m_{\rm i})$; and $v^2 < 2v_{\rm Ti}^2\Gamma$ for $\Gamma \gg 1$. The physical reason for manifestation of negative friction is clear: the cross-section for ion absorption by grain increases with the relative velocity between the ion and grain decreases due to the charge-dependent part of the cross-section. Therefore, for a moving highly-charged grain ($\Gamma \gg 1$) the total momentum transfer from ions to the grain in the direction of grain velocity could be larger than in the opposite direction.

Naturally, the Coulomb scattering and particle friction related to the ion-grain and neutral-grain elastic collisions will increase the threshold for negative friction and even can suppress it for some plasma parameters. However, the latter are rather sophisticated processes, and their theoretical description within various approximations can shadow the physics of the phenomenon. Some estimate of the effect of such processes will be also done below, but the detailed description of their action should be a matter of further consideration. On the other hand, the problem of microscopic investigation of negative friction in dusty plasmas is so fundamental that it deserves a description of its simplest manifestation (which occurs in the case of dominant charging collisions) to be considered in this paper.

We start from the Fokker-Planck equation for the spherical grains in dusty plasma with typical narrow charge distribution around a negative value $q = e_{\rm e} Z_{\rm g}$, which permits to put all grain charges to be equal. We also ignore the increase of the grain mass [10, 11] assuming that neutral atoms generated in the course of the surface electron-ion recombination escape from the grain surface into a plasma. As it was mentioned already in [9], where the subsequent kinetic theory with ion absorption in dusty plasmas has been developed, the problem of correct momentum transfer due to ion absorption and further surface recombination is important for description of stationary state. Strictly speaking the stationary state can be reached only by inclusion of these both processes. To avoid this complication, however, usually it is suggested that the mass of grain should be constant, due to the big difference of ion and grain masses. Recently in [12] it was shown that this assumption leads to the violation of the Galiley invariance of the particle transition probability generated by collisions. Nevertheless, the correct kinetic theory of ion absorption with further surface ion recombination is still in the stage of development and will be published elsewhere.

Then, for the conditions $l_{\rm i} \gg \lambda_{\rm D}$, a, where $l_{\rm i}$ is the ion mean free path length and $\lambda_{\rm D}$ is the plasma screening (Debye) length, the friction and diffusion coefficients in the Fokker-Planck kinetic equation $\beta(q, \mathbf{v})$ and $D(q, \mathbf{v})$ are given by [9]:

$$\begin{aligned}
\beta(q,\mathbf{v}) &= -\sum_{\alpha=e,i} \frac{m_\alpha}{m_{\rm g}} \int d\mathbf{v}' \frac{\mathbf{v}\cdot\mathbf{v}'}{v^2} \sigma_{\alpha{\rm g}}(q, |\mathbf{v}-\mathbf{v}'|)|\mathbf{v}-\mathbf{v}'| f_\alpha(\mathbf{r},\mathbf{v}',t), \\
D(q,\mathbf{v}) &= \sum_{\alpha=e,i} \frac{1}{2}\left(\frac{m_\alpha}{m_{\rm g}}\right)^2 \int d\mathbf{v}' \frac{(\mathbf{v}\cdot\mathbf{v}')^2}{v^2} \sigma_{\alpha{\rm g}}(q, |\mathbf{v}-\mathbf{v}'|)|\mathbf{v}-\mathbf{v}'| f_\alpha(\mathbf{r},\mathbf{v}',t), \\
\sigma_{g\alpha}(q,v) &= \pi a^2 \left(1 - \frac{2e_\alpha q}{m_\alpha v^2 a}\right) \theta\left(1 - \frac{2e_\alpha q}{m_\alpha v^2 a}\right). \qquad (1)
\end{aligned}$$

Here $\sigma_{g\alpha}(q, \mathbf{v})$ is the cross-section for grain charging within the orbital motion limited (OML) theory. In this approximation all electrons and ions approaching the grain on the distance smaller than a are assumed to be absorbed. Subscript $\alpha = e, i$ labels plasma particle species. The rest of notation is traditional. To include the processes of the electron, ion and atom scattering we have to summarize the appropriate coefficients on the different type of the processes. In the case of dominant charging collisions integration in equation (1) can be performed explicitly. The ion part of β (which exceeds the electron one at least in $(m_{\rm i}T_{\rm e}/m_{\rm e}T_{\rm i})^{1/2}$ times) is as follows:

$$\begin{aligned}
\beta_{\rm i}(q,\mathbf{v}) &= -\sqrt{2\pi}\frac{m_{\rm i}}{m_{\rm g}} a^2 n_{\rm i} v_{\rm Ti} \left[I_1(\eta) + I_2(\eta,\Gamma)\right], \\
I_1(\eta) &= \left(-\frac{1}{2} + \frac{1}{4\eta}\right)\sqrt{\frac{\pi}{\eta}}{\rm Erf}\sqrt{\eta} - \frac{1}{2\sqrt{\eta}} e^{-\eta}, \\
I_2(\eta,\Gamma) &= \frac{\Gamma}{\eta}\left[\frac{1}{2}\sqrt{\frac{\pi}{\eta}}{\rm Erf}\sqrt{\eta} - e^{-\eta}\right]. \qquad (2)
\end{aligned}$$

Here n_a are the densities of the electrons and ions, ${\rm Erf}\sqrt{\eta}$ is the error function and $\eta \equiv \eta(v) = v^2/2v_{\rm Ti}^2$. It is easy to see that the terms $I_1(\eta)$ and $I_2(\eta, \Gamma)$ describe the parts of $\beta_{\rm i}(q, \mathbf{v})$ related to purely geometrical and charge-dependent collecting cross-sections, respectively. The integration for $D_{\rm i}(q, \mathbf{v})$ leads to the expression:

$$\begin{aligned}
D(q,\mathbf{v}) &= \frac{4}{3}\sqrt{2\pi}\left(\frac{m_{\rm i}}{m_{\rm g}}\right) a^2 n_{\rm i} v_{\rm Ti}\left(\frac{T_{\rm i}}{m_{\rm g}}\right)\Big[K_1(\eta) + K_2(\eta,\Gamma)\Big], \\
K_1(\eta) &= \frac{3}{16\eta}\left[2(\eta-1)e^{-\eta} + (2\eta^2 + \eta + 1)\sqrt{\frac{\pi}{\eta}}{\rm Erf}\sqrt{\eta}\right],
\end{aligned}$$

$$K_2(\eta,\Gamma) \;=\; \frac{3\Gamma}{8\eta}(\eta+1)\left[-2e^{-\eta}+\sqrt{\frac{\pi}{\eta}}\mathrm{Erf}\sqrt{\eta}\right]. \tag{3}$$

As follows from the physical reason and directly from equation (3), the coefficient $D(q,\mathbf{v})$ is always positive. At the same time for some values of η and Γ the friction coefficient $\beta_{\mathrm{i}}(q,\mathbf{v})$ can be negative. To find the root $\eta(\Gamma)$ of the equation $\beta_{\mathrm{i}}(\eta,\Gamma)=0$ let us consider two limiting cases $\eta \ll 1$ and $\eta \gg 1$. For $\eta \ll 1$ equation (2) gives

$$\beta_{\mathrm{i}}(\eta,\Gamma) = 2A\left[1-\Gamma-\frac{\eta}{5}(1-3\Gamma)\right], \tag{4}$$

where

$$A = \frac{1}{3}\sqrt{2\pi}\left(\frac{m_{\mathrm{i}}}{m_{\mathrm{g}}}\right)a^2 n_{\mathrm{i}} v_{\mathrm{Ti}}.$$

Equation (4) has a root $\eta(\Gamma)$, which exists and is small (according to the conditions of applicability for this expansion) only for $\Gamma > 1$, but $\Gamma - 1 \ll 1$:

$$\eta(\Gamma) = 5\frac{(\Gamma-1)}{3\Gamma-1} \simeq \frac{5}{2}(\Gamma-1). \tag{5}$$

Therefore the function $\beta_{\mathrm{i}}(\eta,\Gamma)$ is negative at $\eta < \eta(\Gamma)$, if $\Gamma > 1$. Equation (4) shows that for $\Gamma < 1$, when β_{i} is positive for all η, the derivative $(\mathrm{d}\beta_{\mathrm{i}}(\eta,\Gamma)/\mathrm{d}\eta)|_{\eta=0}$ changes its sign at $\Gamma = 1/3$. Near $\eta = 0$ the friction β_{i} decreases as function η if $\Gamma < 1/3$ and increases if $\Gamma > 1/3$.

For $\eta \gg 1$ and arbitrary values of Γ equation (2) gives

$$\beta_{\mathrm{i}}(\eta,\Gamma) = \frac{3}{2}A\sqrt{\frac{\pi}{\eta}}\left[1-\frac{1+2\Gamma}{2\eta}\right]. \tag{6}$$

It means that for $\Gamma \gg 1$ and large η the equation $\beta_{\mathrm{i}}(\eta_1\Gamma) = 0$ has a root $\eta_1(\Gamma) \simeq (1+2\Gamma)/2 \sim \Gamma$. Naturally, for all $\Gamma > 1$ there exists appropriate $\eta(\Gamma)$, which is the root of that equation.

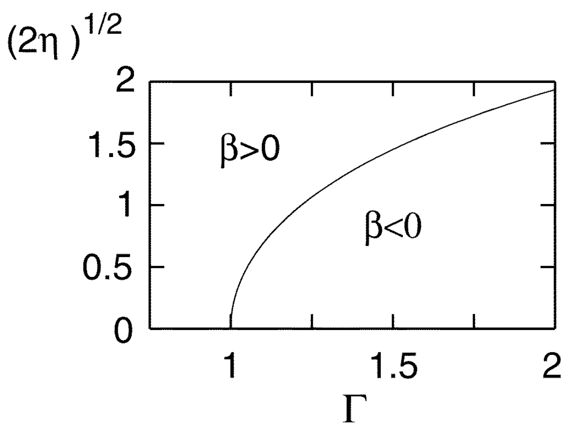

Figure 1. Numerical solution of the equation $\beta(\eta,\Gamma)=0$ separating the positive and negative values of the friction coefficient.

This conclusion is confirmed by the exact numerical calculations of $\eta(\Gamma)$ (figure 1). Asymptotically (for $\eta \gg \max(1,\Gamma)$) $\beta_{\mathrm{i}}(\eta,\Gamma)$ tends to zero as $\sqrt{\eta^{-1}}$ and is positive for $\eta > \eta(\Gamma)$. For the case of the negative friction ($\Gamma > 1$) the maximum of the coefficient $\beta_{\mathrm{i}}(\eta,\Gamma)$ is located at the point $\eta_{\mathrm{m}}(\Gamma) \gg 1$:

$$\eta_{\mathrm{m}} \simeq \frac{3}{2}(1+2\Gamma), \qquad \beta_{\mathrm{i}}(\eta_{\mathrm{m}},\Gamma) \simeq A\sqrt{\frac{2\pi}{3(1+2\Gamma)}}. \tag{7}$$

For the diffusion coefficient $D_{\rm i}(\eta,\Gamma)$ the expansions for $\eta \ll 1$ and $\eta \gg 1$ lead to

$$D_{\rm i}(\eta,\Gamma) = 4A\left(\frac{T_{\rm i}}{m_{\rm g}}\right)\left[1+\frac{\Gamma}{2}+\frac{\eta}{10}(1+2\Gamma)\right], \qquad \eta \ll 1$$

$$D_{\rm i}(\eta,\Gamma) = 4A\left(\frac{T_{\rm i}}{m_{\rm g}}\right)\frac{3}{8}(\pi\eta)^{1/2}\left[1+\frac{1}{2\eta}(1+2\Gamma)\right], \qquad \eta \gg 1. \tag{8}$$

The typical behaviour of $\beta_{\rm i}(\eta,\Gamma)$ and $D_{\rm i}(\eta,\Gamma)$ calculated numerically based on the equations (2), (3) is shown in figures 2, 3.

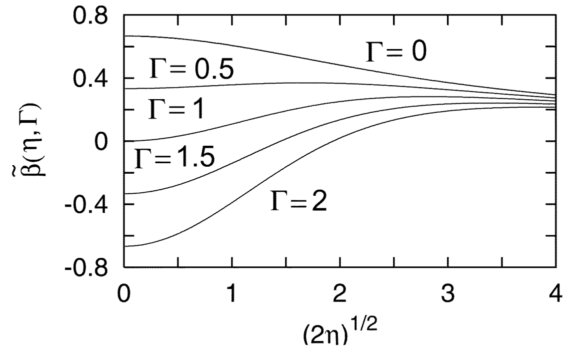

Figure 2. The velocity dependences of dimensionless friction coefficient $\widetilde{\beta}(\eta,\Gamma) = I_1(\eta) + I_2(\eta,\Gamma)$ for different values of Γ.

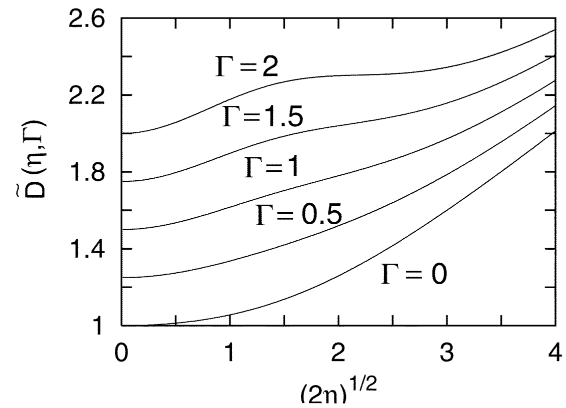

Figure 3. The same dependences for the dimensionless diffusion coefficient $\widetilde{D}(q,v) = K_1(\eta) + K_2(\eta,\Gamma)$.

The stationary solution of the Fokker-Planck equation (which is of the Itoh's form) with the kinetic coefficients (2), (3) for the grain distribution function $f_{\rm g}(q,v)$ is

$$f_{\rm g}(q,v) = \frac{C}{D_{\rm i}(q,v)}\exp\left[-\int\limits_0^v {\rm d}vv\frac{\beta_{\rm i}(q,v)}{D_{\rm i}(q,v)}\right], \tag{9}$$

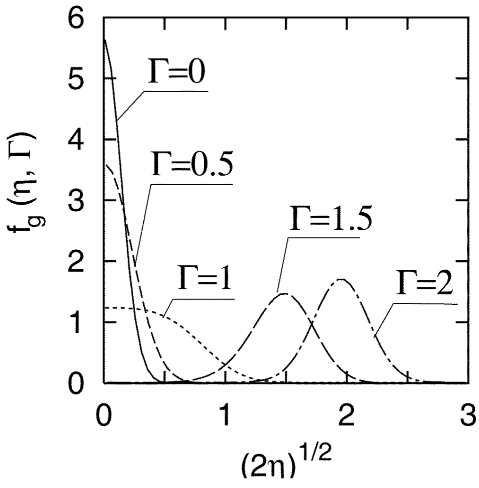

Figure 4. The same dependences for the distribution function $f_{\mathrm{g}}(\eta,\Gamma)$.

where C is a constant, providing normalization $\int \mathrm{d}\mathbf{v} f_{\mathrm{g}}(q,v)=1$. The velocity dependence of this solution for different values of Γ is shown in figure 4.

In order to get some analytical estimates let us consider the vicinity of the point $\Gamma=1$. In such a case the integration in equation (9) leads to the non-Maxwellian distribution function, which for $\Gamma>1$ possesses a maximum at $v\neq 0$:

$$\begin{aligned} f_{\mathrm{g}}(q,\mathbf{v}) &= \left(\frac{m_{\mathrm{i}}}{2\pi T_{\mathrm{i}}}\right)^{3/2} \frac{\sqrt{\pi}}{2Y_{3/2}(1+\Gamma/2)}\left[1+\frac{m_{\mathrm{i}}v^2}{2T_{\mathrm{i}}^*}\frac{1}{5}(1+2\Gamma)\right]^{-1} \\ &\quad\times \exp\left\{-\frac{m_{\mathrm{g}}v^2}{2T_{\mathrm{i}}^*}\left[(1-\Gamma)+\frac{m_{\mathrm{i}}v^2}{2T_{\mathrm{i}}^*}\frac{1}{10}\left(5\Gamma^2+4\Gamma-3\right)\right]\right\}. \end{aligned} \tag{10}$$

Here

$$\begin{aligned} Y_\nu(\Gamma) &= \int\limits_0^\infty \mathrm{d}\eta\,\eta^{\nu-1}\exp\left[-\alpha\eta(1-\Gamma)-\gamma\eta^2\right] = \\ &= (2\gamma)^{-\frac{\nu}{2}}\frac{\sqrt{\pi}}{2}\exp\left[\frac{\alpha^2(1-\Gamma)^2}{8\gamma}\right]D_{-\nu}\left(\frac{\alpha(1-\Gamma)}{\sqrt{2\gamma}}\right), \\ T_{\mathrm{i}}^* &= 2T_{\mathrm{i}}\left(1+\frac{\Gamma}{2}\right), \qquad \alpha=\frac{m_{\mathrm{g}}}{2m_{\mathrm{i}}}\left(1+\frac{\Gamma}{2}\right), \\ \gamma &= \frac{\alpha(5\Gamma^2+4\Gamma-3)}{10(2+\Gamma)} \end{aligned} \tag{11}$$

and D_ν is the cylindrical parabolic function. Equation (10) is relevant, if the integral over η converges, which is valid for positive γ $(\Gamma>0,472)$. At the same time, as was mentioned above, for the applicability of the expansions (4), (8) the inequality $\Gamma<1+\delta$ with $\delta\ll 1$ is required. It is clear from equations (6), (8), that the asymptotic behaviour of the distribution function is non-exponential $f_{\mathrm{g}}(\eta)|_{\eta\to\infty}\sim\eta^{-(m_{\mathrm{g}}/m_{\mathrm{i}})}$, but for the values of Γ under consideration it is not essential for calculations of the averages due to the rapid convergence of the integrals over η. In particular, the average kinetic energy K of grains is

$$K(\Gamma) = \left(\frac{m_{\rm g}}{m_{\rm i}}\right) T_{\rm i} \frac{Y_{5/2}(\Gamma)}{Y_{3/2}(\Gamma)}\,, \tag{12}$$

that gives $K = 2,86(m_{\rm g}/m_{\rm i})^{1/2} T_{\rm i}$ for $\Gamma = 1$. Thus, the ion absorption by grains can lead to grain heating and grain average kinetic energy (or their effective temperature) can be much higher than the electron and ion temperatures. This effect could be treated as the one giving qualitative explanation of the experimental data [13, 14], as it was already suggested in [8,9] based on the velocity independent approximation for β and D. As it was shown in [7,8] for some plasma parameters the effective temperature became negative, which has no physical sense. The same statement was repeated in the recent paper [15], dealing with the hydrodynamic consideration of anomalous heating of grains in dusty plasmas. As is clear from the present paper and was shown already earlier (see e.g. [16]), the effective temperature, as well as the average kinetic energy of grain, calculated consequently based on the theory developed in [8,9], with velocity dependent friction coefficient, is always positive, as it follows from equations (10)–(12). At the same time the statement on the possibility of anomalous temperature (or more exactly average kinetic energy) of grains as the result of momentum transfer due to ion absorption is valid in the model under consideration. The final conclusion about the anomalous heating due to the ion absorption can be done based on the theory, taking into account surface ion recombination and regeneration of neutral atoms.

Let us consider now the case $\Gamma \gg 1$. In this case the ratio $Q = (T_{\rm i}/m_{\rm i})\beta(\eta,\Gamma)/D(\eta,\Gamma)$ can be represented for all values of η with a good accuracy by the function $Q = (m_{\rm g}/m_{\rm i})(\Gamma-\eta)/[(1+\Gamma)(\eta+\Gamma)]$ that gives

$$f_{\rm g} = \frac{C}{D(\eta,\Gamma)} \exp\left\{\frac{m_{\rm g}}{m_{\rm i}}\left[\frac{(3\Gamma+1)}{(2\Gamma-1)}\ln\frac{\Gamma(\eta+1)}{\Gamma+\eta} - \frac{1}{2}\ln\frac{\eta^2+(\Gamma+1)\eta+\Gamma}{\Gamma}\right]\right\}, \tag{13}$$

where

$$D_{\rm i}(\eta,\Gamma)|_{\Gamma\gg 1} \simeq \frac{3}{2}A\frac{T_{\rm i}}{m_{\rm g}}\sqrt{\pi\eta}\left[1+\frac{\Gamma}{(\eta+\frac{3}{4}\sqrt{\pi\eta})}\right]. \tag{14}$$

It is clear that $f_{\rm g}(\eta,\Gamma)$ given by equation (13) is non-exponential.

Finally for the domain $\Gamma < 0,472$ we can omit the term $\sim v^4$ in equation (10) and the distribution $f_{\rm g}(q,v)$ becomes Maxwellian with the effective grain temperature:

$$T_{\rm eff} = \frac{2T_{\rm i}^*}{1-\Gamma} = \frac{2T_{\rm i}(1+\frac{\Gamma}{2})}{1-\Gamma}. \tag{15}$$

Equation (15) is similar to equation (69) for the effective temperature, which was found in [8]. At the same time equation (15) is valid only for $\Gamma < 0,472$, where the effective temperature is positive. The above conclusions are in a good agreement with the results of numerical calculations of $f_{\rm g}(\eta,\Gamma)$ based on the equation (9) (figure 4).

If we take into account the processes of atom-grain and ion-grain elastic scattering the coefficients $\beta(\eta,\Gamma)$ and $D(\eta,\Gamma)$ should include additional terms. To estimate, for example, contribution of atom-grain and ion-grain scattering we can consider $\beta(\eta)$ calculated based on the equation (2) with the appropriate transport cross-sections. For the case $\eta \ll 1$ the friction coefficient is

$$\beta_{\rm i}(\eta,\Gamma) = 2A\left[1-\Gamma+4\frac{n_a}{n_{\rm i}}\left(\frac{T_a m_a}{T_{\rm i} m_{\rm i}}\right)^{1/2}-\frac{\eta}{5}(1-3\Gamma)+\Gamma^2\ln\Lambda_{\rm i}\right]. \qquad (16)$$

The negative friction can exist for small η, if the Coulomb scattering is strongly suppressed, when the Coulomb logarithm $\ln\Lambda_{\rm i}$ is small [17–19], which is typical of strong interaction. The root of the function $\beta(\eta,\Gamma)$ is shifted to the region of large, Γ, which is determined by the atom density n_a. The result of rough estimate for the case $\lambda_{\rm Li} \geqslant \lambda_{\rm D} > a$ ($\lambda_{\rm Li}$ is the Landau length) gives

$$\eta(\Gamma) = \frac{5}{3\Gamma-1}\left[\Gamma-1-3\sqrt{\frac{Z_{\rm g}\Gamma}{\pi n_{\rm i} a^3}}-\frac{4n_a}{n_{\rm i}}\left(\frac{T_a m_a}{T_{\rm i} m_{\rm i}}\right)^{1/2}\right]. \qquad (17)$$

It is necessary to point out that completely positive Coulomb logarithm for dusty plasma (for arbitrary relations between the values $\lambda_{\rm D}, \lambda_{\rm Li}, a$) has been introduced in [20]. For $a \ll \lambda_{\rm D}$ the respective logarithm was found in [18].In general, a more detailed consideration of scattering processes and the effects of strong interactions between the grains [19], is needed for the exact description of the region of negative friction. Of course, the realization of the negative value of the total friction coefficient in the physical experiment requires special conditions, since other mechanisms of positive friction exist. At the same time at the presence of ion flow there is an effective mechanism of negative friction due to the ion scattering by grains [21], which can be very important for dusty plasmas even when ion absorption is small.

In conclusion, ion absorption by grains can generate negative friction and provide a substantial increase of the average grain kinetic energy in comparison with the temperatures of the other plasma components. Microscopical justification of negative friction on the kinetic level is presented in the model without surface ion recombination and thus deviation of the grain distribution function from the Maxwellian distribution is found. Further development of the theory with regard to surface recombination is needed in order to answer the question whether the conditions for negative friction in real dusty plasmas do exist. For dusty plasma, as for an open system, the fluctuation-dissipation theorem in the form of Einstein relation is not applicable. Finally, we would like to point out once more that negative friction is a widely spread phenomenon in physics, chemistry and biology, which is usually described phenomenologically. Therefore microscopical approach to this problem for specific models seems to be important and useful.

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Броунівський рух зернин і від'ємне тертя у запорошеній плазмі

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В наближенні домінуючих контактних зіткнень порошинок з частинками плазми виконано макроскопічні розрахунки кінетичних коефіцієнтів Фокера-Планка для заряджених порошинок, що рухаються у плазмі. Показано, що завдяки зіткненням з йонами коефіцієнт тертя може бути від'ємним, і знайдено порогове значення заряду порошинки. Знайдено розв'язок стаціонарного рівняння Фокера-Планка з кінетичними коефіцієнтами, що залежать від швидкості, і встановлено суттєву відмінність отриманих розв'язків від максвелових розподілів.

Ключові слова: *запорошена плазма, від'ємне тертя, ефективна температура, рівняння Фокера-Планка*

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