Coalescence of complex plasma clouds

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Abstract. Observations, taken in an experiment performed under microgravity conditions with the PKE-Nefedov Laboratory on the International Space Station (ISS), show that two colliding complex plasma clouds ‘sense’ each other long before they touch. The plasma clouds get deformed, the surfaces flatten and the approach velocity reduces with decreasing gap width until—at the point of coalescence—it becomes zero. This behaviour suggests the existence of a plasma sheath surrounding such a complex plasma cloud. Using a simple dynamical model the complex plasma properties are estimated.

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1. Introduction

The physical processes at the interface of a complex plasma and a particle-free plasma have been investigated theoretically in a number of papers [1]–[7]. The microparticles in the complex plasma carry some of the negative charge, and as a consequence the plasma potential—which may be taken to be zero in the particle-free regime—becomes more negative in the complex plasma cloud. The difference between the plasma potentials increases with increasing ‘Havnes parameter’ \( P = Z n_d / n_e \), where \( Z \) is the charge on a microparticle (in units of \( e \), the electronic charge) and \( n_\alpha \) are the number densities of species \( \alpha \) (\( d \) = ‘dust’ particles, \( e \) = electrons, \( i \) = ions and \( n \) = neutrals). Another contribution to the lowering of the complex plasma potential is due to the plasma recombination on the microparticles. The removal of charges requires a replenishment from the particle-free plasma into the complex plasma cloud. This in turn requires a potential step, retarding electrons and attracting the (much slower) ions. This additional step in the plasma potential is therefore a consequence of the property that complex plasmas are thermodynamically open systems [8, 9].

Sheath formation at solid walls is well known and well studied. It may come as a surprise, however, that complex plasma surfaces—where the solid (particle) surface fraction only amounts to \( 10^{-3} \text{–} 10^{-4} \), i.e. the system is extremely porous—also can have a plasma sheath phenomenon. There have been a number of attempts to measure the complex plasma sheath properties, both outside the complex plasma cloud and in the surface layer inside. The most promising technique of using microparticles as tracers has so far not been accurate enough [10, 11], so that other measurements are of great interest.

The reason for studying this surface phenomenon is twofold: on the one hand, it may provide an insight into the kinetics of collective processes such as surface tension, surface elasticity and surface waves—on the other hand, it may give us evidence for collective self-confinement processes, which in turn will allow us to study the onset of cooperative behaviour in ‘small’ systems, research that can only be carried out at the kinetic level. In both cases, the discreteness
of the kinetic measurements will eventually be studied and lead to a better understanding of nano-structures and (classical) nano-processes (e.g., nano-fluidics [12, 13]). Without the ability to conduct ‘observable’ measurements, we are limited to theoretical predictions and simulations—both can never replace experiments.

In this paper, we study the coalescence of two complex plasma clouds observed under microgravity conditions, and propose a simple dynamical model which describes the major physical mechanisms governing the process. By comparing results of the model with measured characteristics, we deduce principal complex plasma parameters, such as the particle charge, coupling parameter and the effective screening length in the complex plasma sheath.

2. Observations

Measurements were conducted with the PKE-Nefedov setup which has been described in [14]. Experiment parameters were: gas (argon) pressure 0.48 mBar, rf discharge power 0.24 W, particle diameter $3.4 \, \mu m$ and mass $3.1 \times 10^{-11} \, g$. Particles were recorded by a video camera at 25 fps frame rate. The complex plasma at rest assumed the typical shape of a spheroid with a central void [14]. By decreasing the pressure from 0.72 to 0.48 mBar at high power the complex plasma cloud was separated into two concave sub-clouds as shown in figure 1(a). After the abrupt decrease to the above-mentioned power level, the clouds coalesced in the outer midplane region (marked by the dashed frame in the figure), eventually reverting back to the equilibrium spheroid structure with a central void (figure 1(b)). There was a gradient of the particle density established in the coalescing clouds, with the mean interparticle distance being about two to three times larger at the surfaces ($\Delta \simeq 0.4 \, \text{mm}$) than in the interior of the clouds.

In figure 2, a sequence of images is presented from the region marked in figure 1, showing the two plasma clouds during the approach up to the point where they coalesce (see the movie). It is seen that the plasma clouds ‘sense’ each other long before they coalesce—the most obvious feature is the flattening of the surface, and there is also a gradual slowing down of the boundary.
velocity as the clouds approach. Figure 3 shows the cloud separation versus time, which was obtained by measuring the minimal distance between the clouds at each video frame.

3. Model

In order to reveal the essential physical mechanisms governing the coalescence, we employ a simplified one-dimensional (1D) model with two semi-infinite complex plasma clouds separated by a (time varying) gap of width $2L(t)$. We assume that the force driving the two complex plasma clouds together is the gradient of the internal pressure produced by interacting microparticles, $F_p$. The forces acting against the pressure are friction (drag) produced by the neutrals, $F_{nd}$, ions, $F_{id}$, and electrons, $F_{ed}$, as well as an electrostatic force due to the mutual repulsion between the clouds, $F_s$. Thus, the force equation for a microparticle of mass $m_d$ reads

$$m_d \frac{du}{dt} = F_p - F_s - F_{nd} - F_{id} - F_{ed}. \tag{1}$$

At high rf power, the ionization rate and hence the ion drag force is relatively strong, so that the clouds are kept apart. Once the power is reduced, the ion drag is decreased and the clouds start to
coalesce due to the internal pressure. Similar to, e.g., conventional gas expansion into vacuum, the density gradient is naturally established in both approaching clouds providing the pressure force.

Based on the force equation (1), we calculate the dynamics of the cloud coalescence and derive the properties of the complex plasma sheath by matching the theoretical results to observations.

3.1. Pressure force

The internal pressure of the complex plasma due to the electrostatic interaction between microparticles (in the nearest-neighbour approximation) is \cite{14, 15}

\[ P_d \sim N_{nm}(1 + \kappa) \Gamma n_d T_d, \]

where \( n_d \) is the volume density and \( T_d \) is the kinetic temperature of the microparticles (which is assumed to be room temperature), \( N_{nm} \) is the number of the nearest neighbours (e.g., in the crystalline state, \( N_{nm} = 12 \) for the fcc and hcp lattices and \( N_{nm} = 8 \) for the bcc lattice), \( \Gamma \) is the coupling parameter defined by

\[ \Gamma = \frac{Z^2 e^2}{T_d \Delta e^{-\kappa}}, \]

with \( \Delta \) the mean separation between microparticles, and \( \kappa \) is the lattice parameter,

\[ \kappa = \frac{\Delta}{\lambda_D}, \]
with $\lambda_D$ the screening length in the complex plasma, which is usually put equal to the local Debye length. In most typical complex plasmas investigated with PKE-Nefedov, we find that $\kappa \simeq 2–4$ [14]. The force due to this internal pressure acting on a single microparticle at the surface is

$$F_p(\Delta) \sim N_{nn}(1 + \Delta / \lambda_D)(Ze)^2 / \Delta^2 e^{-\Delta / \lambda_D},$$

where we used $n_d \simeq \Delta^{-3}$ for the volume density at the surface.

### 3.2. Repulsion of clouds

Now we consider the forces opposing the pressure, starting with the mutual repulsion of the clouds. From figure 2, we see that (at least) for the flattened surface regime we can approximate the cloud as a crystal with parallel lattice planes. If the gap width $2L$ exceeds $\Delta$ considerably—so that particular lattice structure does not play significant role in the interaction between the clouds—we can 'smear out' the particle charges over the plane, with the average surface charge density being

$$\Sigma \sim -Ze / \Delta^2.$$

The corresponding potential profile perpendicular to the lattice plane, in the $x$-direction, is

$$\varphi \sim -Ze / \Delta^2 \lambda e^{-x / \lambda},$$

where $\lambda$ is the relevant screening length. This expression for the potential is valid at distances $x$ much longer than the interparticle distance. Therefore, as long as $2L \gg \Delta$ (‘free plasma’ gap), we have the following representation for the force (per particle) acting on the surface plane due to mutual repulsion of the merging clouds:

$$\tilde{F}_s(L, \Delta) \sim (Ze)^2 / \Delta^2 e^{-2L / \lambda_s},$$

where $\lambda_s$ is the screening length in the complex plasma sheath = the effective Debye length in the gap, which should generally be different from $\lambda_D$, (see next subsection), and the tilde denotes the ‘free plasma’ asymptote.

As the clouds approach each other, the lattice structure becomes important and the electrostatic repulsion between the clouds should naturally be balanced by the pressure, namely, $F_s$ should tend to $F_p$ when $2L \to \Delta$. Therefore, the repulsion force can be identically written as $F_s \equiv F_p + (\tilde{F}_s - F_p) f_F(2L - \Delta)$, so that the combination of the pressure and electrostatic forces in equation (1) is

$$F_p - F_s \equiv [F_p(\Delta) - \tilde{F}_s(L, \Delta)] f_F(2L - \Delta),$$

where the transition between the force asymptotes is described by some interpolation function $f_F$, with $\lim_{x \to 0} f_F(x) = 0$ and $\lim_{x \to \infty} f_F(x) = 1$. In order to ensure stability at $2L = \Delta$, the interpolation function should scale as $f_F \propto x$ at small $x$. The spatial scale of the transition between the
asymptotes is determined by the plasma screening, so that generally one can write \( f_F \sim x/\lambda \) for small \( x \).

### 3.3. Screening length

The screening length in the gap, \( \lambda_s \), cannot be independent of the gap width \( 2L \) as well. We should have \( \lambda_s = \tilde{\lambda}_D \) when \( 2L \) reduces to \( \Delta \), since nothing then distinguishes the gap from any other lattice region. A full analysis of this problem, including effects such as charge-exchange collisions and plasma flows, is beyond the scope (and intention) of the paper. Similar to equation (3), we can write

\[
\lambda_s = \lambda_D + (\tilde{\lambda}_D - \lambda_D) f_s(2L - \Delta),
\]

(4)

where \( \tilde{\lambda}_D \) is the ‘free plasma’ screening length, i.e., the asymptotic value of \( \lambda_s \) for \( L \to \infty \). Function \( f_s(x) \) should not necessarily coincide with \( f_F(x) \). However, for sufficiently small \( x \), we can naturally assume \( f_s \sim f_F \sim x/\tilde{\lambda}_D \).

### 3.4. Drag forces

Other opposing forces acting on the cloud are the drag forces,

\[
F_{ad} = -m_d v_{ad} \delta u_{ad},
\]

where \( \delta u_{ad} \) is the relative velocity between species \( \alpha \) and the microparticles in the complex plasma cloud, and \( v_{ad} \) is the momentum exchange rate [16],

\[
v_{ad} = \frac{4 \sqrt{2\pi} m_\alpha}{3 m_d} \alpha^2 n_\alpha v_{T\alpha} \Phi_{ad}.
\]

Here, \( m_\alpha \) is the mass and \( v_{T\alpha} \) is the thermal speed of \( \alpha \)-species, \( \alpha \) is the radius of the microparticle and \( \Phi_{ad} \) are cross-section functions given in [16]. For collisions with neutrals, the relative velocity, \( \delta u_{ad} = u \), is the cloud velocity (neutrals are at rest), and also \( \Phi_{nd} = 1 \). For ions we have \( \delta u_{id} = u + u_i \simeq u_i \), where \( u_i \) is the ion drift velocity, and for electrons \( \delta u_{ed} = u + u_e \simeq u_e \), where \( u_e \) is the electron drift velocity.

Ions and electrons are created both inside and outside the complex plasma cloud. This implies that with potential variations there will be corresponding (electron and ion) drifts and drifts produced. Deep inside the cloud these will balance, but at the cloud surface we need to consider them separately. We employ the nearest-neighbour approximation and, hence, consider only two consecutive lattice planes. In accordance with equation (2), the resulting potential difference between the surface plane and a point located half-way to the second lattice plane inside the cloud is

\[
|\delta \varphi_{cl}| \sim \frac{Ze}{\Delta^2} \lambda_D (1 - e^{-\Delta/2\lambda_D})^2.
\]

(5)
Ions are attracted towards the lattice plane (they are attracted to the particles, to be precise, but not all ions are absorbed) and the average drift velocity calculated in this way is given by

\[ \frac{1}{2} m_i u_{cl}^2 \sim e|\delta \varphi_{cl}|. \]

Here, we assume that the ion mean-free path is not shorter than half the interparticle distance—this allows us to use the ballistic relation between the velocity and potential (for the conditions of the experiment, both the mean-free path and the half-interparticle distance are about 200 \( \mu \)m). In a similar way, the drift velocity of ions entering the complex plasma cloud from the gap (sheath) can be calculated. We obtain

\[ |\delta \varphi_s| \sim \frac{Ze}{\Delta^2} \lambda_s (1 - e^{-L/\lambda_s})^2 \] (6)

and the corresponding ion velocity is determined by

\[ \frac{1}{2} m_i u_{is}^2 \sim e|\delta \varphi_s|. \]

Thus, the resulting ion drag is

\[ F_{id} = m_i u_{id} (u_{is} - u_{cl}). \]

Clearly, as long as the complex plasma clouds are separated by more than \( \Delta \), the ion velocity into the cloud exceeds that leaving, and therefore the ion drag tends to drive the clouds apart. Since the electron and ion fluxes in the cloud should be (on average) the same, the electron drag is much smaller than the ion drag and we will not consider it further.

### 3.5. Ionization and loss rates

It is noteworthy that, in fact, there are two sources of plasma in the gap between the complex plasma cloud: collisional (volume) ionization in the rf field and ion drift from the central (void) region into the gap. The collisional ionization by electrons of temperature \( \simeq 2 \text{ eV} \) in an Ar gas with ionization potential 15.8 eV gives the plasma production rate [17], \( Q_I \sim 3 \times 10^{-11} n_e n_e s^{-1} \text{ cm}^{-3} \). Assuming a global ion drift in the gap (parallel to the cloud boundaries) with a speed \( u_{ig} \), we get a ratio of the volume-to-drift plasma source terms,

\[ \frac{Q_I H}{n_i u_{ig}} \sim \frac{3 \times 10^{-11} n_e H}{u_{ig}}, \] (7)

where \( H \sim 10 \text{ cm} \) is the width of the coalescing clouds in the horizontal direction (see figure 1). This ratio is of the order of unity for gas pressure \( \simeq 0.5 \text{ mBar} \), assuming \( T_i \sim 300 \text{ K} \) and \( u_{ig} \sim v_{Ti} \). Interestingly, this ratio is independent of the gap width. Hence it should hold even for \( 2L = \Delta \), i.e., drift supply of ions is possible everywhere and only limited by absorption on the grains. In turn, the absorption length of the drifting ions can easily be estimated. It is

\[ \ell_i \sim \frac{u_{ig} \Delta^3}{\sqrt{8 \pi a^2 v_{Ti} (1 + Ze^2/T_i a)}}, \]
which, physically, is the ion flux through area $\Delta$ divided by the flux on to a microparticle (we used OML theory [8, 9])—multiplied by $\Delta$. Rewriting this in terms of the coupling parameter, $\Gamma$, and substituting typical values ($T_d \sim T_i$, $Z \sim 3 \times 10^3$, $\kappa \sim 3$) yields

$$\ell_i / \Delta \sim 3 \times 10^3 u_{ig} / (1 + 0.1 \Gamma v_{Ti}).$$

Even for extremely strong coupling, $\Gamma \sim 1000$, the absorption length of ions (and electrons) is much larger than $\Delta$—they may drift right through the system without appreciable losses. We will discuss the surface physics of complex plasmas using this approach elsewhere and will not develop it further here.

As regards the plasma loss, in addition to the conventional (ambipolar) diffusion towards the chamber walls there exists also the plasma absorption on the surface of microparticles. The corresponding loss rate is [8]

$$L_A = 2\sqrt{2\pi a^2 n_i v_{Ti} \left( 1 + \frac{Z e^2}{T_i d} \right)}.$$  

We obtain that the absorption loss rate corresponding to the surface particle density ($\Delta \simeq 0.4 \text{ mm}$) is about an order of magnitude smaller than the estimated ionization rate $Q_I$. On the other hand, in the bulk of the cloud, where the density is an order of magnitude larger, the absorption on particles may even exceed the local ionization rate.

### 4. The analysis

Applying the proposed model to the observations, we can extract from the data such parameters as the ‘free plasma’ screening length, $\lambda_D$, the typical charge on the microparticles $Z$, and the coupling strength $\Gamma$. These are all quite difficult properties to determine experimentally—nevertheless they are the most important quantities to know for the equation of state of complex plasmas.

By assembling the results derived in the previous section, we can finally rewrite the force equation (1) in the following form:

$$\dot{u} + v_{nd} u = \frac{2}{\omega_{pd}^2} \left[ \frac{Z e^2}{n_i m_d} - e^{-2L/\lambda_D} f_F (2L - \Delta) \right]$$

$$- v_{id} \omega_{pi} \sqrt{\frac{Z e^2}{n_d m_d} \frac{\lambda_s}{\lambda_D}} \left[ \frac{\lambda_s}{\lambda_D} (1 - e^{-L/\lambda_D}) - (1 - e^{-\Delta/2\lambda_D}) \right],$$  

where $\omega_{pd} = \sqrt{e^2 Z^2 n_d / m_d}$ and $\omega_{pi} = \sqrt{e^2 n_i / m_i}$ are the dust and ion plasma frequencies, respectively, and $\lambda_s/\lambda_D$ is given by equation (4). Equation (9) describes a nonlinear damped oscillator: the net force—the right-hand side of equation (9)—pushes the complex plasma clouds to coalesce and vanishes when the clouds have approached to the equilibrium distance $2L = \Delta$.

The system may be weakly damped (if gas friction is small) and perform oscillations, or be overdamped and describe the asymptotic approach to equilibrium. For the conditions of the experiment, the slowdown of the complex plasma clouds should occur on the neutral drag time scale $1/v_{nd} \sim 3 \times 10^{-3} \text{ s}$. Figure 3 shows that the slowdown actually takes about 1–2 s, implying that the system is strongly overdamped—the inertial term $\dot{u}$ can be omitted, so that the internal
pressure of microparticles drives the system in equilibrium with neutral gas friction, thus pushing the clouds to coalesce.

The estimates of the internal pressure and the ion drag forces—the first and second terms in the right-hand side of equation (9), respectively—using plasma properties which are either measured or deduced from simulations \((a = 1.7 \times 10^{-4} \text{ cm}, Z \sim 3 \times 10^3 e)\)—as from the OML theory with \(T_e \simeq 2 \text{ eV}, m_d \simeq 3 \times 10^{-11} \text{ g}, \Delta \simeq 400 \mu \text{ m}, N_{nn} \simeq 10, \kappa \simeq 3, n_i \sim 10^8 \text{ cm}^{-3}, m_i = 5 \times 10^{-23} \text{ g}, v_{Ti} = 3 \times 10^4 \text{ cm s}^{-1},\) cross-section function \(\Phi_{id} \sim 300\) suggest that the former force is an order of magnitude stronger than the latter. In principle, the experimental observations can be considered as a manifestation of such a force imbalance. Indeed, the pre-factor of the pressure term, \(\omega_{pd}^2\), does not depend explicitly on the ion density, whereas for the ion drag term we have \(\propto v_{id} \omega_{pd} \propto n_i^{3/2}\). Therefore initially, when the rf discharge power and, hence, the plasma density were much higher than at the coalescence stage, one can expect that the forces had an opposite imbalance—the ion drag force exceeded the internal pressure thus pushing the clouds apart.

Note that figure 3 immediately provides us with a rough estimate for the ‘free plasma’ screening length \(\tilde{\lambda}_D\): the slowdown and flattening of the merging clouds indicates that the electrostatic force of the mutual repulsion becomes comparable with the pressure force—thus suggesting that the magnitude of the screening factor \(e^{-L/\lambda_s}\) is comparable to unity. Hence, the separation between the two complex plasma clouds at that point may be taken to define \(\lambda_s \sim \tilde{\lambda}_D\). From figure 3, we obtain a value of the ‘free plasma’ screening length \((=\text{half the distance between the complex plasma clouds when their approach velocity changes})\) of \(\sim 1 \text{ mm}\), which corresponds to \(\sim 10 \lambda_D\). We can compare this estimate with the (more exact) solution of equation (9).

### 4.1. Free expansion phase

Initially, at large \(L\), we have the case of an isolated cloud with internal overpressure—the terms \(\propto e^{-L/\lambda_s}\) can be neglected in equation (9). The ‘free expansion’ velocity of the clouds is constant, as seen from figure 3. The least-squares fit of the measured gap width with a linear function yields the velocity \(\tilde{u} \simeq 2.7 \text{ mm s}^{-1}\) for each cloud. From equation (9), letting \(L \to \infty\) we derive the free expansion velocity

\[
\tilde{u} = N_{nn}(1 + \kappa) e^{-x} \frac{\omega_{pd}^2}{v_{nd}} \Delta. \tag{10}
\]

Substituting \(\Delta \simeq 0.4 \text{ mm}, v_{nd} \simeq 200 \text{ s}^{-1}\) and \(\kappa \simeq 3\), we derive the following charge on each microparticle:

\[Z \simeq 10^3 e\]

from which we calculate the coupling parameter,

\[\Gamma \sim 100\]

and, using equation (5), the potential of the ‘particle plane’ in the cloud, \(\sim (Ze/\Delta \kappa)(1 - e^{-x/2})^2 \sim 10^{-2} \text{ V}\).
4.2. Interactive phase

As can be seen from figure 3, the dynamics of the complex plasma clouds changes considerably once they start to interact. Whilst this is not surprising, it does give us added opportunities to determine the complex plasma properties—namely, the ‘free plasma’ screening length $\tilde{\lambda}_D$. In order to describe the slowdown stage, we can linearize equation (9) and then the equation for the cloud boundary $L(t)$ reads $\ddot{L} + s\dot{L} = s\Delta / 2$, and the boundary velocity is

$$u(t) \equiv -\dot{L}(t) = u_0 e^{-s(t-t_0)},$$  \hspace{1cm} (11)

where the slowdown rate is $s = N_{nm}(1 + \kappa)e^{-\kappa}\left(\omega_{pd}^2 / \nu_{nd}\right)\left(\Delta / \tilde{\lambda}_D\right) \equiv \tilde{u} / \tilde{\lambda}_D$ with $\tilde{u}$ defined in equation (10), and $u_0$ is the velocity magnitude at the moment $t_0$ when the transition to the interactive stage occurs. Figure 3 shows the least-squares fit of the exponential profile (11) to the experimental data. This yields the slowdown rate $s \simeq 1.85$ s$^{-1}$ and, hence, the ‘free plasma’ screening length of $\tilde{\lambda}_D \simeq 1.5$ mm, which is about 12 times larger than the (assumed) complex plasma value $\lambda_D$. Now, using equation (6) we can estimate the ‘free’ surface potential of the cloud (complex plasma sheath potential),

$$|\phi_s| \sim \left(Z e / \Delta^2\right)\tilde{\lambda}_D \sim 0.1 \text{V}.$$  

This implies that upon transition between the complex plasma cloud and free plasma, the resulting variation of the particle electrostatic energy is $\sim |Z\phi_s| \sim 100 \text{eV}$, which is similar to what was measured in [10] using the particle image velocimetry (PIV) technique. Note that the velocity magnitude from the fit is $u_0 \simeq 0.92\tilde{u} \simeq 2.5$ mm s$^{-1}$, which implies that the asymptotic expressions (10) and (11) very well represent the entire measured profile of $L(t)$.

5. Conclusion

In summary, an experiment has been performed under microgravity conditions with the PKE-Nefedov Laboratory on the ISS, where the coalescence of two complex plasma clouds was investigated. The two well-defined phases of the coalescence were observed: the ‘free expansion’ phase (the cloud boundaries approach each other with a constant velocity, due to a balance between the inner pressure gradient and the neutral gas friction acting on the particles), and the ‘interactive’ phase (the boundaries slow down when the gap between the clouds becomes small enough, and the clouds eventually merge). The slowdown occurs because of the mutual repulsion produced by the complex plasma sheaths formed at the boundaries.

We proposed a simple dynamical model which describes major physical mechanisms governing the coalescence and, by comparing results of the model with measured characteristics, we deduced principal complex plasma parameters, such as the particle charge, coupling parameter, and the effective screening length in the complex plasma sheath.

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