Charging of insulating and conducting dust grains
by flowing plasma and photoemission

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New Journal of Physics 11 (2009) 043005 (20pp)
Received 24 November 2008
Published 3 April 2009
Online at http://www.njp.org/
doi:10.1088/1367-2630/11/4/043005

Abstract. The charging of conducting or alternatively insulating dust grains in
a supersonic plasma flow with a directed photon flux is studied by the particle-
in-cell method. The electron emission modifies the charge distribution on
the grain surface and in the surrounding plasma. The charge and potential
distributions on and around a dust grain are studied for different photon fluxes
and different angles of the incident unidirectional photons with respect to the
plasma flow velocity vector. Continuous and pulsed radiations are considered.
It is shown that photoemission allows the charge on conducting grains to be
controlled, and that interactions between positively charged grains can be strong.
The charging of stationary and spinning insulating grains is discussed. The
simulations are carried out in two spatial dimensions, treating ions and electrons
as individual particles.
1. Introduction

The charging of an object in a plasma is one of the basic problems in plasma physics. The understanding of this process is important in studies of interactions between the plasma and the object, or between many objects in a plasma. The question is particularly important in studies of dusty plasmas, where a number of charged dust grains can form ordered structures, such as dust clusters, strings and crystals [1]–[3]. In dusty plasma experiments, grains are usually levitated in the sheath region above the electrode, and they are charged negatively due to the high mobility of electrons. Such grains can be exposed to an ion flow. A wake and characteristic regions of enhanced ion density (ion focus) are observed behind dust grains immersed in flowing plasmas [4]–[7]. The ion focusing and the corresponding potential enhancement are more conspicuous for supersonic flows [7], and can lead to the alignment of grains in the flow direction [8]–[12]. Analogous problems can be formulated also for larger objects moving with respect to a plasma, such as spacecraft or meteoroids [13, 14].

In a space environment, dust is often exposed to electromagnetic radiation [15]. Radiation can be directional or isotropic, either due to background radiation or scattering of directed light [16]. The situation is relevant not only for dust in space, but also for dusty surfaces of larger lunar bodies. In the latter case, the dust on the surface is charged by the plasma and directed solar radiation. It has been argued that the shadowing of light can lead to strong electric fields and transport of dust above the lunar surface [17]. In laboratory plasmas, the radiation is either due to the plasma glow or an external light source, thus it is either isotropic or directed, similarly to the space environment [18]. If the energy of incoming photons is larger than the work function of the dust surface material, the photoelectron current contributes to the net current to the grain and should be included in the charging analysis [3], [19]–[21]. In several respects, the physics of this process resembles that for electron emissive probes [22]. The differences between the two physical processes are found in the mechanisms for the electron emissions and to some extent also in the velocity distributions of the emitted electrons. Photoemission will change the total charge on the dust grain and the surface charge distributions, and it can lead to new types of interactions between dust grains. Structures comprising positively charged dust grains in a plasma in the presence of UV radiation have been discussed theoretically [23, 24], and observed in experiments [25, 26].
A theory describing the charging of dust grains with photoemission in a self-consistent way is difficult to develop. In particular, photoelectrons can modify the plasma in the vicinity of dust grains. Several theoretical studies consider simplified models [24, 27, 28]. To study more realistic problems, one should employ numerical simulations, which can account for nonlinear and other possible phenomena, and model the charging of a dust in plasma in a self-consistent way. By numerical simulations, it was demonstrated that the charge of the dust cloud in a plasma discharge can be modified by UV radiation [29]. The potential structures around a positively biased spacecraft were also studied numerically [30]. In these works, the electrons were treated as a Boltzmann distributed background. Neither of these studies considered the self-consistent charging of isolated objects.

One of our intentions here is to provide a more realistic model by including electrons (photoelectrons in particular) in the analysis. In a recent communication, we have demonstrated that UV radiation allows for an accurate control of the charge on an isolated conducting grain [31]. We have also shown that photoelectrons can modify and polarize the surrounding plasma, enabling strong interactions between positively charged conducting dust grains.

In the present paper, we study numerically, with the use of the particle-in-cell (PIC) method, the charging of conducting or alternatively insulating dust grains in a supersonic plasma flow with a directed photon flux. We analyze the charge, density and potential distributions for different fluxes and energies of photons, and for different angles between the incoming unidirectional photons and the plasma flow. Continuous as well as pulsed radiation is considered. Pulsed radiation may trigger collective phenomena in complex plasmas, such as plasma waves or oscillations of dust structures. Unidirectional radiation is relevant for dust in space exposed to the solar radiation, as for example lunar dust, or a laboratory experiment with an external radiation source.

2. Numerical code

We have modified the PIC code used in our previous studies [7], [32]–[34], by including a photon flux and the photoelectric effect [31]. We consider collisionless plasmas in a two-dimensional (2D) system in Cartesian coordinates. Both electrons and ions are treated as individual particles, with the electron to ion temperature ratio $T_e/T_i = 100$ and $T_e = 0.18 \text{eV}$. The ion to electron mass ratio is $m_i/m_e = 120$ in most of the simulations. As a control case, we analyze also results for a conducting grain with $m_i/m_e = 36720$ (to represent neon). The plasma density is $n = 10^{10} \text{m}^{-2}$, and the plasma flow velocity is $v_d = 1.5C_s$, with $C_s$ denoting the speed of sound. Because of the large thermal velocity of electrons, the plasma flow is represented solely by the ion drift at the boundaries of the simulation box.

A circular dust grain of radius of $R = 0.375$ in units of the electron Debye length $\lambda_{De}$ is placed inside a simulation box of size of $50 \times 50 \lambda_{De}^2$. The grain is assumed to be massive and immobile, except for the simulations of the spinning insulator. Initially, the grain is charged only by the collection of electrons and ions. For the perfectly insulating grain, a plasma particle hitting the dust grain surface remains at this position at all later times and contributes to the surface charge distribution. To model a small conductor in this work, the charge is redistributed equally on the dust grain surface at each time step. Such an algorithm is simple to use and is also found in other numerical studies [35, 36]. It does not, however, account for the electric dipole moment on conducting dust grains as induced by the anisotropic potential distribution.
in flowing plasmas. The equally distributed surface charge will not necessarily cancel electric fields inside the grain, and thus the algorithm is not adequate for grains larger than the Debye length or for grains of shapes different than spherical (or circular). This algorithm is different from the one used in our previous studies of the charge distribution on larger dust grains, which enforced constant potential within the dust grain \[7,32\]. The computational expenses of that algorithm were lengthy simulations and strict constraints on shapes and sizes of simulated dust grains.

A directed photon flux is switched on after approximately 40 ion plasma periods \(\tau_i\). At this time, we can assume that the surface charge distribution on a grain has reached a stationary level. The code is run typically up to 50 ion plasma periods. Three different angles between the incoming photons and the direction of the ion drift are considered: \(\alpha = \{0^\circ, 90^\circ, 180^\circ\}\). This, together with photon energies \(E_{hv}\) of 4.8, 5.5 and 7.2 eV, gives a photon flux power density \(H \in (1.9, 28.8) \text{ W m}^{-2}\). The work function \(W\) of the conducting dust grains is taken to be \(W = 4.5\) eV, which is close to work functions of many metallic materials \[24\]. For insulating grains, the photon energies \(E_{hv}\) are 10.3, 11.0 and 12.7 eV. This, together with the photon fluxes as for the case of conducting grains, gives a photon flux power density of \(H \in (4.5, 50.8) \text{ W m}^{-2}\). The work function of the insulating grain is taken as \(W = 10\) eV, which implies photoelectrons with the same energies as for the conducting case. For studies of the grain and plasma response to pulsed radiation, a single pulse of duration of \(\tau_i\), as well as a series of three pulses of durations of \(\tau_i\) and separating intervals of 0.5, 1.0 and 1.5 \(\tau_i\) are considered. For radiation pulses, we have \(\alpha = 0^\circ\).

When a photon hits the surface of the dust grain, a photoelectron of energy \(E = E_{hv} - W\) is produced at distance \(l = s v \Delta t\) from the dust grain surface, where \(s\) is an uniform random number \(s \in (0, 1]\), \(\Delta t\) is the computational time step and \(v\) is the photoelectron speed. Photoelectron velocity vectors are uniformly distributed over the hemisphere and directed away from the dust grain surface, in accordance with Lambert’s law.

To investigate the stability of the surface charge distribution on insulating grains, we also simulate spinning grains. Continuous rotation with angular velocities \(\Omega = \{\pi/2, \pi, 2\pi\}\) in units of \(\text{rad/} \tau_i\) (corresponding to the grain rotation by angles of 90°, 180°, and 360° within \(\tau_i\), respectively) is considered. The rotation starts at approximately one ion plasma period after the onset of radiation. As a control case we also rotate the grain throughout the whole simulation.

The two dimensionality of the system implies that our analysis can be directly applied for grains, or rods, extended in the direction perpendicular to the simulated plane, yet with a careful interpretation, the results are also relevant for the three dimensional (3D) case \[14\]. For point-like particles it seems well established that the differences between results in two and in three spatial dimensions can be of little significance \[37\].

### 3. Numerical results

The present section is in four parts. We consider first the charging of a conducting and then an insulating grain in the presence of continuous radiation. This problem is followed by the results from the simulations with pulsed radiation.
Table 1. The total charge \( q_t \) on a conducting dust grain for different photon energies \( E_{hv} \) and different photon fluxes \( \Phi_{hv} \) for \( \alpha = 0^\circ \), averaged over 10\( \tau_i \). The charge fluctuations \( \Delta q_t \) are also shown. The total charge \( q_t \) is normalized with the unitary 2D charge \( q_0 = e[n_0(3D)]^{1/3} \), where \( e \) is an elementary charge and \( n_0(3D) \) is the plasma density in the corresponding 3D system. The unit of \( q_0 \) is \( [q_0] = \text{C m}^{-1} \).

| \( \Phi_{hv} \) \((10^{19} \text{ m}^{-2} \text{s}^{-1}) \) | \( E_{hv} = 4.8 \text{ eV} \) | \( E_{hv} = 5.5 \text{ eV} \) |
|---|---|---|
| \( q_t \) \((q_0)\) | \( \Delta q_t \) \((q_0)\) | \( q_t \) \((q_0)\) | \( \Delta q_t \) \((q_0)\) |
| 0.0 | -755 | 30 | -755 | 30 |
| 0.25 | -163 | 30 | -168 | 29 |
| 0.50 | 19 | 33 | 12 | 31 |
| 1.25 | 251 | 45 | 273 | 50 |
| 2.50 | 795 | 61 | 1330 | 93 |

3.1. Continuous radiation, conducting grain

The charge on a conducting dust grain exposed to a continuous photon flux becomes more positive with the onset of radiation and saturates within one ion plasma period. Some of the results for a conducting grain have been presented before [31]. We include some of these also in the present work for completeness. The saturation charge on a conducting grain, summarized in table 1, depends on the flux density and photon energy. For a sufficiently high photon flux, the grain becomes positively charged. For low fluxes, the saturation charge does not depend significantly on the photon energy. For higher fluxes, high-energy photoelectrons lead to a more positive dust grain. The results for the total charge for \( E_{hv} = 7.2 \text{ eV} \) are very similar to the case of \( E_{hv} = 5.5 \text{ eV} \), and therefore they are not presented in table 1.

The floating potential on a positively charged grain for the two highest photon fluxes is shown in table 2 together with the corresponding results from analytical calculations. The analytical results are calculated for a balance of the photoemission \( i_{hv} \), ion \( i_i \) and electron \( i_e \) currents to the grain: \( i_e = i_i + i_{hv} \). For consistency, we restrict our analysis to the 2D case. The photoemission current can for this case be expressed by [1]

\[
i_{hv} = A_{hv} e \Phi_{hv(2D)} \exp \left( -\frac{e \Psi}{k T_{hv}} \right),
\]

(1)

It is assumed that the photoelectric yield and photoemission efficiency both equal unity. In the present 2D model with unidirectional photons, \( \Psi \) is the potential of the grain, \( T_{hv} \) is the photoelectron temperature, \( A_{hv} = 2R \), and \( \Phi_{hv(2D)} = c(\Phi_{hv(3D)})^{2/3} \), with the physical dimension of \( [\Phi_{hv(2D)}] = \text{m}^{-1} \text{s}^{-1} \). Subscripts (2D) and (3D) stand for two-dimensional and three-dimensional cases, respectively. The ion current to a plane surface segment with area \( A_i \) due to singly charged ions drifting at supersonic speed \( v_d \), can be approximated by

\[
i_i = A_i n_{0(2D)} e v_d \exp \left( -\frac{e \Psi}{k T_i} \right),
\]

(2)

where we define the ion cross section for supersonic ion flow as \( A_i = 2R \). The ion current is consistent with the current to a probe for retarding fields [38], but we replaced the ion...
remains negatively charged, we observe an ion focusing in the wake and energy of the photons. For photon fluxes of between the wake and the undisturbed plasma, see figure 3. The ion focusing is destroyed for positively charged grains. In this case, ions are slowed down and deflected in front of the grain. Consequently, a region of enhanced ion density in front of the dust grain, where electrons are underrepresented. A surplus of electrons corresponds to a region of an enhanced ion density in front of the positively charged conducting dust grain. The polarization of the plasma is negative behind, and positive in front of the grain. The spatial extent of the rarefaction of the ion density in front of the grain increases with increasing photon flux and energy, while the shape of the enhanced density region depends on \( \alpha \) [31].

The potential around the positively charged conducting dust grain is polarized for higher photon fluxes. In figure 2, we illustrate the difference \( \delta \) between the density of Boltzmann distributed electrons that would correspond to the calculated potential and the actual electron density: 
\[
\delta = n_0 \exp[e\Phi/kT_e] - n_e.
\]
Before the onset of the photon flux, the electrons can be well approximated by the Boltzmann distribution. With photoemission, the electrons are no longer Boltzmann distributed. The largest discrepancies for conductors are associated with a surplus of electrons due to the photoelectron emission, and to a region of an enhanced ion density in front of the dust grain, where electrons are underrepresented. A surplus of electrons corresponds to a region of an enhanced ion density.

**Table 2.** The floating potential on a grain for different photon energies \( E_{hv} \) and different photon fluxes \( \Phi_{hv} \) for \( \alpha = 0^\circ \). The results from the simulations \( \Psi_{fl, sim} \) as well as from analytical calculations \( \Psi_{fl, calc} \) are shown.

| \( \Phi_{hv} \) \((10^{19} \text{ m}^{-2} \text{s}^{-1}) \) | \( E_{hv} = 4.8 \text{ eV} \) | \( E_{hv} = 5.5 \text{ eV} \) |
|---|---|---|
| \( \Psi_{fl, sim} \) \((\text{V}) \) | \( \Psi_{fl, calc} \) \((\text{V}) \) | \( \Psi_{fl, sim} \) \((\text{V}) \) | \( \Psi_{fl, calc} \) \((\text{V}) \) |
| 1.25 | 0.17 | 0.17 | 0.28 | 0.27 |
| 2.50 | 0.16 | 0.36 | 0.56 | 0.66 |

thermal velocity by \( v_d \) and neglected a numerical constant by assuming that ion velocities are unidirectional and normal to the probe surface at the sheath edge. We note that the ion current to the positively charged grain is negligible due to the small thermal velocity of ions, but nevertheless we include it in the calculations for completeness. Since the grain radius \( R \) is comparable to the electron Debye length, we use a general expression for the orbit-motion-limited current to a conducting cylinder, to calculate the electron current to the grain [38]:

\[
i_e = \frac{1}{4} A_e n_0(2D) v_{\text{she}} \frac{r_s}{R} \left[ \text{erf} \left( \sqrt{\frac{-\chi}{r_s^2/R^2 - 1}} \right) + \frac{R}{r_s} \exp(-\chi) \left( 1 - \text{erf} \left( \sqrt{\frac{-\chi r_s^2}{r_s^2 - R^2}} \right) \right) \right],
\]

where \( \chi = -e\Psi/kT_e \), \( A_e = 2\pi r_s \) and \( r_s \) is the sheath radius, which in our calculations is set to \( r_s = 3R \). We introduced the error function as \( \text{erf}(x) = 2/\sqrt{\pi} \int_0^x \exp(-y^2)\,dy \).

The density and potential distributions around a conducting dust grain depend on the flux and energy of the photons. For photon fluxes of \( \Phi_{hv} = 0.25 \times 10^{19} \text{ m}^{-2} \text{s}^{-1} \), when the grain remains negatively charged, we observe an ion focusing in the wake [7]. The ion density in the focusing region is \( n_i \approx 1.2 n_0 \), where \( n_0 \) is the undisturbed ion density far from the grain. This result is smaller than for the corresponding case without photoemission where we had \( n_i \approx 2.2 n_0 \). The ion focusing is destroyed for positively charged grains. In this case, ions are slowed down and deflected in front of the grain. Consequently, a region of enhanced ion density is formed in front of the grain, while downstream from the grain there is a distinct boundary between the wake and the undisturbed plasma, see figure 1. The spatial extent of the rarefaction in the ion density behind the conducting grain increases with increasing photon flux and energy, while the shape of the enhanced density region depends on \( \alpha \) [31].

The potential around the positively charged conducting dust grain is polarized for higher photon fluxes. In figure 2, the potential distribution around the conducting dust grain is shown for different angles of incidence of photons with the flux \( \Phi_{hv} = 2.5 \times 10^{19} \text{ m}^{-2} \text{s}^{-1} \). The potential is negative behind, and positive in front of the grain. The polarization of the plasma is most conspicuous for \( \alpha = 180^\circ \).

In figure 3, we illustrate the difference \( \delta \) between the density of Boltzmann distributed electrons that would correspond to the calculated potential and the actual electron density: 
\[
\delta = n_0 \exp[e\Phi/kT_e] - n_e.
\]
Before the onset of the photon flux, the electrons can be well approximated by the Boltzmann distribution. With photoemission, the electrons are no longer Boltzmann distributed. The largest discrepancies for conductors are associated with a surplus of electrons due to the photoelectron emission, and to a region of an enhanced ion density in front of the dust grain, where electrons are underrepresented. A surplus of electrons corresponds to...
Figure 1. The ion density around a conducting dust grain exposed to the photon flux $\Phi_\nu = 2.5 \times 10^{19} \text{ m}^{-2} \text{s}^{-1}$ of energy $E_\nu = 4.8 \text{ eV}$ averaged over nine ion plasma periods $\tau_i$. The plasma flow is in the positive $x$-direction and $\alpha = 0^\circ$. The white region corresponds to ion densities below $0.5n_0$.

negative values, a deficiency to positive values in figure 3. We note in passing that for insulators the electric dipole governs the potential in vicinity of the grain.

To validate the simulations with reduced ion mass, we also ran the code for a realistic ion mass representing neon. The results from the latter simulations are in accordance with the results obtained with the reduced ion mass. For the realistic ion mass, the total charge $q_t$ on a grain without photoemission is more negative than for simulations with reduced ion mass. The ratio of the saturation charges for the different ion masses is $q_{t,1}/q_{t,2} = 2.5$, where indices 1, 2 refer to ion masses $m_i = 36720m_e$ (for neon), and $m_i = 120m_e$, respectively. The measured charge ratio is close to the ratio

$$\frac{Q_{0,1}}{Q_{0,2}} = \frac{\ln \left( \gamma_1/2\pi + 1 \right)}{\ln \left( \gamma_2/2\pi + 1 \right)} = 2.9,$$

where $Q_0$ is the theoretical charge on a grain in a stationary plasma in a 2D system, given by

$$Q_0 = 2\pi \epsilon_0 \Psi_{\text{fl}} \frac{R}{\lambda_D} \frac{K_1(R/\lambda_D)}{K_0(R/\lambda_D)}.$$  (5)

In (5), $K_0$ and $K_1$ are modified Bessel functions, $R$ is the radius of a grain, $\gamma = m_i/m_e$, and $\Psi_{\text{fl}}$ is a floating potential of the grain, here given by

$$\Psi_{\text{fl}} = -\frac{\kappa T_e}{2e} \left[ \ln \left( \frac{\gamma}{2\pi} + 1 \right) \right].$$  (6)

In (6) it is assumed that cold ions are reaching the surface of the large conducting object at the Bohm speed. A more detailed discussion on equations (5) and (6) is given elsewhere [32].

Without photoemission, ions are streaming out of the ion focus with a wider angle for realistic ion masses as compared to the case with a reduced ion mass. This is due to different ion drift velocities for the two cases. In both cases the ion drift is $v_d = 1.5C_s$, with the speed of sound given by $C_s = \sqrt{\kappa(T_e + 5T_i/3)/m_i}$, in the plasma far from the grain. With photoemission,
Figure 2. The potential around a conducting dust grain exposed to the photon flux $\Phi_{h\nu} = 2.5 \times 10^{19} \text{ m}^{-2} \text{ s}^{-1}$ of energy $E_{h\nu} = 4.8 \text{ eV}$ for $\alpha = 0^\circ$ (a), $\alpha = 90^\circ$ (b) and $\alpha = 180^\circ$ (c). The data were averaged over a time interval of nine ion plasma periods $\tau_i$. The plasma flows in the positive $x$-direction.

the saturation charge and the wake are similar for both ion masses. The length of the wake is the same, while the width for the realistic ion mass is larger by 5%. The charge saturates within one ion plasma period for both cases. The ion plasma period for the realistic ion mass is approximately 20 times larger than for the reduced ion mass. In another simulation for grains with radius $2R$, we find that the saturation charge is approximately twice the charge value for the grain with radius of $R$.

3.2. Continuous radiation, insulating grain

The charging of an insulating grain exposed to continuous radiation differs from the conducting case. With the onset of the photon flux, we observe the development of an electric dipole moment on the grain. This moment is antiparallel to the direction of the incident photons. A saturation in the charging characteristics is observed for photon fluxes $\Phi_{h\nu} = 0.25 \times 10^{19}$ and
Figure 3. The difference $\delta$ between the density of Boltzmann distributed electrons that would correspond to the calculated potential and the actual electron density is shown for the case with (solid line) and without (dashed line) photoemission. A conducting grain is considered with $\Phi_{h\nu} = 2.5 \times 10^{19} \text{m}^{-2} \text{s}^{-1}$ and $\alpha = 0^\circ$.

Figure 4. The total charge on an insulating dust grain as a function of time for different photon fluxes and angles of photon incidence $\alpha$. Squares correspond to the photon flux $\Phi_{h\nu} = 2.5 \times 10^{19} \text{m}^{-2} \text{s}^{-1}$ and triangles to $\Phi_{h\nu} = 0.5 \times 10^{19} \text{m}^{-2} \text{s}^{-1}$. The photon energy is $E_{h\nu} = 11.0 \text{eV}$. The results are smoothed with a moving box average filter to improve the presentation.

For photon fluxes and energies high enough to change the sign of the total charge on the grain, the photoelectric effect is significant. In fact, the charge does not saturate within the timespan of our simulations. In all cases, the charging depends on the angle of incidence, see figure 4. For lower fluxes the charge is getting less negative with increasing $\alpha$. For higher fluxes the results are more uncertain since we do not observe the equilibration of charge. The reason for this may be the electric dipole moment that does not saturate for high photon fluxes.
Figure 5. The total charge on an insulating dust grain rotating by an angle $\pi$ over one ion plasma period $\tau_i$ for $\Phi_{hv} = 2.5 \times 10^{19} \text{ m}^{-2} \text{ s}^{-1}$ (a) and continuously rotating after the start of photoemission (b). In (b), continuous radiation with $\Phi_{hv} = 1.25 \times 10^{19} \text{ m}^{-2} \text{ s}^{-1}$ (squares correspond to $\Phi_{hv} = 2.5 \times 10^{19} \text{ m}^{-2} \text{ s}^{-1}$) is switched on at $t = 39\tau_i$. Triangles correspond to the dust grain spinning throughout the whole simulation. We have chosen $\alpha = 0^\circ$ for $\Omega = \pi$, $\alpha = 90^\circ$ for $\Omega = 2\pi$, and $\alpha = 180^\circ$ for $\Omega = 0.5\pi$ in units of $\tau_i^{-1}$. The results are smoothed with a moving box average filter to improve the presentation.

Since no charge equilibration is observed for strongly illuminated stationary insulating grains, we also investigated rotating grains. Rotation by an angle $\pi$ within one ion plasma period and then arresting the rotation, modifies the charging of the grain, see figure 5(a). The rotation of the grain redistributes the charge on the dust grain surface, and lowers the total charge on the grain. For $\Phi_{hv} = 1.25 \times 10^{19} \text{ m}^{-2} \text{ s}^{-1}$, after arresting the grain rotation, the charge becomes more negative for $\alpha = 0^\circ$ and $90^\circ$, while for $\alpha = 180^\circ$ a quadrupole moment can develop in the surface charge distribution.

With steady-state rotation, the total charge oscillates in time with the mean charge value lower than on the conducting grain with the corresponding parameters for the photon flux,
The period of oscillation reflects the period of full rotation of the grain. There is little difference in the grain charging characteristics for different starting conditions of the rotation of the grain. For a grain spinning throughout the whole simulation, the total charge before the onset of photoemission is less negative than on a stationary grain and close to the charge value on the corresponding conducting grain. The charge redistribution on spinning insulating grains tilts the electric dipole moment on the grains. The strength of the electric dipole moment oscillates together with the total charge on the grain. Simultaneously, the wake becomes asymmetric and its size oscillates in time.

3.3. Pulsed radiation, conducting grain

The charge on a conducting grain exposed to a radiation pulse is more positive during the illumination. After the pulse, the charge recovers to the previous value within approximately one ion plasma period. The charge recovery is initially fast and then continues at a slower rate. Initially, the charge recovery can be well approximated by an exponential function of the form $q = q_0 \exp(-t/\tau)$, with the time constant depending on the radiation level: $\tau = 3.45 \times 10^{-9}$ s for $\Phi_{hv} = 2.5 \times 10^{19}$ m$^{-2}$ s$^{-1}$, and $\tau = 4.56 \times 10^{-9}$ s for $\Phi_{hv} = 0.5 \times 10^{19}$ m$^{-2}$ s$^{-1}$. These time constants are comparable with the electron plasma period $\tau_e = 3.53 \times 10^{-9}$ s. This suggests that initially the charge recovery is primarily due to electrons. The time constant for $\Phi_{hv} = 0.5 \times 10^{19}$ m$^{-2}$ s$^{-1}$ is larger than $\tau_e$ because the maximum charge is close to zero in this case, and the ions contribute initially to the charge recovery. After a time interval of $2\tau_e$, the time constant $\tau$ increases and reaches $\tau \approx 0.5\tau_e$ at the end of the recovery for both cases. The charging is shown in figure 6(a) together with exponential fits. For clarity of presentation the exponential fits are shown with an offset with respect to both $x$ and $y$ axes. Approximately one ion plasma period after the switch off, a small overshoot in the charging characteristic is observed for higher photon energies. The charging for given photon fluxes depends only weakly on the photon incident angles $\alpha$.

The charging after a series of three pulses is similar to the case of a single pulse with the similar photon flux and energy. Each pulse corresponds to a peak in the charging characteristics in figure 6(b). The height of each peak does not change much with the time interval between the pulses.

The electrostatic potential around the conducting dust grain exposed to radiation pulses becomes polarized as in the case of continuous radiation. The region behind the grain remains negatively charged also between the pulses. After approximately $1.5\tau$, after the last pulse, the positive potential region in the grain wake is rebuilt: first in the vicinity, and then further away from the grain. At the same time, the region with net negative charge becomes less pronounced and moves further downstream from the dust grain, slower than the ion drift speed. During the pulses, the wake potential in the vicinity of the grain oscillates with the frequency of the pulses, see figure 7. These oscillations propagate into the wake, but are heavily damped further away from the dust grain, and diminish after the last pulse, see again figure 7.

Between the pulses, the ion density behind a conducting dust grain relaxes towards the ion density distribution characteristic for negatively charged grains. This process is terminated by the start of a successive pulse, and we see the rebuilt enhanced ion density region behind the grain moving away from the dust grain in the ion drift direction, but at a lower speed. The full recovery of the ion focus occurs after approximately one ion plasma period from the last pulse. The ion density depletion in front of the grain is located closer to the grain for $\alpha = 180^\circ$ than

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Figure 6. The total charge on a conducting dust grain exposed to a single radiation pulse (a) and to a pulse series with different time intervals between pulses $\Delta t_p$ (b). In plot (a) $\Phi_{h\nu} = 2.5 \times 10^{19} \text{m}^{-2} \text{s}^{-1}$, and the local exponential fits with different time constants $\tau$ are shown with the offset from both $x$ and $y$ axes for a better presentation. In plot (b) triangles correspond to $\Phi_{h\nu} = 0.5 \times 10^{19} \text{m}^{-2} \text{s}^{-1}$ and squares to $\Phi_{h\nu} = 2.5 \times 10^{19} \text{m}^{-2} \text{s}^{-1}$. In both cases $\alpha = 0^\circ$, and $E_{h\nu} = 5.5 \text{eV}$. The results are smoothed with a moving box average filter to improve the presentation.

for $\alpha = 0^\circ$. There is also a depletion in the electron density in the region corresponding to the ion wake originating from a positively charged grain. After the switch-off, photoelectrons are rapidly redistributed, but the depletion in the region corresponding to the wake is present until the wake is filled with ions.

3.4. Pulsed radiation, insulating grain

For an insulating grain exposed to pulsed radiation we do not encounter problems with the charge saturation. An insulating dust grain exposed to a single radiation pulse has charging
Figure 7. The potential variations at different distances $\Delta x$ from the rear of the conducting grain for $y = 25.8$ in units of $\lambda_{De}$ exposed to the radiation pulses as a function of time. $\Phi_{hv} = 2.5 \times 10^{19} \text{ m}^{-2} \text{ s}^{-1}$, $E_{hv} = 5.0 \text{ eV}$, $\alpha = 180^\circ$ and $\Delta t_p = 0.5\tau_i$. The curves are selectively displaced along $y$-axis for presentation.

characteristics similar to the conductor only for low photon fluxes and $\alpha = 0^\circ$, when the total grain charge remains negative, see figure 8(a). For the low photon flux and $\alpha = 180^\circ$, the charge does not recover within one ion plasma period after the pulse, but reaches half of the charge value prior to the pulse. Longer simulations show that the charge recovery time is approximately 20 ion plasma periods in this case.

For high photon fluxes, the charge during the pulse is positive, see figure 8(b). After a single pulse, the charge is more negative than before the pulse when $\alpha = 0^\circ$, and it is less negative when $\alpha = 180^\circ$. For a series of pulses, the charge between the subsequent pulses gets more negative for $\alpha = 0^\circ$. For $\alpha = 180^\circ$, it can also be less negative for short intervals between the pulses. In all insulating cases, the total charge after a series of pulses can be more negative than after a corresponding single pulse, and the full charge recovery takes usually several plasma periods.

The radiation pulses modify the potential and density patterns in the vicinity of the insulating grain. For photons with $\alpha = 0^\circ$, we observe an enhancement in the electric dipole moment on the grain. For photons with $\alpha = 180^\circ$, a positive surface charge is accumulated on both the front and rear sides of the dust grain. During such pulses, the electric dipole moment is parallel to the ion flow (antiparallel to the photon direction), and after the pulses, a quadrupole moment in the surface charge distribution develops, see figure 9. The quadrupole moment diminishes in time, faster for low photon fluxes and the dipole moment in the surface charge distribution as well as the electrostatic potential distribution are recovered.

After the radiation pulses, the ion focus is rebuilt in the wake behind the dust grain. For $\alpha = 0^\circ$, this region is similar to the one before the pulses, while for $\alpha = 180^\circ$, at the same time instances it is wider and weaker for low fluxes, and spatially narrower with stronger focusing for high fluxes.
4. Discussion

Photoemission provides an electron source on the irradiated side of the grain and modifies the dust grain charge. For sufficiently high photon fluxes, the charge on a conducting grain becomes positive and saturates within one ion plasma period. Positively charged conducting grains slow down and deflect incoming, flowing ions. As a result, a region of enhanced ion density forms in front of the grain, while behind the grain a substantial wake in the ion density is formed, see again figure 1. The wake in the ion density scales with the photon energy and flux, being larger for higher fluxes and energies. Hence, the wake size is proportional...
Figure 9. The potential around an insulating dust grain after a series of pulses with $\Delta t_p = 1.0 \tau_i$, $\Phi_{hv} = 2.5 \times 10^{19} \text{ m}^{-2} \text{s}^{-1}$, $E_{hv} = 11.0 \text{ eV}$ and $\alpha = 180^\circ$. The minimum in the potential is $\Psi = -12.2$ in units of $kT_e/e$ on the surface of the dust grain. Potentials lower than $\Psi = -5 kT_e/e$ are colored black.

to the charge on the grain. Photoelectrons with energies comparable to the electron thermal velocity can easily be reabsorbed on the grain surface, while with higher energies they are more likely to escape the trapping potential of the grain. This, together with the photoemission rate, which is proportional to the flux, explains the development of more positive charge on the grain for high-energy photons and high fluxes \[31\]. The angle between incoming photons and plasma flow direction has little effect on the potential distribution around conducting grains. Photoelectrons contribute in neutralizing enhanced ion density regions. The electrons penetrate into the ion wake, due to their high mobility. The resulting imbalance between ion and electron densities in the vicinity of the grain leads to polarization of the plasma, see figure 2. This allows for strong interactions between many positively charged grains in flowing plasmas. The electrons are no longer Boltzmann distributed in the vicinity of the illuminated grain.

To calculate the charge on positively charged conducting grains, we use a 2D capacitance model \[32\]. In (5), we use the simulation results for the floating potential $\Psi_0$. For photon fluxes of $\Phi_{hv} = 1.25 \times 10^{19} \text{ m}^{-2} \text{s}^{-1}$, the calculated charge is $q_t \approx 250q_0$ for both photon energies. This result is close to the data shown in table 1. For $\Phi_{hv} = 2.5 \times 10^{19} \text{ m}^{-2} \text{s}^{-1}$, we have $q_t \approx 410q_0$ for $E_{hv} = 4.8 \text{ eV}$, and $q_t \approx 732q_0$ for $E_{hv} = 5.5 \text{ eV}$, which is lower than the simulation results. However, if in equation (5) we formally substitute $\lambda_D$ by $\lambda_{De}$ for $\Phi_{hv} = 2.5 \times 10^{19} \text{ m}^{-2} \text{s}^{-1}$, then the results are close to the simulation results: $q_t \approx 734q_0$ for $E_{hv} = 4.8 \text{ eV}$, and $q_t \approx 1312q_0$ for $E_{hv} = 5.5 \text{ eV}$. This suggests that for lower photon fluxes, and low positive potentials of the grain, the ions can effectively contribute to the shielding of the grain, while for more positive potentials, the grain potential is predominantly shielded by electrons.

The analytical solution for the floating potential is in good agreement with the simulation results for low-energy photons. For high-energy photons, the analytical calculations give more positive potentials than obtained from the simulations. This is due to thermalization of the photoelectrons. The temperature of low-energy photoelectrons is higher than the mean
temperature of the background electrons, but the corresponding velocities are still within the thermal spread of the background electrons. Photoelectrons are effectively slowed down by the grain and interact with background electrons. We find that the average deceleration of the high-energy photoelectrons is 25% of their initial energy. With this correction for high-energy photons, the analytical calculations for the floating potential give values close to the numerical results.

The simulations for illuminated insulating grains differ significantly from the case of conducting grains. An electric dipole moment, which is antiparallel to the photon direction, develops on stationary insulating grains. Neither the electric dipole moment nor the charge saturate for high photon flux on such grains within the simulation time. It is suggested that the electric dipole moment governs the plasma dynamics around the grain. Since it does not saturate in the unidirectional photon flux, the charge on the grain does not reach equilibrium.

Rotation of the insulating dust grain redistributes the charge on the grain surface. Without photoemission but retaining the directed ion flow, the total charge on fast spinning grains becomes less negative. The electric dipole moment on the surface diminishes and the value of the total charge is similar to the conducting case. The electric dipole moment on the spinning dust grain is tilted by an angle with respect to the direction of radiation, and the positive charge on the irradiated side of the dust grain is being neutralized when it reaches the shadow side. Due to the depletion of the ion density in the wake, the total charge oscillates on fast spinning grains, see figure 5. With the spinning grain, the symmetry in the ion wake is destroyed near the grain surface. For sufficiently fast rotation, the redistribution of the negative charge bends ion trajectories and leads to wake erosion. This process continues until the ion density is rebuilt in the vicinity of the grain and the region of reduced ion density detaches from the grain. The electron density in the wake increases, and so does the total electron current to the grain. When the charge becomes less positive, and the electron current to the grain decreases, photoemission leads to the formation of the new wake in the ion density. The closing of the wake and the resulting oscillating total charge will occur only if the erosion of the ion wake is substantial. If the rotation of the grain is slow and the photoemission rate is high, the wake will not close and will not detach.

The present discussion considers grains with spherical, cylindrical and oblate shapes. Irregularities on the dust grain surface can lead to variations in the surface charge distribution and the wake can be perturbed also for smaller angular velocities of the grain. It is also noted that grains with inhomogeneous surface properties will spin due to angular momentum transfer from ions [39]. Such inhomogeneities can also lead to complicated surface charge distributions [32].

Due to the high inertia of the dust grains, the rotation of the grains will be slow in most experiments [40]. Hence, irregularities in the dust grain surface that give rise to a complicated surface charge distribution will be the main factor for the charge saturation on the insulating grain in the presence of directed radiation. This will be valid also for slowly spinning grains. On the other hand, the surface charge distribution on the insulating dust grain exposed to directed radiation will lead to strong electric fields within the grain. This effect will be more pronounced on grains with surface irregularities and may eventually lead to the breakup of the grain. This will be valid also for stationary plasmas and is similar to the sterilization and destruction of bacteria by means of plasma used as a source of UV radiation [41, 42].

Photoemission provides a method for controlling the charge on conducting grains both in vacuum [18] and in plasma [31]. For stationary, perfectly insulating grains, we are in our

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simulation only able to achieve total charge saturation in cases where the total grain charge remains negative. The charging depends on the photon incidence angle. The charge oscillates for fast spinning insulating grains, however.

After pulses of radiation, the charge relaxes toward the value from before the onset of radiation. The recovery is initially mainly due to electrons, and then both electrons and ions. The charging of conducting grains can be approximated by an exponential function with the time constant \( \tau \) that is initially comparable with the electron plasma period and is larger later. This is in agreement with the previous results from dust grain charging simulations \([32]\). The charge on conducting grains recovers within one ion plasma period. The small overshoot in the charging characteristics at approximately one ion plasma period after the switch off can be attributed to the ion response due to the reduced ion mass \([32]\), and the formation of the ion focus. For insulators, the charge recovery may take up to 20 ion plasma periods. This is due to a complicated surface charge distribution on the grain. For \( \alpha = 180^\circ \), the positive charge on the rear of the grain is reduced slowly after the pulse because of the reduced density in the wake, while the front side of the grain is positively charged by the ion flow. For insulating grains, a quadrupole moment is present in the surface charge distribution for several ion plasma periods after a pulse. Consistently, the charge between and after the pulses is less negative for \( \alpha = 180^\circ \) than for other angles. For other angles, the surface charge distribution leads to a more negative total charge after the switch off as compared to the dust grain charge prior to photo-emission.

The photon fluxes considered for conductors in this work can be achieved by commercially available sources of UV radiation (e.g. low-pressure mercury lamps) \([41]\). In the case of lamps it would be necessary to collimate the light. In the case of UV lasers the energies of \( E_{hv} \in (4.8, 7.2) \) eV, corresponding to \( \lambda \in (172, 258) \) nm, can be achieved for instance by excimer lasers used in photolithography \([43]\). Photon energies in the range \( E_{hv} \in (10.3, 12.7) \) eV, corresponding to \( \lambda \in (97, 120) \) nm, are more difficult to obtain, but can still be achieved for instance by free electron lasers \([44]\). In our work we assumed the work function for the insulator to be \( W = 10 \) eV, which is similar to values from experiments with ice \([45, 46]\). However, insulating grains can have lower work functions. The work function of pure ice is \( W \approx 8.7 \) eV, and can be significantly lower if the ice contains impurities, as is often expected \([47]\). Sodium silicate glass can have a work function as low as \( W = 6 \) eV \([48]\). Therefore, on such insulators effects similar to those demonstrated in this study should be achieved by lower photon energies.

By illuminating a grain using short pulsed lasers, it is possible to modify the charge on the grain and excite potential oscillations in the wake. The corresponding waves in the surrounding plasma are heavily damped away from the grain. Other dust grains located in the wake will experience oscillations. Since the charge on the illuminated grain is determined by photoemission, the motion of the particles in the wake would provide non-invasive diagnostics for measuring the charge on the grain located in the wake of the illuminated one. On the other hand, continuous illumination of the grain placed in the wake will fix the grain charge and allow for accurate study of the wake of other grains.

The results presented here do not depend on the ion to electron mass ratio, except for the saturation charge on the dust grain without photoemission and the charging rate. However, the charge on the grain with photoemission is less dependent on the ion to electron mass ratio. This is because the photoemission current does not depend on the ion mass, and the ion current to positively charged grains is negligible.
There are certain limitations of our model. The radiation in the present study was assumed to be unidirectional and photoelectrons to be monoenergetic. In many problems in laboratory plasmas, UV radiation will be more isotropic due to the plasma glow and light scattering, while photoelectron energies will be statistically distributed. Isotropic radiation will cause more homogeneous distributions of the photoelectrons and the grain surface charge. We considered perfectly insulating and perfectly conducting dust grains. Finite conductivity due to impurities and resistivity can modify the results, especially for insulators. These issues were not considered in this work.

Two dimensionality implies that our results are directly relevant for infinitely long rods, and qualitatively relevant for grains or objects extended in the direction perpendicular to the simulated plane (e.g. elongated grains or cylindrical satellites). We expect that for spherical grains in a 3D system, the plasma dynamics around the grain will lead to stronger ion focusing for negatively charged grains, as well as to a stronger polarization of the plasma surrounding the grain and a more distinct wake behind positively charged grains. We are planning to address these questions with the 3D code that is now being developed.

5. Conclusions

The results from numerical simulations of charging of isolated dust grains in flowing plasmas with photoemission were presented for perfectly insulating and perfectly conducting grains. By means of photoemission, the total charge on conducting grains can be effectively controlled. We do not observe charge saturation on strongly illuminated stationary insulating grains. This is most probably due to the development of a strong electric dipole moment during photoemission. For insulating grains, surface charge irregularities and rotation of the dust grain can redistribute the surface charge on the grain. Fast spinning of the grain results in oscillations of the value of the total grain charge and the density wake behind it.

During photoemission, the electrons are non-Boltzmann distributed in the vicinity of the grain. This makes a theoretical analysis of the problem difficult. The plasma is polarized, which can give rise to strong interactions between dust grains. For insulators the interactions are controlled by a strong electric dipole moment on the surface, antiparallel to the direction of radiation. After pulses of radiation, the charge, density and potential distributions relax toward the conditions prior to photoemission. The recovery takes approximately one ion plasma period for conducting grains, and several times longer for insulating grains, due to the complicated charge distributions on the dust grain surface.

By a fine adjustment of the charge with the use of photoemission, the coagulation of dust grains can be induced due to fluctuations of the charge when the total charge on a grain is small. Both continuous and pulsed radiation should allow for non-invasive diagnostics of the charge and wake structure in dusty plasma experiments.

Acknowledgments

This work was in part supported by the Norwegian Research Council, NFR, and by the Australian Research Council, ARC. Two of the authors (WJM and HLP) wish to thank Dr Jørgen Schou for useful discussions on the photoemission from insulating materials.
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