Dynamics of cluster particles in a dense plasma

V V Yaroshenko, B M Annaratone, T Antonova, H M Thomas and G E Morfill
Max-Planck-Institut für Extraterrestrische Physik, D-85740, Garching, Germany
E-mail: viy@mpe.mpg.de

New Journal of Physics 8 (2006) 203
Received 21 July 2006
Published 20 September 2006
Online at http://www.njp.org/
doi:10.1088/1367-2630/8/9/203

Abstract. Detailed analysis of the particle dynamics in a small cluster (4 grains), suspended inside a dense plasma, is presented. Estimations of the forces providing the particle equilibrium and main plasma parameters are discussed. The simple ‘dipole’ model of binary dust–dust interactions has been employed to explain the observed distortion of the cluster by a particle rotating beneath. The measured force acting on two different cluster particles revealing pronounced vertical motion is compared with a theoretical model. The good qualitative agreement between measured and theoretical curves strongly suggests that attractive forces due to particle–wake interactions can be communicated not only downstream in the flowing plasma, but in the case of very dense plasmas, also upstream. This is a phenomenon not previously observed.

Contents

1. Introduction 2
2. Determination of the plasma parameters 3
3. Model of the binary dust–dust interactions in the presence of an ion flow 5
4. Interacting forces in small clusters 7
5. The shadow force 10
6. Results 11
References 12

1 Author to whom any correspondence should be addressed. Permanent address: Institute of Radio Astronomy of National Academy of Science of Ukraine, Chervonopraporna 4, Kharkov 61002, Ukraine.
1. Introduction

Small three-dimensional (3D) microparticle clusters, combining properties of nonideal plasmas and condensed matter, provide an ideal model for studying dynamical properties of the complex plasmas. The grains in the plasma acquire high negative charges of the order of $10^3$–$10^4$ elementary charges because of the high electron mobility. They are usually trapped in the space charge sheath of the discharge due to the balance of gravitational and electrostatic forces. The dynamical properties of finite (mostly 2D) clusters have been widely studied theoretically and experimentally, see e.g. [1, 2]. Small clusters have the advantage that the individual particles can be observed and investigated in great detail, shedding light on the fundamental problem of the nature of the interaction forces between the microparticles. In particular, the investigation of the possible attractive force between like-charged particles in a plasma, which has been predicted theoretically, has now become one of the high priority problems [3]–[7].

Up to now, cluster structures have been observed only under sheath conditions. In this paper, we study the particle dynamics in a small cluster suspended in a different plasma regime—the specific secondary plasma produced by an adaptive electrode inside the sheath of the discharge [8]. The main features of our observations are: (i) the particles levitate inside a modified sheath region; (ii) weak external confinement; (iii) the clusters particle number can be externally controlled; (iv) perturbations of the cluster introduced by a rotating grain trapped below the main cluster; (v) use of simultaneous 3D diagnostics for reconstruction of the particle dynamics. As we will show below, these features allow us to estimate interaction forces among the cluster particles, reveal the role of the ion flows (ion–wake interactions) and determine the main plasma parameters.

The experiment is performed in the so-called PKE (plasma-crystal experiment)-Nefedov chamber, which consists of two parallel plate electrodes and glass walls [8]. The upper electrode is radiofrequency (rf) driven at 13.56 MHz, while the lower one is grounded, apart from a small central electrode (‘pixel’), $3.8 \text{ mm} \times 3.8 \text{ mm}$, which can be independently driven in direct current (dc) and rf regimes (for more details about this adaptive electrode, we refer to [9]). In our experiment, the ‘pixel’ is dc grounded through the inductance. Typical experimental conditions are: argon pressure in the range from 35 to 70 Pa, rf voltage on the upper electrode in the range of 200–300 $V_{pp}$ and about 150 $V_{pp}$ in ‘push–pull’ on the lower central pixel. In this case, there is an electric current flowing out from the plasma, and anode-like structures are produced [9].

When spherical melamine–formaldehyde particles (mass density $1.51 \text{ g cm}^{-1}$ and $3.4 \mu\text{m}$ diameter) are injected into the plasma, they are captured in the plasma sheath. By suitably adjusting dc and rf on the central pixel, a glow region inside the sheath is formed, which is much brighter than the surrounding discharge plasma. By simultaneously varying the adaptive electrode voltage, the particles can be transported one by one to a new lower levitation height inside the secondary plasma, where they self-organize into a stable cluster structure [8]. We have investigated 3D clusters with four up to 73 particles. For the present purpose, we focus on the smallest structure of four grains, observed at gas pressure $p \sim 57 \text{ Pa}$. The interesting feature of this observation is the existence of a fifth particle ‘rotating’ in a circular orbit below the small cluster. (Such rotating particles have often been observed in dc and rf experiments [10, 11].)

The particles in the cluster were illuminated by three laser sheets and the particle motion was recorded by three CCD cameras, providing instantaneous particle coordinates $(xyz)$ with a rate of 25 frames per second [8]. The uncertainties in the particle positions are of the order of $3 \mu\text{m}$ in the horizontal plane and $21 \mu\text{m}$ in the vertical direction.
The secondary plasma produced by the adaptive electrode inside the (initial) sheath is interesting by itself, e.g. can it be considered as an almost quasineutral (presheath type) plasma or are the sheath properties still predominant? Moreover, what type of ion flows occur, which can affect the screening length of the plasma, provide the direct dragging influence on highly negatively charged dust particles (the ion drag force), lead to the formation of a region of enhanced ion density downstream of the suspended microparticles, the so-called ‘ion wake’ [12]–[16], and hence affect the cluster particle dynamics?

The plasma density of the bright cloud is measured by spatially resolved spectroscopy which gives \( n_0 \approx 10^{10} \text{ cm}^3 \) [17]. The electron temperature in capacitively-coupled argon plasmas is always about 3–4 eV. We have performed measurements with Langmuir probe and found a higher electron temperature, \( T_e \approx 5 \text{ eV} \) at a gas pressure \( p = 57 \text{ Pa} \). Although the probe was slightly affecting the shape of the plasma bright region, the higher electron temperature, can be explained by the fact that electrons are accelerated out of the plasma in a part of the rf cycle and may thermolize at a temperature somehow higher than in the bulk plasma. As a result, the electron Debye length, \( \lambda_{\text{De}} \approx 160 \mu\text{m} \), is much larger that the ion mean free path \( l_i \approx 30 \mu\text{m} \).

The electric field, \( E \), inside the secondary plasma can be roughly estimated as \( E \approx T_e/eL \), where \( L \) is a characteristic length of the bright plasma cloud (\( L \approx 0.5 \text{ cm} \)). This yields \( E \lesssim 10 \text{ V cm}^{-1} \). At small values of \( E/p \), the ion temperature is generally assumed to coincide with the neutral gas temperature \( (T_i \approx T_n \approx 0.03 \text{ eV}) \). To exclude the ion heating inside the secondary plasma, we calculate \( T_i \) from the mean energy of ions. According to Robertson and Sternovsky [18], one can derive the ion temperature as \( T_i = T_n + (\pi - 2)m_iu^2/6 \) with \( u \) being the ion-drift velocity. The latter is determined by \( u = \mu_iE \), where \( \mu_i \) is the ion mobility, which in general is a function of the electric field through \( \mu_i(E) = \mu_0p^{-1}(1 + \alpha(E/p))^{-1/2} \) with the coefficients \( \mu_0 \) and \( \alpha \) depend on the type of ions [19]. The upper limit of the ion-drift velocity corresponding to the electric field \( E \approx 10 \text{ V cm}^{-1} \) is \( u \approx 1.8 \times 10^4 \text{ cm s}^{-1} \), and hence \( T_i \approx 0.03 + 0.004 \text{ (eV)} \). Therefore the corrections to \( T_i \) do not exceed 10% for the upper limit of the drift velocity and will be even less for real values of \( u \) specifying below, and we are now safe to assume \( T_i \approx T_n \approx 0.03 \text{ (eV)} \).

The resulting ion Debye length is \( \lambda_{\text{Di}} \approx 15 \mu\text{m} \), which is smaller than \( l_i \). In the presence of an ion flow, the effective screening length, \( \lambda_D \), is given by [20]

\[
\lambda_D(u) = [\lambda_{\text{Di}}^{-2}(1 + U^2)^{-1} + \lambda_{\text{De}}^{-2}]^{-1/2},
\]

where \( U = u/V_{\text{Ti}} \) is the thermal Mach number or the ratio of the ion flow velocity \( u \) to the ion thermal velocity \( V_{\text{Ti}} = \sqrt{T_i/m_i} \approx 2.7 \times 10^4 \text{ cm s}^{-1} \). In a plasma with subthermal ion drift, \( U \ll 1 \), the screening length is given by the linearized Debye radius \( \lambda_D \approx [\lambda_{\text{Di}}^{-2} + \lambda_{\text{De}}^{-2}]^{-1/2} \approx \lambda_{\text{Di}} \) because typically \( T_e \gg T_i \), while for suprathermal flow, \( U \gg 1 \), the ions cannot participate in the screening process, and \( \lambda_D \approx \lambda_{\text{De}} \). In our parameter regime, the estimated electric fields cannot accelerate the ions to suprathermal velocities, and the thermal Mach number obeys \( U = u/V_{\text{Ti}} \lesssim 1 \). The effective screening length, \( \lambda_D \) is then close to the ion Debye length and we have a weakly collisional regime, where the ion mean free path is comparable with the plasmas screening length, \( \lambda_D \lesssim l_i \).

New Journal of Physics 8 (2006) 203 (http://www.njp.org/)
To estimate the particle charge, we use an analytical model applicable in the weakly collisional regime for the ions [21, 22]. The ion flux in this case can be approximated as

\[ I_i = \sqrt{8\pi a^2 n_0 V_{Te}^2} \left( 1 + 0.1 \tau \frac{\lambda_D}{l_i} \right). \]  

(2)

Here \( a \) and \( Z_d \) denote the particle radius and charge number, and the dimensionless coefficients \( z \) and \( \tau \) are defined as \( \tau = T_e/T_i \), and \( z = Z_d e^2 / (a T_e) \). Equating the ion flux (2) with the electron flux given by

\[ I_e = \sqrt{8\pi a^2 n_0 V_{Te}^2} \exp(-z) \]  

(3)

one can obtain the particle charge, \( Z_d \). The results of the calculations of the dimensionless particle charge \( z \) for different \( \lambda_D/l_i \) are shown in figure 1. This yields \( Z_d \sim (5/6) \times 10^3 \) for the range of the collisionality parameter under consideration \( \lambda_D/l_i \sim (1/0.5) \).

In the following we will need the neutral drag coefficient, \( \gamma = \nu_{dn} M \), where \( M \) is the particle mass. For the momentum transfer frequency in dust–neutral collisions, \( \nu_{dn} \), we adopt the standard theory [23] yielding \( \gamma = \nu_{dn} M = (8\sqrt{2\pi/3})\delta a^2 n_m n_n v_{Tn} \) \((m_n, n_n \text{ and } v_{Tn} = \sqrt{T_n/m_n} \) are the mass, density and thermal velocity of neutrals, respectively and the factor \( \delta \approx 1.4 \) accounts for diffuse scattering [24]).

The rotating particle gives a unique possibility to estimate the forces acting on the particles and determine the plasma parameters related to this specific region of the secondary plasma. Analysing the coordinates of this grain, we found that its motion occurs almost horizontally (in \((x, z)\) plane) along a slightly disturbed almost circular orbit with an average radius of \( r \approx 215 \mu m \), while the rotational frequency corresponds to \( \Omega_0 \approx 12 \text{ s}^{-1} \). This immediately yields an average value of the centripetal (radial confining) force \( F_{ext} = M \Omega_0^2 r \approx 9.6 \times 10^{-16} \text{ N} \), which is at least two order of magnitude lower than the forces providing equilibrium in the vertical direction (e.g. the gravitational force \( F_g = M g \approx 3 \times 10^{-13} \text{ N} \)).

For a gas pressure of 57 Pa and for maximal measured grain velocity \((V \sim 0.3 \text{ cm s}^{-1})\), the neutral drag force \( F_n = \nu V \) is of the order of \( \sim 1.8 \times 10^{-14} \text{ N} \). To keep the particle moving, we need a driving force in azimuthal direction of the same order of magnitude as the neutral drag force. There have been several attempts for theoretical interpretation of the reason for such grain rotation [10, 11]. All these theories invoke a nonspherical grains. In the following we will not focus on the driving force, but consider the force balance in radial \((r)\) and vertical \((y)\) directions

\[ Mr\Omega_0^2 = Z_d e E_r - F_{id,r}, \]  

(3)

\[ Mg = Z_d e E_y - F_{id,y}. \]  

(4)

Figure 1. The dimensionless particle charge \( z = Z_d e^2 / (a T_e) \) as a function of the collisionality parameter \( \lambda_D/l_i \).
Here $E_r$ and $E_y$ refer to the radial and vertical components of the electric field, while $F_{id,r}$ and $F_{id,y}$ denote the corresponding components of the ion drag force.

Despite the important role of the ion drag force in complex plasmas, a complete self-consistent model of this force has not yet been constructed. We will use the model that is believed to be most relevant for our parameter regime of subthermal ions ($U = u/V_{Ti} < 1$). Therefore the interactions between the ions and the particles can be characterized by the so-called thermal scattering parameter $\beta_T = Z_d e^2/(T_i \lambda_D)$, which in our case $\beta_T \sim 15–20$ is large enough to provide the strong interactions between the ions and dust particles. In this case, the ion drag force can be estimated as

$$F_{id} \simeq \pi \lambda_D^2 (\ln \beta_T)^2 n_i T_i \mu_i E/V_{Ti}. \quad (5)$$

Inserting this in (3) and (4) gives equations determining the electric field components. Taking $Z_d \sim (4.5/5.5) \times 10^3$ and the plasma parameters mentioned earlier yields $E_r \simeq (1.3–1.6) \times 10^{-2}$ V cm$^{-1}$, and $E_y \simeq (4.8–6)$ V cm$^{-1}$. The corresponding ion-drift velocities then satisfy $u_r \lesssim 40$ (cm s$^{-1}$) $\ll V_{Ti}$, $u_i \lesssim 1.1 \times 10^4$ (cm s$^{-1}$) $< V_{Ti}$ and give the thermal Mach number $U \lesssim 0.4$, which justifies the initial assumption $U < 1$. One finds that the condition $U \lesssim 0.4$ leads to an effective plasma screening length $\lambda_D(u) \simeq 20 \mu$m, which is close to the ion Debye radius. Hence, as has been supposed earlier, the secondary plasmas reveal properties peculiar to the presheath region of the discharge rather than to a typical sheath plasma. Furthermore, with the charge recovered from the analytical model, we find that the ratio of the radial force components $F_{id,r}/Z_d e E_r$ does not exceed 0.2 and for vertical components it becomes $F_{id,y}/Z_d e E_y \lesssim 0.3$ at the gas pressure $p = 57$ Pa. This means that in our case, the electric field force remains the main force which determines the particle equilibrium state. Finally, we point out that in the following, we adopt the estimated complex plasma parameters for the cluster region as a whole and assume that all the cluster particles are identical (i.e. they are of the same mass $M$ and charge $Q = Z_d e$).

3. Model of the binary dust–dust interactions in the presence of an ion flow

Considering the particle dynamics in the cluster, we firstly need to model the binary grain interactions. The usual assumption of the pure symmetric Coulomb screened dust–dust interaction is too simple for plasmas with ion flows. In fact, the ion flow redistributes the ions inside the Debye sphere leading to the formation of a region of enhanced ion density downstream of the suspended microparticles—an ‘ion wake’ [12]–[16]. This effect modifies the dust–dust interaction potential introducing an anisotropy in particle interactions. In particular, according to recent laboratory studies of ‘particle pairing’ in a plasma sheath, the ion wake can cause an attractive force between grains [12, 26, 27].

Following the standard approach, we model the ion wake as a positive point-like effective charge $q$ located at some small distance $l$ downstream, within the Debye sphere [26], [28]–[30]. The combined system ‘particle charge and its ion density enhancement’—in a further simplification [31]—treated as an uncompensated residual particle charge $Q$, plus an electric dipole $P = ql$, which accounts for the anisotropy of the plasma due to the ion flow. Accordingly the force describing binary dust–dust interactions can be represented as a combination of the electrostatic force due to the repulsion of the residual like particle charges and a dipole force
Figure 2. Scheme of the binary particle interactions. The vertical coordinate $y$ is directed to the lower electrode (along the ion flow).

due to the streaming ions

$$F = \frac{Q^2}{\lambda_D^2} \left[ \frac{(1 + \kappa) \exp(-\kappa)}{\kappa^2} + 3\zeta \left(1 - 3\sin^2 \chi \right) \right],$$

(6)

where $F$ denotes the force component along the radius-vector between two particles (see figure 2), $\kappa = \Delta / \lambda_D$, is the so-called lattice parameter, $\Delta$ is the interparticle distance, and the dimensionless coefficient $\zeta = q^2 l^2 / (Q \lambda_D)^2$ specifies the value of the electric dipole moment. The quantity $\chi$ denotes the angle between the normal to the ion flow and the radius-vector connecting the two particles (figure 2).

The use of model (6) implies that the dipole interactions are unshielded. The latter can be justified by the calculations of a test particle potential in the case of anisotropic (in velocity space) plasmas [32, 33]. These results (both mathematically and numerically) support the ‘unshielded dipole’ interaction at distances larger than the screening length. Therefore, the approximation (6) is valid for the interparticle distances $\Delta > \lambda_D$.

The mutual particle interactions are now determined by the force, which can be attractive or repulsive depending on the plasma parameters and relative positions of the two dipoles: when $(1 - 3 \sin^2 \chi) > 0$, so that $|\chi| < 0.62$, grains are repelled since both of the forces act in the same direction, while for $|\chi| > 0.62$, the electrostatic and dipole forces compete with each other, and the resulting force corresponds to either attraction or repulsion. Figure 3 shows the normalized force $f = F \lambda_D^2 / Q_n^2$ as a function of $\kappa = \Delta / \lambda_D$ for $\chi = 2\pi/3$ at different coefficients $\zeta$. The interactions are repulsive only for a rather narrow region of $\kappa$, which is specified by the contribution of the dipole term in the total force (or in other words by a value of $\zeta$). It can be easily verified that the interparticle force (6) changes its sign when

$$\zeta = \frac{1}{3} (1 + \kappa) e^{-\chi} \frac{\kappa^2}{3 \sin^2 \chi - 1}.$$  

(7)

Later we will use this expression to estimate the dipole contribution (the coefficient $\zeta$) in the total force for given particle positions.
Figure 3. Normalized interacting force, \( f = F_{\lambda D}^2/Q^2 \) for \( \chi = 2\pi/3 \) as a function of \( \kappa = \Delta/\lambda_D \) for \( \zeta = 0.3 \) (solid line), \( \zeta = 0.2 \) (dashed curve) (a) for small interparticle distances (\( \kappa \lesssim 4 \)) and (b) for large distances (\( \kappa \gtrsim 5 \)).

Figure 4. A cluster of four particles: (a) the typical particle positions in the cluster; (b) the cluster distortion of the orbiting grain B. Units on the axis are in mm. Reprinted with permission from T Antonova et al 2006 Phys. Rev. Lett. 96 115001. © 2006 by the American Physical Society [8].

4. Interacting forces in small clusters

In the undisturbed phase, the cluster particles were suspended in the weak electric field \( (E_y \simeq (4.8–6) \text{ V cm}^{-1}) \) of the secondary plasma and form almost a tetrahedron with the average interparticle distances of the order of 160–200 \( \mu \text{m} \). Apart from thermal vibrations, the cluster particles are apparently periodically disturbed by the motion of the lower orbiting particle. Unfortunately, typical particle displacements are of the same order as the uncertainties of the measurements. Because of this we consider the dynamical behaviour of the cluster particles at the specific stage, when the interparticle interactions lead to strong displacements and distortion of the cluster structure. Let us denote the cluster particles lying almost horizontally as A, A1, A2, an upper grain as D, and the grain orbiting below the cluster as B (figure 4(a)). As the (perturbing) particle B approaches, particles, A and A1, break the bond with the upper grain D and start to move towards the rotating particle B (almost along the vertical direction) (see figure 4(b)). When the particle distances A–B and A1–A become small enough, the particles return back to their places inside the tetrahedron structure to restore equilibrium. Variations in the vertical coordinates of A and A1 during this distortion phase (frames 151–155 correspond
to motion downstream, while frames 155–158 to the upward motion) are shown in figure 5. The error bar is due to the pixel discretization of the particle orbit segment.

The motion of dust particles in clusters and in solid state lattices can be explained by an external confinement potential and the influence of the neighbour particles’ field. The particle B has an orbital motion almost unchanged in time. Its horizontal displacement is limited by the external confinement (and allow us measuring of this quantity, ∼10⁻¹⁵ N), which is assumed to hold for all the cluster particles. As for the vertical confinement, it is reasonable to assume that this force is also weak, smaller than our experimental error. The latter assumption can be justified by a few points: (i) the almost symmetric down and up motion of the particle A and A1 during distortion phase (while the vertical confining force would be in one direction and it would destroy the symmetry); (ii) the lower electrode is in fact too far from the cluster; (iii) the dimension of the cluster is more than ten times smaller than the dimensions of secondary plasma, which can generate the external confinement. Therefore, one can suppose that positions of any cluster particle is mainly determined by the interactions with other particles, including the grain B.

In order to obtain quantitative results from the experiment, we consider the equation of motion of these two particles (subscript \( n \) refers to A and A1)

\[
\mathbf{F}_n = M\ddot{\mathbf{r}}_n + \gamma \dot{\mathbf{r}}_n, \tag{8}
\]

where \( \mathbf{F}_n \) is the force acting on the \( n \)-th particle. At the high gas pressures used in the experiment the inertia term is always much smaller than the neutral drag force, so that we may estimate the force acting on the individual particles during snapshots. From the 3D trace of the particle, one can determine the vector velocity and, using equation (8) reconstruct the instantaneous vector force acting during the attraction–repulsion phase, when the particle displacements are much larger than the uncertainties of the measurement. Since the motion in the vertical (\( y \)) direction is most pronounced, we consider the vertical component of the total force, \( \mathbf{F}_n \). Figure 6 shows the computed vertical projection of the measured force acting on cluster particles A and A1 due to interactions with all other grains during frames 151–158.

\[\text{Figure 5. Variations in vertical coordinates } y \text{ of the cluster particles A and A1 during the distortion phase (the axis } y \text{ is directed to the lower electrode).}\]
Using the results of the force measurements, we are now able to investigate whether the
dust–dust interactions described by (6) are compatible with observations. As the interparticle
distances vary between 140–220 µm, and are much larger than the effective screening length
\( \lambda_D \sim 20 \mu m \) (so that \( \kappa \sim 7–11 \)), we can safely use the dipole approximation (6), for modelling
the dust–dust interactions. From the measured particle coordinates, one can easily calculate the
instantaneous interparticle distance and angle \( \chi \). In particular, at the closest approach of grains,
A and A1, we compute \( \kappa_{\text{min}} \sim 7 \) and the angle \( \chi \) be in the range \( \chi_{\text{min}} \sim 1–0.75 \). For the grains
A and B, we obtain \( \kappa_{\text{min}} \sim 8 \), and the angle \( \chi_{\text{min}} \) lies between 0.65–0.7. One can consider the
values of \( \kappa_{\text{min}} \) as points where the transition from repulsion (due to like-charges) to the dipole
attraction occurs (the left-hand side of equation (6) goes to zero). Hence, substituting \( \kappa = \kappa_{\text{min}} \) and
\( \chi = \chi_{\text{min}} \) in (7) gives the dipole coefficient in the binary interactions, \( \varsigma \). We get \( \varsigma \sim 0.15–0.35 \n\) (interactions between A and A1) and \( \varsigma \sim 0.35–0.65 \) (interactions between A and B). To calculate
the corresponding force, we take the average values \( \varsigma \simeq 0.25 \) for the binary interactions between
the cluster particles and \( \varsigma \simeq 0.5 \) for the pair A and B. The difference in the dipole coefficients
\( \varsigma \) can be attributed to specific properties of the rotating particle B; (i) its nonspherical form and
thus another value of the particle and wake charges; (ii) its location in the vicinity of the lower
electrode (the distance from the electrode is only 387 µm), and particle B-electrode interactions
probably dominate.

Figure 6. Vertical projections of the force acting on the cluster particles A and
A1 during the distortion phase.

New Journal of Physics 8 (2006) 203 (http://www.njp.org/)
Substituting the values of $\varsigma$, $\lambda_D$ and the particle charge $Q$ in equation (6), yields the binary interparticle force, $F_{nj}$, for a given particle’s position (at $\kappa$ and $\chi$). The total force acting on any grain, $n$, can be then obtained by computing the sum $F_n = \sum_{(j)} F_{nj}$, where $F_{nj}$ is determined by the binary force (6) and $j$ denotes the number of neighbour cluster grains including the orbiting particle, B. Limiting ourselves to the three nearest neighbour contributions ($j = 3$), we can compare the calculated force $F_n$, acting on the cluster particles, A and A1, with the direct force measurements. If they agree, we may take this as a verification of our dipole model for binary interactions.

The results of computation of the vertical projections of the total force $F_n$ acting on the cluster particles A and A1 during the attraction–repulsion phase are shown in figure 6. Both curves demonstrate the functional dependence close to the measured force (except for the frames 156 and 157 for A-particle). For $Z_d = 4.5 \times 10^3$, the calculated force is decisively below the experimental data for both particles. The best quantitative agreement between the measured force and the dipole model is found for the particle charge corresponding to the highest theoretical proposed value $Z_d = 5.5 \times 10^3$. An interpretation of the measurements using only screened Coulomb interactions ($\varsigma = 0$) does not yield any quantitative agreement, the qualitative trend is even opposite to the observations!

Considering the crudeness of the dipole model (assumption of the constant coefficient $\varsigma$, constant screening length $\lambda_D$, the estimate of the ion flow velocity $u$, identical particle charge $Z_d$ etc.) the agreement shown in figure 6 is surprisingly good. Note also, that the external confinement force, $F_{\text{ext}}$, can also affect the interaction force (especially at small particle displacements, when $F_n$ may be of the same order as $F_{\text{ext}} \sim 10^{-15}$ N).

Considering separately the contributions of the neighbouring grains in the total force acting on cluster particle A1, we found that its motion is mostly governed by the dipole attraction to the downstream grain A, while the dynamics of particle A is more complicated and depends on the contributions of all close neighbours: A1, A2 and B, (however it is mostly determined by the dipole attraction to grain B in frames 152–154, when their intergrain separation achieves its minimum). This observation independently supports the finding that at large intergrain separations ($\kappa \gg 1$), the particles indeed interact rather like dipoles than separated charges. The implication is that, we observe a very different situation to that considered e.g. in [13] for the sheath plasma, where the particle–wake attraction was strongly asymmetric and was communicated only downstream along the ion flow (the upper grain experienced no attractive force from the lower one). The difference can be explained by the specific properties of the dense plasma, produced by the adaptive electrode: (i) the small screening length provides the condition $\kappa \gg 1$ and the particle interaction resembles point dipoles rather than lengthy distributed charges, (ii) the ion flow is subthermal, hence information can pass between the interacting particles in both directions and modify the dynamics. Determination of the vertical projection of the force acting on the orbiting particle B and comparison with the model prediction would shed more light on this, but unfortunately, particle B rotates too close to the lower electrode, so that its influence probably dominates in this particle dynamics.

5. The shadow force

Finally, we compare our force estimates with another type of attraction, which is known as ‘shadow force’ [6, 7]. This is related to the shadowing of the plasma or neutral atom flux towards
one of the interacting dust grains by another one. In low-temperature plasmas, the flux of neutral atoms on the dust grain surface substantially exceeds the ion flux \((n_i \ll n_n)\) and we consider only the force caused by neutral particle shadowing

\[
F_{\text{shad}} \simeq -\frac{3}{4\Delta} \pi a^4 n_n T_n \left( \kappa_{\text{at}} - \kappa_{\text{det}} \right) \sqrt{\frac{T_s}{T_n}}.
\] (9)

Here \(\kappa_{\text{at}}\) and \(\kappa_{\text{det}}\) denote the attachment and detachment coefficients respectively and \(T_s\) refers to the dust surface temperature, which is often higher than the neutral gas temperature \(T_n\). Under laboratory conditions, both the temperature difference and the difference between the attachment and detachment coefficients are very small. Thus the sign of the shadow force is determined by the difference between two small parameters \(\delta T/T_n = T_s/T_n - 1 \ll 1\) and \(\delta \kappa/\kappa_{\text{at}} = 1 - \kappa_{\text{det}}/\kappa_{\text{at}} \ll 1\). The value \(\delta T/T_n\) is typically of the order \(\sim 10^{-3}\), while \(\delta \kappa/\kappa_{\text{at}}\) can be slightly larger than this value, especially for rough dust surfaces [6]. Taking \((\kappa_{\text{at}} - \kappa_{\text{det}})\sqrt{T_s/T_n}\sim 10^{-2}\), one obtains an upper value of the shadowing force \(F_{\text{shad}} \simeq 10^{-15} \text{ N}\) (this corresponds to the maximal approach \(\sim 140 \mu m\) between grains A and A1). This estimate is at least five times less than the dipole contribution to the vertical projection of the total force for the same grain separations (frame 154 in figure 6(b)) and is an order of magnitude less than the total attraction force \(F \simeq 10^{-14} \text{ N}\) measured in [8]. Unless the surface properties of the particles are very much different from those used here (perhaps as a result of plasma etching), the shadowing mechanism is unlikely to be responsible for the observed particle dynamics.

6. Results

We have presented a detailed analysis of the particle dynamics in a small cluster (four grains), which was formed inside a dense plasma region created by an adaptive electrode, modifying the typical sheath plasma. The highly resolved 3D measurements were compared with a dipole interaction model and a shadow force interaction scenario. This comparison has provided an important conclusion with respect to binary interactions between particles embedded in subthermal plasma flows and has enabled us to determine main plasma parameters. It turns out that the secondary plasma has properties peculiar to the presheath region of a discharge rather than to a typical sheath plasma.

To model the dust–dust interactions in the presence of subthermal ion flows, we have introduced a ‘hybrid’ type of interaction, containing the screened Coulomb repulsion of the particle charges and attraction due to the positive charge of the ion wake. This simple ‘dipole’ model of binary dust–dust interactions was then employed to describe and model the particle dynamics in small clusters. The resulting theoretically derived force was compared with the measured one for two different cluster particles that exhibited pronounced vertical motion. Good qualitative agreement between the measured and theoretically predicted particle trajectories was obtained. For the experimental plasma conditions (small Debye length and thus \(\kappa > 1\) or even \(\kappa \sim 10\)), the transition from attraction to repulsion develops according to a scenario presented in figure 2(b): the particle interactions are attractive for separations \(\kappa \gtrsim \kappa_{\text{min}}\) (there is a local extremum of the attractive dipole force as indicated in figure 2(b)), then the attractive force rapidly decreases \(\propto \kappa^{-4}\) and at small intergrain distances.
when $\kappa \lesssim \kappa_{\text{min}}$, the particles repel each other. We believe that in large strongly coupled plasma crystals (or liquid complex plasmas), where ion flows into the particle cloud have to exist due to enhanced recombination inside the cloud, dipole interactions, as we have established here, will be a natural consequence—always leading to an attractive force component. The properties of such a system will be discussed in future publications.

References

[1] Melzer A 2003 Phys. Rev. E 67 016411
[2] Ivanov Y and Melzer A 2005 Phys. Plasmas 12 072110
[3] Annaratone B M 1997 J. Phys. IV 7 C4–C155
[4] Ignatov A M 1996 Plasma Phys. Rep. 22 585
[5] Khrapak S A, Ivlev A V and Morfill G 2001 Phys. Rev. E 64 046403
[6] Tsytovich V N and Morfill G E 2002 Plasma Phys. Rep. 28 195
[7] Morfill G E and Tsytovich V N 2003 Plasma Phys. Rep. 29 1
[8] Antonova T, Annaratone B M, Goldbeek D D, Yaroshenko V, Thomas H M and Morfill G E 2006 Interaction force between particles in three-dimensional plasma clusters Phys. Rev. Lett. 96 115001
[9] Annaratone B M, Antonova T, Goldbeck D D, Thomas H M and Morfill G E 2004 Plasma Phys. Control. Fusion B 46 495
[10] Annaratone B M and Morfill G E 2003 J. Phys. D: Appl. Phys. 36 1
[11] Paeva G V, Dahiya R P, Kroessen G M W and Stoffels W W 2004 IEEE Trans. Plasma Sci. 32 601
[12] Schweigert V A, Schweigert I V, Melzer A, Homann A and Piel A 1996 Phys. Rev. E 54 4155
[13] Melandset F and Goree J 1995 Phys. Rev. E 52 5312
[14] Melzer A, Schweigert V A and Piel A 1999 Phys. Rev. Lett. 83 3194
[15] Vladimirov S V, Maiorov S A and Cramer N F 2003 Phys. Rev. E 67 016407
[16] Maiorov S A, Vladimirov S V and Cramer N F 2001 Phys. Rev. E 63 017401
[17] Antonova T, Annaratone B M, Sato T, Thomas H M and Morfill G E 2006 Spectroscopic investigation of 3D cluster environment Abstracts of 13 ICPP in preparation
[18] Robertson S and Sternovsky Z 2003 Phys. Rev. E 67 046405
[19] Frost L S 1957 Phys. Rev. 105 354
[20] Khrapak S A, Ivlev A V, Zhdanov S K and Morfill G E 2005 Phys. Plasmas 12 042308
[21] Lampe M, Goswami R, Sternovsky Z, Robertson S, Gavrishchaka V, Ganguli G and Joyce G 2003 Phys. Plasmas 10 1500
[22] Khrapak S A et al 2005 Phys. Rev. E 72 016406
[23] Epstein P 1924 Phys. Rev. 61 710
[24] Liu B, Goree J, Nosenko V and Boufendi L 2003 Phys. Plasmas 10 9
[25] Khrapak S A, Ivlev A V, Morfill G and Zhdanov S K 2003 Phys. Rev. Lett. 90 225002
[26] Steinberg V, Sutterlin R, Ivlev A V and Morfill G 2001 Phys. Rev. Lett. 86 4540
[27] Samarian A A, Vladimirov S V and James B M 2005 Phys. Plasmas 12 022103
[28] Melzer A, Schweigert V A, Schweigert I V, Homann A, Peters S and Piel A 1996 Phys. Rev. E 54 R46
[29] Ivlev A and Morfill G E 2000 Phys. Rev. E 63 016409
[30] Yaroshenko V V, Ivlev A and Morfill G E 2005 Phys. Rev. E 71 046405
[31] Yaroshenko V V, Thomas H M and Morfill G E 2006 New J. Phys. 8 54
[32] Montgomery D, Joyce G and Sugihara R 1968 Plasma Phys. 10 681
[33] Cooper G 1969 Phys. Fluids 12 2707