Brownian dynamic of laser cooling and crystallization of electron-ion plasma

A. P. Gavriliuk,[1] I. L. Isaev,[2] S. V. Karpov,[2,*] I. V. Krasnov,[1] and N. Ya. Shaparev[1]
[1]*Institute of Computational Modeling, Russian Academy of Sciences, Krasnoyarsk, Russia*
[2]*L.V. Kirenskiy Institute of Physics, Russian Academy of Sciences, Krasnoyarsk, Russia*
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Laser cooling and crystallization of electron-ion plasma is studied using the Brownian dynamics simulation technique and taking into consideration the interaction of ions with the electron subsystem. It has been shown that the nonlinear dependence of laser friction force on the velocity of ions has to be taken into account in order to simulate in an adequate manner the cooling dynamics and obtain a correct estimate for minimum temperatures. It has been found that times required for formation of an ordered ionic structure can be much longer than the typical plasma cooling time.

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I. INTRODUCTION

Obtaining of a strongly nonideal plasma and investigation of its properties has been the subject of a research interest since fairly long ago [1,2]. Electron-ion ultracold plasma (UP) with temperatures of charged particles ≤ 100 K and concentrations $n \leq 10^{10}$ cm^{-3} is an interesting object of investigation. Such plasma was first obtained in [3–5] by means of a near-threshold photoionization of cold atoms. In earlier papers [6,7], laser cooling was suggested for obtaining strongly nonideal ions in electron-ion plasma. Those papers also showed that the conditions for Wigner crystallization of the ion subsystem could be achieved in plasma. Papers [3–5] encouraged a further study of UP properties [8–15], which revealed, in particular, that the strong nonideality required for initiation of crystallization of an electron or ion subsystem could not be achieved due to the fast relaxation of electrons and ions into equilibrium energy distribution followed by a three-body recombination. This finding had motivated the authors of [11–14] to employ the molecular-dynamics methods for description laser cooling of quasineutral plasma. Such cooling resulted in the formation of a quasicrystalline ionic structure in the shape of coaxial spheres. While the results of those papers were quite successful there were still questions left as to the adequacy of description of the laser friction force and electron-ion energy exchange. In particular, in the first papers [6,7] as well as in the later ones [11–16], coefficient of laser friction was used that was independent of the ion velocity. A more precise description of the laser friction force is generally required which takes into consideration the nonlinear dependence of the force on the velocity of ions. For low energy particles, the role of electron-ion energy exchange can be significant even in a low-density plasma as the cross section of elastic collisions of charged particles is $\sigma_{\text{el}} \sim 1/\varepsilon_{\text{kin}}^2$, where ε_{kin} is the electron kinetic energy, which is not high in UP. And, as was shown in [6,7,11–14], the electron-ion energy exchange is one of the key factors limiting the lowest achievable temperature of ions under laser cooling. So a correct description of this energy exchange is important for simulation of plasma behavior under laser cooling.

The primary goals of the paper are the modification of well-known model of one-component plasma [1] by taking into account the thermal interactions of negative electron background with ions and carrying out of the following investigations by means of this modified model: (a) regularities of cooling and crystallization of electron-ion plasma in the field of laser radiation; (b) the effect of nonlinear dependence of laser friction force on the velocity of ions and the influence of such dependence on the dynamics of cooling and crystallization of plasma.

In Sec. III we present the results of simulations for beryllium plasma in which as our estimates showed [16] the maximal value of the parameter of nonideality can be achieved as well as the most pronounced electron-ion heat exchange.

II. MODEL OF LASER COOLED PLASMA

In this paper we deal with a quasineutral electron-ion plasma having a weakly nonideal electron subsystem, the ion subsystem however can be strongly nonideal

$$\Gamma_\alpha = \frac{q_\alpha^2}{4\pi\varepsilon_0 a k_{\text{B}} T_\alpha}, \quad \alpha = \text{e,i}$$

$$\Gamma_{\text{e}} \ll 1, \quad \Gamma_{\text{i}} \geq 1, \quad 4\pi a^3 n = 1, \tag{1}$$

where Γ_α is the nonideality parameter of electrons (e) or ions (i), q_α is the charge of the respective particle ($q_{\text{e,i}} = \pm \text{e}$), a is the Wigner-Seitz radius, k_{B} is the Boltzmann constant, T_α is the temperature of particles, ε_0 is the electric constant. Interaction between particles is of a Coulomb nature. Laser cooling has the same effect on ions as application of a friction force resulting in a reduced kinetic energy of the ions.

The behavior of such systems is conveniently described by means of the molecular-dynamics method. It involves solving the motion equation for particles with very different velocities due to the large difference in mass and temperatures between electrons and ions. This situation calls for the use a small step time discretization to describe the motion of electrons on a long time scale for which shifts of ions are noticeable. This factor combined with the long-range Coulomb interaction for a large ensemble of particles poses a problem when the molecular-dynamics method is applied to

*Also at Siberian Federal University, Krasnoyarsk.

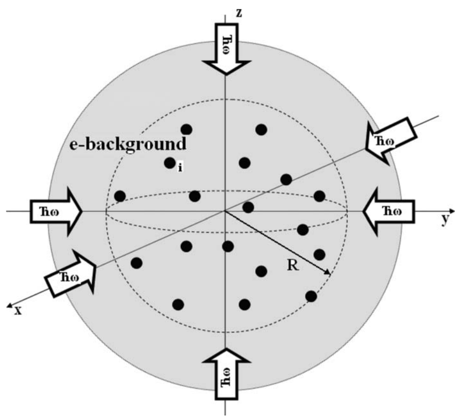

FIG. 1. The 3D laser cooling scheme of an ion cloud in a uniform spherical symmetric negative background.

describe the behavior of an electron-ion plasma.

For a weakly nonideal electron subsystem, the problem can be eased (as suggested in [7]) by considering the motion of ions on a negative background that neutralizes their charge, where the role of the background is played by electrons. Ignoring the electron-ion energy exchange and taking into account only the Coulomb interaction of ions with each other and with the background we obtain a model of a one-component plasma [1], where ions are point charges. The electron-ion energy exchange due to elastic collisions is disregarded in this model. To allow for this exchange, a viscosity of the electron background is introduced, and the motion of ions in this background is treated as the motion of Brownian particles.

Let a uniform spherically symmetric negative (electron) background with the time-constant charge density $\rho = ne$ contains N ions with the same macroscopic density distribution in a sphere of R radius $N = 4\pi R^3 n/3$ (Fig. 1). In this situation, the Coulomb interaction of ions with the background keeps them localized in the center. Ions are cooled by six laser fields of a three-dimensional (3D) symmetric configuration (Fig. 1) having equal amplitudes and the frequency ω shifted by $\Delta = \omega - \omega_{21} < 0$ to the red range from the frequency ω_{21} of the quantum transition of ions.

The temperature of the electron subsystem is assumed to be constant and fairly high (~ 100 K) so that the subsystem would be weakly nonideal in the concentration range of interest ($10^5 \div 10^9$ cm^{-3}) and the three-body recombination would be insignificant. The cooling of electrons due to collisions with ions can be compensated by heating them up with a low-intensity microwave radiation [6,7]. Under these conditions ions get involved in the following interactions.

A. Coulomb interaction

The Coulomb interaction of ions includes the interaction with a uniform electron background and interaction with each other. The force acting on the kth ion in a uniform spherical by symmetric background is described by a simple expression

$$(\mathbf{F}_{\text{bg}})_{\text{k}} = -\frac{ne^2}{3\varepsilon_0}\mathbf{r}_{\text{k}}, \tag{2}$$

where $\mathbf{r}_{\text{k}}$ is the radius vector of the kth ion. Note that according to Eq. (2), the force acting from the background corresponds to the harmonic potential similar to the trapping potential in ion traps [17]. In addition, every kth ion is subject to a Coulomb force acting from other ions

$$\mathbf{F}_{\text{k}} = \frac{e^2}{4\pi\varepsilon_0}\sum_{\text{n}\neq\text{k}}^{\text{N}}\frac{(\mathbf{r}_{\text{k}} - \mathbf{r}_{\text{n}})}{|\mathbf{r}_{\text{k}} - \mathbf{r}_{\text{n}}|^3}. \tag{3}$$

B. Thermal interaction of ions with the background

The change of the average energy of ions ε_{i} resulting from the energy exchange with weakly nonideal electron system due to elastic Coulomb collisions is described by the equation [18,19].

$$\frac{2}{3k_{\text{B}}}\frac{\delta\varepsilon_{\text{i}}}{\delta t}\bigg|_{\text{ie}} = -\frac{2m}{M}\nu_{\text{ei}}T_{\text{i}} + \frac{2m}{M}\nu_{\text{ei}}T_{\text{e}} \tag{4}$$

where M and m are the masses of ion and electron, respectively; and ν_{ei} is the rate of elastic electron-ion collisions.

When performing numerical simulation of ions dynamics, every ion is treated as a Brownian particle (bearing in mind that $m/M \ll 1$), moving with $\boldsymbol{v}$ velocity in an electronic gas. This approach allows us to introduce a friction force affecting the ion $\mathbf{F}_{\text{fr}} = -\eta\boldsymbol{v}$, where η is the friction coefficient η

$$\eta = m\nu_{\text{ei}}. \tag{5}$$

This force accounts for the first term in the right-hand part of Eq. (4). In order to allow for the heating of ions by electrons [the second term in Eq. (4)], we introduce a random force $\tilde{\mathbf{F}}_{\text{r}}$ [20], defined as a δ-correlated Gaussian process. For the purpose of a numerical experiment, a random force $\mathbf{F}_{\text{r}}$ can be used, which is an average over the integration step Δt, with projections onto the coordinate axes

$$(F_{\text{r}})_{\text{j}} = \frac{1}{\Delta t}\int_{t}^{t+\Delta t}(\tilde{F}_{\text{r}})_{\text{j}}dt, \quad j = x,y,z$$

exhibiting a Gaussian distribution of the probability density

$$P[(F_{\text{r}})_{\text{j}}] = [2\pi\langle(F_{\text{r}})_{\text{j}}^2\rangle]^{-1/2}\exp\left[-\frac{(F_{\text{r}})_{\text{j}}^2}{2\langle(F_{\text{r}})_{\text{j}}^2\rangle}\right]. \tag{6}$$

Their dispersion is derived from the equation of heat balance between the ion and electron subsystems at $T_{\text{i}} = T_{\text{e}}$ and is equal to

$$\langle(F_{\text{r}})_{\text{j}}^2\rangle = 2\frac{\eta k_{\text{B}}T_{\text{e}}}{\Delta t}. \tag{7}$$

The time step Δt here has to satisfy the following requirement: $\Delta t \ll (\eta/M)^{-1}$ [20].

C. Interaction with the light field: spontaneous light pressure force

The optical field is formed by three pairs of mutually orthogonal counterpropagating small-amplitude light beams (Fig. 1), satisfying the condition

$$|V_0| \ll \gamma, \tag{8}$$

where $|V_0|/2$ is the Rabi frequency for an individual beam, γ is the spontaneous decay rate of the excited ion state. In this

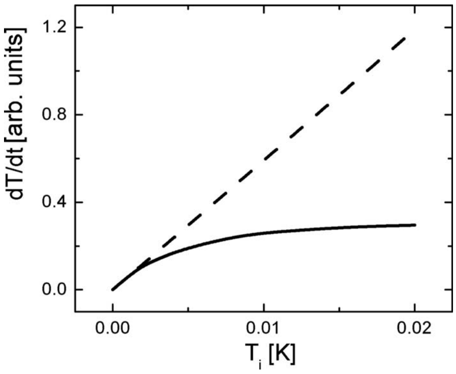

FIG. 2. The laser cooling rate of Be^+ at the detuning $\Delta=0.5\gamma$ versus the ion temperature as found from Eq. (9) (the solid line) and Eq. (11) (the dashed line).

case, every ion experiences friction force along the direction of the unit basis vector $\mathbf{e}_j$ ($j=x,y,z$) of the Cartesian coordinate system; this friction force is the sum of projections of spontaneous light pressure forces induced by each light beam [21–23]

$$(F_l)_j = -\chi(v_j)v_j, \quad \chi(v_j) = \frac{\hbar k_j^2 \gamma |\Delta||V_0|^2}{[(\Delta - k_j v_j)^2 + g^2][(\Delta + k_j v_j)^2 + g^2]}, \tag{9}$$

where $k_j=k=\omega/c$, $g^2 \approx \gamma^2/4$, $v_j=(\boldsymbol{v}\cdot\mathbf{e}_j)$ is the projection of the ion velocity $\boldsymbol{v}$ on the respective coordinate axis, c is the light speed. Under condition (8), this model provides a good approximation for ions with the quantum transition $F=0 \rightarrow F=1$ or when the method of alternating light beams [24] [1] is used.

As already mentioned above, cooling is normally analyzed in the approximation of slow ions, i.e., when

$$k\langle v\rangle \ll \gamma, \tag{10}$$

$\langle v\rangle$ is the mean thermal velocity. Then the force in $(F_l)_j$ can be represented as

$$(F_l)_j = -\chi_0 v_j, \quad \chi_0 = \chi(0) = \frac{\hbar k^2 \gamma |\Delta||V_0|^2}{[\Delta^2 + g^2]^2}, \tag{11}$$

where χ_0 is the light-induced viscous friction coefficient. However in practice condition (10) may appear unachievable, particularly at the initial stages of cooling. This condition of slowness Eq. (10), for example, for Be ions studied herein and in [13,14] is violated already at temperatures $T_i \geq T_i^* = 0.012$ K. Our estimates based on the results of [8–10] suggest that ion temperatures reach the values $T_i \gg T_i^*$ in time $\tau \ll (\chi_0/M)^{-1}$ due to rapid correlation heating. Figure 2 illustrates the difference in the cooling rates for forces determined by Eq. (9) and (11) and proves that Eq. (9) is the correct one for an adequate simulation of the process of laser cooling of UP. Therefore all our further calculations were made with the use of Eq. (9).

[1] In this case, each of the six laser beams is applied one at a time with the repetition rate Ω ($\chi_0/M \ll \Omega \ll \gamma$) so that ions interact alternately with one of the light beams for approximately 1/6 of the $2\pi\Omega^{-1}$ period. So ions are exposed to just one light beam at any time. Then the term $|V_0|$ in Eq. (9) has the meaning of the Rabi frequency averaged over the $2\pi\Omega^{-1}$ period.

D. Fluctuation heating

Quantum fluctuation of the radiation force [22] is another important heating mechanism of ions. The heating rate (the rate of change of the mean kinetic energy) for a single low-intensity standing wave and a two-level ion is determined by diffusion of ions in the velocity space and equals [22]

$$\Lambda_1 = \frac{(\hbar k)^2 \gamma |V_0|^2}{2M(\Delta^2 + \gamma^2/4)}. \tag{12}$$

Considering that the fields in our model of six counterpropagating waves have symmetrical configuration, we will take the total heating rate to be equal to $\Lambda = 3\Lambda_1\xi$, where ξ is the coefficient ~ 1 dependent on the mode of exposure to light fields and on the field polarization. Equation (12) disregards the particle velocity factor since the effect of fluctuation heating on the cooling process becomes noticeable only at low temperatures ($T \sim \hbar\gamma/k_B$), when condition (10) is certainly satisfied. In order to account for this heating, introduce a random force $\mathbf{F}_{fl}$ [23], exactly as force $\mathbf{F}_r$ has been introduced to simulate the heating of ions by electrons (see Sec. II B). The force $\mathbf{F}_{fl}$ dispersion is defined as

$$\langle \mathbf{F}_{fl}^2 \rangle = \frac{2M\Lambda}{\Delta t}. \tag{13}$$

With all the above interactions taken into account, a system of N motion equations acquires the form

$$M\frac{d\boldsymbol{v}_k}{dt} = [(\mathbf{F}_{fr})_k + (\mathbf{F}_r)_k] + [(\mathbf{F}_{bg})_k + \mathbf{F}_k] + [(\mathbf{F}_l)_k + (\mathbf{F}_{fl})_k], \quad k = 1 \ldots N. \tag{14}$$

III. RESULTS

System of stochastic differential Eq. (14) was solved by using the fourth-order Runge-Kutta method with a Δt step under the condition

$$\Delta t \ll (\eta/M)^{-1}, \chi_0^{-1}, \omega_i^{-1},$$

where ω_i is the ion plasma frequency. The numerical results are given for $N=1000$, $T_e=100$ K for various plasma concentrations n. The temperature of ions hereinafter will be

TABLE I. The range of detuning, minimum temperatures, and nonideality parameters of ions for various plasma concentrations.

n, cm^{-3}	10^5	10^6	10^7	2×10^7
$\lvert\Delta\rvert/\gamma$	$0.55\div0.6$	$0.55\div0.6$	$1.0\div1.2$	$1.4\div1.6$
$T_{i,min}$, K	0.00065	0.00127	0.013	0.07
Γ_i	193	213	45	10.5

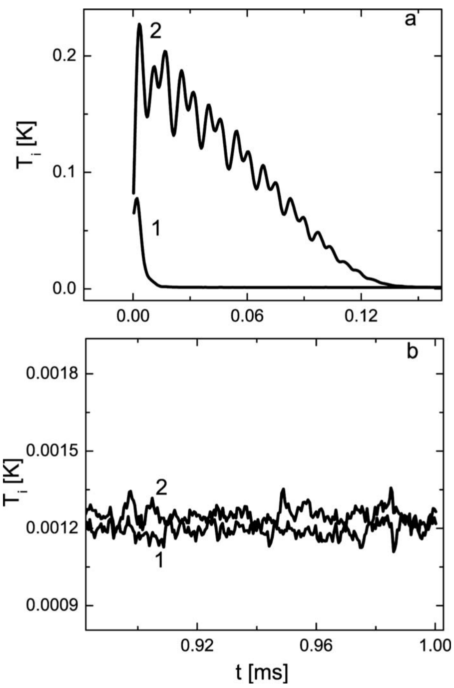

FIG. 3. The effective ion temperature curve at the transient (a) and stabilized (b) stages of the cooling process obtained with the use of Eq. (11) (curve 1) and Eq. (9) (curve 2).

understood as the effective temperature derived from their average kinetic energy $\varepsilon_k: T_i = 2\varepsilon_k / 3k_B$.

The initial space distribution of ions was assumed randomly homogeneous with the Maxwell velocity distribution corresponding to T_{i0} temperature at which $\Gamma_i \sim 1$. For instance, $T_{i0}=0.07$ K was taken for concentrations $n=10^5$ and 10^6 cm^{-3} and $T_{i0}=0.5$ K for concentrations $n=10^7$ cm^{-3}, 2×10^7 cm^{-3}. Calculations continued until an ordered spatial structure of ions was reached. A range of laser radiation detunings was preliminary determined for every concentration within which the lowest temperatures of ions can be achieved (see Table I) during cooling, provided the field amplitude is $|V_0|/\gamma = 0.3$. As can be seen from Table I, the ion subsystem becomes strongly nonideal as a result of the cooling in all the above cases.

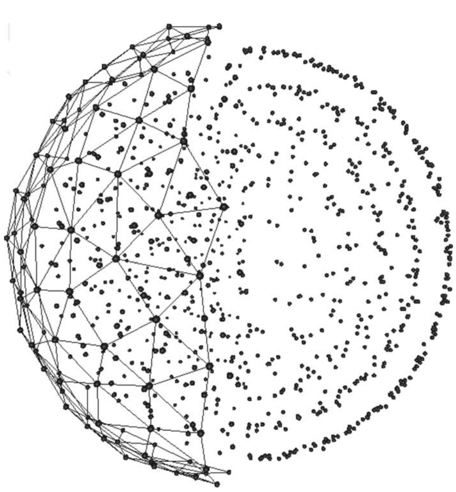

FIG. 4. The ordered space distribution of ions under laser cooling. The left-hand side of the drawing shows particle distribution in the external layer.

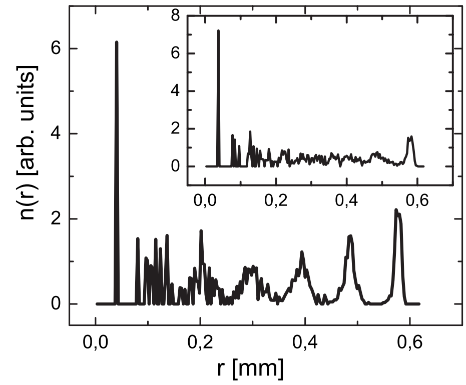

FIG. 5. Final distribution of ion density $n(|\mathbf{r}|)$. Inset: $n(|\mathbf{r}|)$ at the time t_1 when the minimum temperature is achieved.

Note that the nonideality requirement $\Gamma_i \geq 1$ could not be satisfied at concentrations $n > 2\times 10^7$ cm^{-3} and under the conditions discussed ($T_e \leq 100$ K and $|V_0|/\gamma \leq 0.3$), whereas with linear approximation (11) applied to the laser friction force, the values $\Gamma_i \geq 1$ were obtained even for $n \geq 10^9$ cm^3.

Figures 3–7 illustrate the numerical results obtained for the concentrations $n=10^6$ cm^{-3} and optical detuning $\Delta = -0.6\gamma$. As one can see in Fig. 3(a), a rapid growth of the temperature (mean kinetic energy) of ions is observed at the very initial stage (the initial temperature is 0.07 K), followed by a cooling stage. This growth is insignificant and the cooling is monotonic and occurs during $\tau_{cool} \sim 10^{-5}$ s when Eq. (11) is used for the laser friction force, whereas application of Eq. (9) for approximation of the laser friction force yields a much higher growth of energy, and the cooling process exhibits a damped oscillations behavior (curve 2) and lasts longer ($\tau_{cool} \sim 10^{-4}$ s). The minimum temperature [Fig. 3(b)] (≈ 0.0012 K) however is actually the same in both cases. The difference observed in the cooling processes is due to the difference in the cooling rate, which is equal to $2\chi_0/M \approx 7.1\times 10^5$ s^{-1} in the first case. In the second case, the cooling rate is initially 6.6×10^4 s^{-1} and slows down to $\approx 1.5 \times 10^4$ s^{-1} with heating; then as the temperature goes down it

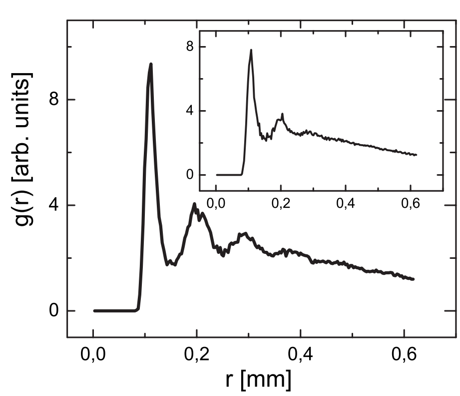

FIG. 6. The pair-correlation function at the final stage of cooling. Inset: the same function at the time t_1 when the minimum temperature has just been reached.

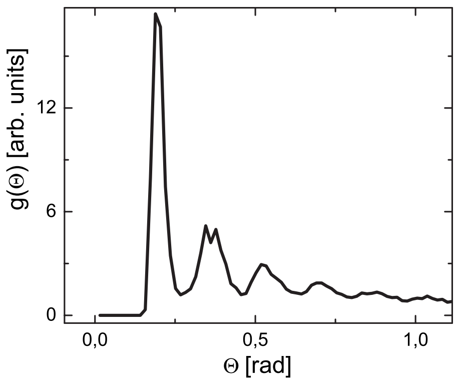

FIG. 7. The angular pair-correlation function $g(\theta)$ for the external layer of ions.

starts growing again reaching the value 7×10^5 s^{-1} at the lowest temperature. Such a nonmonotonic temperature behavior can be attributed to the initial distribution of ions which was assumed to be homogeneous rather than equilibrium (Boltzmann distribution) and to the effect of correlation heating [8]. The potential energy of ionic gas is partially transformed into the kinetic energy of radial motion of ions during relaxation to result in radial oscillations (similar to the "breathing" mode oscillations [25]) of the ion density at the frequency of the order of the ionic plasma frequency ($\omega_i \approx 4.5\times10^5$ s^{-1}). We believe that this might have been one of the reasons for the oscillation of the mean kinetic energy observed in the experiments reported in [26].

The effect of laser friction force on these oscillations is that of damping. As the rate of cooling increases, the effect becomes less noticeable [see Fig. 3(a)]. When Eq. (11) is used, the damping constant $\chi_0/M\sim3.5\times10^5$ s^{-1} is comparable with the oscillation frequency ω_i, therefore the temperature jump is insignificant and the oscillations fade away rapidly. When Eq. (9) is used, the damping constant at the initial stage is by an order of magnitude lower than in the former case. Therefore the temperature curve has the form of damped oscillations and the temperature change occurs in a much longer period of time.

As a result of the cooling and transition into a strongly nonideal state, ions tend to form a number of coaxial spheres (Fig. 4), the so-called "Coulomb ball" [27]. The distance between their surfaces is of the order of the Wigner-Seitz radius (a), with a hexagonal distribution of particles in a layer (the external layer ions are connected by lines for better illustration). This structure is formed in $\tau_{cr}\sim10^{-3}$ s, which is by an order of magnitude longer than the time of cooling $\tau_{cool}\sim10^{-4}$ s.

Note that the time required for formation of the structure is determined solely by the Coulomb interaction (in the absence of random forces and friction forces) equals $\tau_{cr}\sim\omega_i^{-1}$. In the presence of strong friction and random forces the motion of ions in the process of formation of the space structure may have a diffusive character if the effective length of the ion-free path in a "viscous photon medium" satisfies the condition $\lambda_{eff}=M\langle v\rangle/\chi_0\ll a$. The spatial diffusion coefficient (in a cold state) is $D\sim k_B(T_i)_{min}/\chi_0$. Then the time τ_{dif} of formation of layers due to diffusion depends on the time required for ions to pass the distance $s\sim a/2\approx0.31n^{1/3}$: $\tau_{dif}=s^2/D \approx 0.1\chi_0/k_B(T_i)_{min}n^{2/3}$. In a general case, the crystallization time is determined by the longest of these typical times: $\tau_{cr}=\max(\tau_{dif},\ \omega_i^{-1})$. In the case under consideration, $\tau_{cr}=\tau_{dif}\approx0.75\times10^{-3}$ s since $\tau_{dif}\gg\omega_i^{-1}=2.2\times10^{-6}$ s. This result is in good agreement with the crystallization time obtained by the simulation.

The delay of crystallization with respect to cooling is also apparent from the comparison of the radial ion density distribution at different times. The inset in Fig. 5 shows radial distribution at time $t_1=1.5\times10^{-4}$ s, i.e., immediately as soon as the minimum temperature has been reached. It can be seen that only one external layer is formed in the cooling time; and formation of the final structure (consisting of at least four layers) requires $\tau_{cr}\approx10^{-3}$ s.

A similar behavior is observed for the pair-correlation function $g(r)$ (which defines the probability of detecting two ions at a distance r [1]) (Fig. 6): the correlation between ions grows after cooling while the ion temperature remains unchanged. Here r is the distance between particles. The pair-correlation function $g(\theta)$ (is the angular distance between particles) is shown in Fig. 7 for the external layer of ions. By comparing Figs. 6 and 7 one can see that the correlation between ions in the external layer is noticeably higher than the correlation in the entire ensemble.

IV. CONCLUSION

To summarize the above, the Brownian dynamics technique used to simulate laser cooling of a quasistationary electron-ion plasma can provide an adequate description of the cooling dynamics and the structure of an ion subsystem, provided energy exchange with the electron subsystem is taken into account. It has been shown that cooling (in spherically symmetric UP case) results in the formation of a quasicrystalline structure of plasma ions, widely known as the "Coulomb ball" [27]. It has been found that the crystallization in a low-density plasma may take a much longer than the cooling time. This finding can be a very important consideration in experiments on cooling and crystallization in a nonstationary plasma.

An adequate description of the laser cooling behavior of electron-ion plasma requires that the nonlinear dependence of light-induced friction (8) on the particle velocity be taken into account. The use of an approximate expression for this force, which is only valid for "slow" ions [Eq. (9)], yields overestimated cooling rates and concentrations required to obtain large nonideality parameter $\Gamma_i\gg1$, which is the most obvious for light ions. Such a cooling, in particular, for Be^+ is problematic already at concentrations $n\sim3\times10^7$ cm^{-3} and higher.

The approach employed in the paper for modeling of beryllium plasma and the obtained results feature a broad generality: the Brownian character of heat exchange, the influence of nonlinearity of laser friction force, the effect of the delay of crystallization—all these effects can be observed in real experiments on laser cooling of electron-ion plasma of any other elements. For example, as our computations with the presented model of laser cooling showed for electron-ion plasma with heavy ions (Sr^+ or Ba^+), nonlinearity of laser friction force and electron heating of ions make it impossible obtaining the value of the nonideality parameter $\Gamma_i \geq 1$ for concentrations $n \geq 2 \times 10^8$ cm^{-3}.

In closing the discussion we note that the system of mutually orthogonal light beams can be used not only for cooling of ions as suggested herein but also for efficient viscous confinement of ultracold plasma. In other words, the area of intersection of the light beams can play the role of optical molasses (OM) [28] for UP with resonant ions. The confinement time τ_c is determined by the effective ambipolar diffusion coefficient (see [29]) $D_A = T_e/\chi_0$: $\tau_c \sim R_0^2/D_A$, where R_0 is the typical OM size. For a spherically symmetrical cloud confined in OM, the described effect of ionic component crystallization can be observed in the central part of the cloud ($|\mathbf{r}| \ll R_0$), provided τ_{cr} is much less than the viscous confinement time τ_c. This becomes practical when the size of the plasma cloud $\sim R_0$ is fairly large while the degree of nonisothermality of UP $\varsigma = T_e/T_i$ is not too high

$$\left(\frac{a^2}{R_0^2}\varsigma\right) \sim \frac{\varsigma}{N^{2/3}} \ll 1$$

where N is the total number of ions in a plasma cloud.

[1] S. Ichimaru, Rev. Mod. Phys. **54**, 1017 (1982).
[2] V. E. Fortov, A. G. Khrapak, and I. T. Yakubov, *Physics of Strongly Coupled Plasma* (Oxford University Press, New York, 2005).
[3] T. C. Killian, S. Kulin, S. D. Bergeson, L. A. Orozco, C. Orzel, and S. L. Rolston, Phys. Rev. Lett. **83**, 4776 (1999).
[4] S. Kulin, T. C. Kilian, S. D. Bergeson, and S. L. Rolston, Phys. Rev. Lett. **85**, 318 (2000).
[5] T. C. Killian, M. J. Lim, S. Kulin, R. Dumke, S. D. Bergeson, and S. L. Rolston, Phys. Rev. Lett. **86**, 3759 (2001).
[6] A. P. Gavrilyuk, I. V. Krasnov, and N. Ya. Shaparev, Tech. Phys. Lett. **23**, 61 (1997).
[7] A. P. Gavrilyuk, I. V. Krasnov, and N. Ya. Shaparev, Laser Phys. **8**, 653 (1998).
[8] M. S. Murillo, Phys. Rev. Lett. **87**, 115003 (2001).
[9] S. Mazevet, L. A. Collins, and J. D. Cress, Phys. Rev. Lett. **88**, 055001 (2002).
[10] F. Robicheaux and J. D. Hanson, Phys. Rev. Lett. **88**, 055002 (2002).
[11] S. G. Kuzmin and T. M. O'Neil, Phys. Plasmas **9**, 3743 (2002).
[12] T. C. Killian, V. S. Ashoka, P. Gupta, S. Laha, S. B. Nagel, C. E. Simien, S. Kulin, S. L. Rolston, and S. D. Bergeson, J. Phys. A **36**, 6077 (2003).
[13] T. Pohl, T. Pattard, and J. M. Rost, Phys. Rev. Lett. **92**, 155003 (2004).
[14] T. Pohl, T. Pattard, and J. M. Rost, J. Phys. B **38**, S343 (2005).
[15] T. C. Killian, T. Pattard, T. Pohl, and J. M. Rost, Phys. Rep. **449**, 77 (2007).
[16] I. V. Krasnov and A. P. Gavrilyuk, J. Phys. B **41**, 125301 (2008).
[17] D. H. E. Dubin and T. M. O'Neil, Rev. Mod. Phys. **71**, 87 (1999).
[18] J. L. Spitzer, *Physics of Fully Ionized Gases* (J. Wiley and Sons, New York, London, 1962).
[19] G. Hazak, Z. Zinamon, Y. Rozenfeld, and M. W. Dharmawardana, Phys. Rev. E **64**, 066411 (2001).
[20] D. V. Heerman, *Computer Simulation Methods in Theoretical Physics* (Springer-Verlag, Berlin, Heidelberg, New York, London, Paris, 1986).
[21] I. V. Krasnov and N. Ya. Shaparev, Pis'ma Zh. Tekh. Fiz. **2**, 301 (1976).
[22] A. P. Kazantsev, G. I. Sourdoutovich, and V. P. Yakovlev, *Mechanical Effect of Light on Atoms* (World Scientific, Singapore, 1990).
[23] V. G. Minogin and V. S. Letokhov, *Laser Light Pressure on Atoms* (Cordon and Breach, New York, 1987).
[24] P. D. Lett, W. D. Phillips, S. L. Rolston, C. E. Tanner, R. N. Watts, and C. I. Westbrook, J. Opt. Soc. Am. B **6**, 2084 (1989).
[25] D. H. E. Dubin and J. P. Schiffer, Phys. Rev. E **53**, 5249 (1996).
[26] Y. C. Chen, C. E. Simien, S. Laha, P. Gupta, Y. N. Martinez, P. G. Mickelson, S. B. Nagel, and T. C. Killian, Phys. Rev. Lett. **93**, 265003 (2004).
[27] O. Arp, D. Block, A. Piel, and A. Melzer, Phys. Rev. Lett. **93**, 165004 (2004).
[28] S. Chu, L. Hollberg, J. E. Bjorkholm, A. Cable, and A. Ashkin, Phys. Rev. Lett. **55**, 48 (1985).
[29] I. V. Krasnov, Phys. Lett. A **372**, 3118 (2008).