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# Electron heating during discharges driven by thermionic emission

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The heating of plasma electrons during discharges driven by thermionic emission is studied using one-dimensional particle-in-cell Monte Carlo collisions modeling that self-consistently takes the dependence of the thermionic current on the plasma parameters into account. It is found that at a gas pressure of $10^2$ Pa the electron two-stream instability is excited. As a consequence, the electrostatic plasma wave propagates from the cathode to the anode. The trapping of electrons by this wave contributes noticeably to the heating of the plasma. At a larger gas pressure, this instability is not excited. As a consequence, plasma electrons are heated only because of the generation of energetic electrons in ionization events and the scattering of emitted electrons. © *2014 AIP Publishing LLC*. [http://dx.doi.org/10.1063/1.4901571]

## I. INTRODUCTION

Gas discharges driven by thermionic emission play an important role in many technological applications.[1–3] Thermionic energy conversion devices for solar cell applications[2] and hollow cathodes for electric propulsion[3] are examples of such applications. In these discharges, electrons are emitted by a hot cathode, while ions can be either generated at the surface of the hot cathode[4] or by the ionization of the background gas by the thermo-emitted and plasma electrons.[3] The method of ion generation depends on the application and, respectively, on the device's parameters, such as pressure of the background gas, and cathode temperature.

In hollow cathode discharges, for space applications the typical ranges of cathode temperature and gas pressure are 1100–1800 K and $10^2$–$10^3$ Pa, respectively.[3,5] During the first stage of the hollow cathode discharge, thermo-emitted electrons are accelerated in the external electric field and ionize the background gas, generating plasma.[6] When dense plasma has been generated, thermo-emitted electrons are accelerated mainly in the plasma sheath formed near the cathode and propagate through the plasma.[6]

The injection of electrons from the electrodes into the plasma due to thermionic and secondary emission or propagation of the electron beam through the plasma can launch different instabilities (see, for instance, Refs. 7–12). These instabilities can significantly influence the plasma parameters, such as density, electron temperature, and sheaths thickness.

The plasma instability induced by the secondary electron emission was studied experimentally and theoretically (see, for instance, Refs. 9–11 and references therein). The conditions under which this type of instability can be induced were found. These instabilities may cause rapid changes in the plasma parameters and drive spontaneous oscillations, launching plasma waves and producing electromagnetic radiation. Ion acoustic oscillations were observed in experiments with a thermionic energy converter in the presence of an external magnetic field sufficient for electron magnetization.[4] In these experiments, electrons and ions were generated at the surface of the hot cathode having zero potential. Both species traveled from the cathode toward the collector having also zero potential. In this research, the mobility of electrons and ions was comparable, resulting in a small sheath potential near the cathode. The oscillations were obtained when the directed electron beam velocity was greater than the phase velocity of the wave and smaller than the electron thermal velocity.

The low-frequency oscillations in thermionic initiated discharges at low gas pressure ($p < 1$ Pa) were also studied in Refs. 11 and 12 where two different discharge modes were obtained. In the first mode, which is called an anode glow, the discharge glows in the space-charge limited regime, i.e., a virtual cathode forms in this mode and the ions generated near the anode move toward the cathode. These ions compensate the virtual cathode, which causes thermionic current oscillations. Thus, the low-frequency oscillations are associated with the ion transition time through the cathode-anode gap. In the second mode, when the anode voltage exceeds some critical value, a virtual cathode is not formed. This regime was found to be stable.

The excitation of plasma instabilities can significantly influence the parameters of plasma. For instance, plasma electrons, which are in resonance with the instability, can be heated through the particle-wave interaction.[7,8] The heating of electrons increases the ionization rate, which results in an increase in the density of plasma and, respectively, may lead to an increase in the plasma positive potential. The latter can result in the trapping of energetic electrons in the plasma and in a substantial increase in the current of thermo-emitted electrons. The increase in the emitted current could further increase the rate of plasma generation. This process can become uncontrollable and could lead to the device's failure. Thus, the model of a discharge driven by thermionic emission must self-consistently take into account the dependence of thermionic current on the plasma parameters.

 

In discharges driven by thermionic emission, the plasma is heated mainly by the electrons emitted from the cathode and accelerated in the cathode sheath. Emitted electrons may induce plasma instabilities or dissipate their energy through inelastic collisions with the background gas. In the latter case, an increase in plasma electron temperature occurs because of the generation of electrons having non-zero energy and the interaction between plasma electrons and electronically excited neutrals. In addition, if the ionization degree of the plasma is high, the electron-electron collisions play an important role in the energy exchange between thermo-emitted and plasma electrons.

In the present paper, the electron heating in a low-pressure discharge driven by thermionic emission is studied using the one-dimensional self-consistent particle-in-cell Monte Carlo collisions (1D PIC/MCC) model. It is found that at a background gas pressure of $10^2$ Pa, the high-frequency oscillations are excited. These oscillations significantly influence the electron heating. In contrast, oscillations are not induced when the gas pressure is increased up to $10^3$ Pa. In this case, the energy of thermo-emitted electrons is dissipated in collisions with neutrals.

## II. GENERAL CONSIDERATION

First, let us consider the qualitative analysis. In the absence of electron emission from the cathode, the potential of the plasma located between two parallel electrodes, one of which is biased, is defined as[14]

$$\varphi_p \approx \varphi_C + \frac{T_e}{2}\left(1 + \ln\left[\frac{M}{8\pi m_e}\right]\right). \quad (1)$$

Here, $\varphi_C$ is the bias voltage (cathode potential), $T_e$ and $m_e$ are the electron temperature (in eV) and mass, respectively, and $M$ is the ion mass. Thus, at a constant value of $\varphi_C$ the plasma potential depends only on $T_e$. The electron temperature is defined by the electric fields existing in the plasma and electron losses to the walls, as well as by the energy losses in inelastic collisions with the background neutrals. If the value of $T_e$ only slightly depends on pressure $p$ (in the considered range of the pressure), one can conclude that the dependence of the plasma potential on $p$ is insignificant.

The plasma density and, accordingly, the thickness of the plasma sheath near the cathode[13] depend on the gas pressure. Namely, an increase in $p$ results in an increase in the plasma density and decrease in the sheath thickness.[13] Assuming insignificant dependence of $\varphi_p$ on $p$, an increase in gas pressure leads to an increase in the electric field at the cathode surface, which is important in discharges driven by thermionic emission. In such discharges, the density of the electron current emitted from the cathode depends strongly on the electric field at the cathode surface. The emitted current is defined by the Richardson-Dushman law accounting for the Schottky effect[3]

$$J_{th} = DT^2 \exp\left(\frac{-e\varphi_0}{k_B T}\right) \cdot \exp\left(\frac{e}{k_B T}\sqrt{\frac{eE_C(x)}{4\pi\varepsilon_0}}\right). \quad (2)$$

Here, $D = 1.2 \times 10^6\,\mathrm{A\cdot m^{-2}\,K^{-2}}$ is the constant whose value, in general, is defined by the cathode material, $T$ is the cathode temperature (in K), $E_C$ is the electric field at the cathode surface, $e$ is the elementary charge, $\varepsilon_0$ is the permittivity of free space, $k_B$ is the Boltzmann constant, and $e\varphi_0$ is the cathode work function. One can see that the value of $J_{th}$ has an exponential dependence on the electric field. Therefore, even a small change in $E_C$ leads to a drastic change in $J_{th}$.

From this simplified analysis, one can conclude that an increase in the gas pressure results in an increase in the thermionic current. If the plasma potential does not depend on $p$, the emitted electrons are accelerated up to the same energy inside the sheath, which can be considered as collisionless in low-pressure discharges. In the 1D approach, these electrons form an electron beam. Taking into account the larger $J_{th}$ for larger $p$ and almost constant velocity of beam electrons at the sheath edge, one can conclude that the density of beam electrons increases when $p$ increases. Thus, an increase in gas pressure leads to an increase in the densities of both beam ($n_b$) and plasma ($n_{pl}$) electrons. However, the ratio $\alpha = n_b/n_{pl}$ is the non-linear function versus the gas pressure. This ratio implicitly depends on many variables. Therefore, it is difficult to conclude how $\alpha$ changes when $p$ increases.

It is known[7,8] that the propagation of an electron beam through the plasma can launch electron two-stream instability, which could drastically modify the plasma parameters. In the kinetic approach for collisionless plasma and unmovable ions, the kinetic dispersion relation for electrostatic waves is defined as

$$\varepsilon(\omega) = 1 + \frac{e^2}{\varepsilon_0 m_e k^2}\int \frac{k \cdot \partial f_0/\partial V}{\omega - kV} dV. \quad (3)$$

Here, $\varepsilon(\omega)$ is the dielectric permittivity of the plasma, $k$ is the wave number, $\omega$ is the frequency of plasma wave, $V$ is the electron velocity, and $f_0$ is the unperturbed electron velocity distribution function (EVDF). In the case $\omega \gg kV$, the imaginary part of (3) can be presented as[8]

$$\mathrm{Im}\varepsilon(\omega) = -\frac{\pi \cdot e^2}{\varepsilon_0 m_e k^2}\int k\frac{\partial f_0}{\partial V} \cdot \delta(\omega - kV) dV. \quad (4)$$

This equation shows that the complex part of the dielectric permittivity of the plasma is not equal to zero if the derivative of the EVDF is non-zero at velocities near to the phase velocity of perturbation. The criterion of the plasma stability derived in Ref. 8 reads

$$\frac{\partial f_0}{\partial (V)^2} \leq 0 \quad \text{for} -\infty < V < \infty. \quad (5)$$

For Maxwellian plasma, this condition is always satisfied and two-stream instability is not realized. Moreover, if the plasma waves are excited, they are damped in such plasma due to the resonant wave-particle interaction (Landau damping).[8] However, in many laboratory plasmas the processes responsible for the violation of (5) exist. For instance, gas ionization results in the generation of electrons having non-zero energy. In molecular nitrogen, the energy of an electron ejected as a result of ionization is defined as[15]

$$\varepsilon_{eject} = \varepsilon_1 \cdot \tan\left(rnd \cdot \arctan\left[\frac{0.5 \cdot (\varepsilon_e - I)}{\varepsilon_1}\right]\right). \quad (6)$$

Here, $\varepsilon_1$ is some constant, the value of which for several gases is known,[15] *rnd* is the random number in interval (0;1), and $I$ is the threshold ionization energy. Equation (6) shows that the larger is *rnd*, the larger is the energy of ejected electron, and $\varepsilon_{eject}$ can even exceed the ionization threshold. In the latter case, these energetic electrons increase the rate of gas ionization. Further, these electrons can cause the violation of (5) because of the overpopulation of some part of the EVDF and excite two-stream instability. Recently, the instabilities excited due to the violation of the EVDF from the Maxwellian distribution were reviewed by Dyatko *et al*. in Ref. 16.

Another process that can be responsible for the violation of (5) is the thermo-emitted and secondary electrons, which are accelerated in the plasma sheath and enter the plasma as a beam. In the kinetic model, the increment of instability of the beam-plasma system is[8]

$$\gamma \approx \frac{\sqrt{3}}{2}\omega_{pe}\left(\frac{\alpha}{2}\right)^{1/3} \cdot \left[1 - \frac{T_e}{2m_e V^2}\right]. \quad (7)$$

Here, $\omega_{pe}$ is the plasma frequency. One can see that the larger is the value of $\alpha$, the larger is the increment of instability. In the case of thermionic emission, the energy of beam electrons entering the plasma is much larger than $T_e$ because $e\varphi_p > T_e$. Therefore, there is an energy gap between plasma and beam electrons, and one obtains a part of the EVDF where $\partial f_0/\partial(V)^2 > 0$. Further, in this case one can neglect the second term on the right hand side of Eq. (7) and write

$$\gamma \approx \frac{\sqrt{3}}{2}\omega_{pe}\left(\frac{\alpha}{2}\right)^{1/3}. \quad (8)$$

The effect of electron-neutral collisions on the nonlinear Landau damping in low-pressure glow discharges was studied by Kaganovich[17] for the periodic oscillations of plasma electron density propagating through plasma. In this case, the nonlinear increment of instability becomes

$$\gamma_{nl} = \gamma \cdot \tanh(2\nu\tau_r), \quad (9)$$

where $\gamma$ is defined by Eq. (8). In Eq. (9), $\nu$ is the total collision frequency and $\tau_r$ is the bounce time of trapped resonance electrons. One can see that in the limit of $\nu \rightarrow 0$ the increment of instability is zero, which means that there is no electron heating in the collisionless limit. This can be understood for a single wave propagating through the plasma. For a single wave, there is no heating without collisions, because electrons acquire energy from the wave during one half of the wave period and then return it during the other half. Thus, weak collisions are needed in order for electrons to escape the resonance with the wave and be heated.

## III. RESULTS OF PARTICLE-IN-CELL MODELING

1D PIC/MCC modeling is carried out for the gas pressures $p = 10^2$ Pa and $10^3$ Pa, cathode temperature $T = 1300$ K, constant cathode voltage $\varphi_C = -50$ V, grounded anode, cathode-anode (CA) gap $d = 0.1$ cm, and cathode material work function 1.5 eV. The background gas, having constant density distribution in the CA gap, is molecular nitrogen ($N_2$) and the gas temperature is considered to be equal to the cathode temperature. This model was detailed, for instance, in Ref. 18 and it considers only the electron-neutral elastic and inelastic collisions.[19] The electron emission from the cathode is described by the Richardson-Dushman equation (2), which takes the Schottky effect into account. The secondary electron emission from the cathode and the interaction between electrons and excited atoms are not considered. In addition, electron-electron (*e-e*) collisions, which can be important for the thermalization of electrons, are not taken into account in this model. In general, these collisions can be neglected only if the ionization degree of plasma satisfies the condition[20]

$$n_e/n_g < (m_e/M) \cdot (\sigma_{en}/\sigma_{ee}). \quad (10)$$

Here, $M$ is the mass of a molecule, $n_e$ is the electron density, $n_g$ is the gas density, and $\sigma_{en}$ and $\sigma_{ee}$ are the electron-neutral and *e-e* collision cross sections, respectively. The latter is defined as

$$\sigma_{ee} = \frac{e^2 \cdot \ln\Lambda}{16\pi\varepsilon_0^2\varepsilon_e^2}, \quad (11)$$

where $\ln\Lambda \approx 10$ is the Coulomb logarithm. For $\varepsilon_e = 1$ eV in $N_2$ gas, the dominant electron-neutral collision is the elastic scattering[19] for which $\sigma_{en} \sim 10^{-15}$ cm$^2$, while from Eq. (11) one gets $\sigma_{ee} \sim 10^{-14}$ cm$^2$. One can therefore obtain for the gas pressure of $10^2$ Pa and gas temperature of 1300 K that *e-e* collisions become negligible only for the plasma in which $n_e/n_g < 10^{-6}$. For $\varepsilon_e = 100$ eV, one has $\sigma_{en} \sim 5 \times 10^{-16}$ cm$^2$, $\sigma_{ee} \sim 10^{-17}$ cm$^2$ resulting in $n_e/n_g < 10^{-3}$; i.e., the requirement (10) becomes weaker. However, a mean free path of electrons for *e-e* collisions significantly exceeds the typical size of the CA gap, which allows one to neglect this type of collision.

In the model, the electrons are separated into two groups, namely, the electrons emitted from the cathode ($n_{em}$) and plasma electrons generated in the CA gap ($n_{pl}$). This separation is done only for diagnostic purposes. The space charges of both the groups of electrons are taken into account in the Poisson equation.

In $N_2$ gas, the elastic scattering cross-section $\sigma_{el}$ exceeds the ionization cross-section $\sigma_{ion}$ up to electron energy $\varepsilon_e \sim 90$ eV.[19] Therefore, the plasma electrons, the energy of which is much smaller than 90 eV, experience mainly elastic collisions. The minimum electron mean free path $\lambda_{min}$ in $N_2$ gas for collisions with neutrals is obtained for electron energy $\varepsilon_e \approx 22$ eV. The elastic scattering cross section corresponding to this energy is $\sigma_{el} \approx 2 \times 10^{-15}$ cm$^{-2}$. For a gas pressure of $10^2$ Pa and gas temperature of 1300 K, the gas density is $n_g \approx 5.6 \times 10^{15}$ cm$^{-3}$. Thus, one has $\lambda_{min} \sim 0.09$ cm, i.e.,

$\lambda_{min} \sim d$. The results of modeling have shown the energy of the beam electrons propagating through the plasma is $\varepsilon_e \sim 96\,\mathrm{eV}$. For this energy $\sigma_{ion} \approx 2.6 \times 10^{-16}\,\mathrm{cm}^{-2} > \sigma_{el}$, and one obtains for the beam electrons $\lambda \sim 0.68\,\mathrm{cm} \gg d$. Thus, the collisionless analysis presented in Sec. II can be applied in the conditions considered in this study.

The results of 1D PIC/MCC modeling are shown in Figs. 1–3. One can see the oscillations of both the electron density and electric field. Note that Fig. 1(a) shows only the plasma electrons the density of which reaches $n_{pl} \sim 5 \times 10^{12}\,\mathrm{cm}^{-3}$. Figs. 2(b) and 2(c) show that the density of emitted electrons is smaller than the density of plasma electrons. The phase space [see Fig. 3(a)] shows that the main part of these electrons is low-energetic. However, some part of thermo-emitted electrons does not experience collisions and can reach the anode. In the 1D approach, these electrons form an electron beam.

The results of the simulations showed that the EVDF of emitted electrons, which populate the plasma electrons due to neutrals ionization, is Maxwellian. The EVDF of plasma electrons is also Maxwellian with the temperature of $T_e \approx 0.8\,\mathrm{eV}$ [see Fig. 3(c)]. The results of 1D PIC/MCC simulations showed that the emitted electrons, which lost energy in inelastic collisions, only slightly contribute to the temperature of plasma electrons. This is explained by the fact that the density of these electrons is much smaller than the density of plasma electrons [see Figs. 2(b) and 2(c)]. However, the density of beam electrons, which are shown in Figs. 2(a) and 3(a), in accordance with Eq. (8) is sufficient to initiate the two-stream instability.

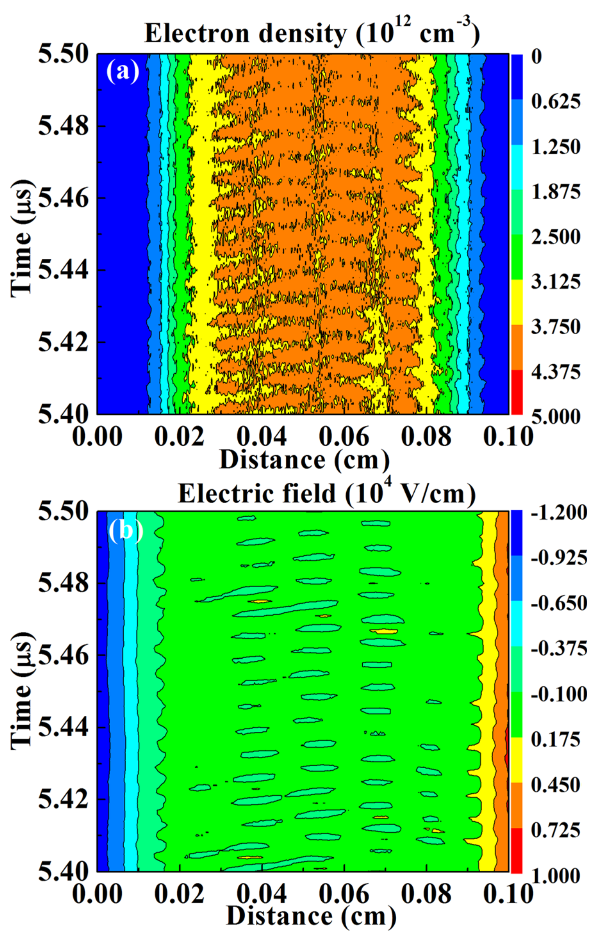


FIG. 1. Phase diagram of the electron density (a), and electric field (b).

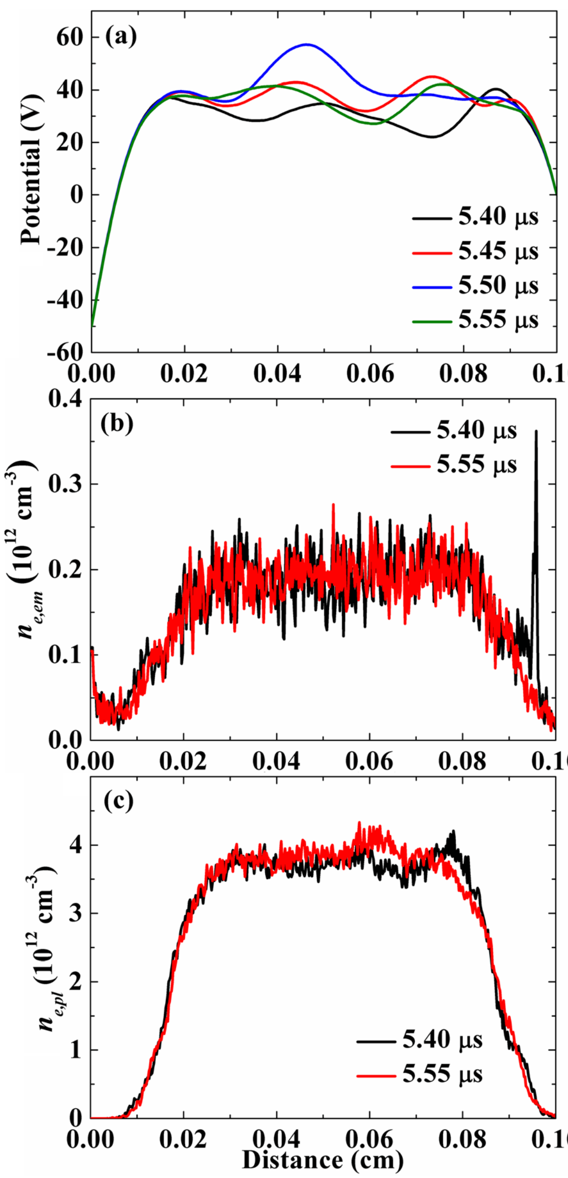


FIG. 2. Potential (a), density of emitted electrons (b), and density of plasma electrons (c), at different times.

The frequency of the electron beam density oscillations is $\omega \sim 10^9\,\mathrm{s}^{-1}$, while the plasma electron frequency is $\omega_{pe} \approx 10^{11}\ \mathrm{s}^{-1}$, i.e., $\omega_{pe} \gg \omega$. The results of the modeling give the velocity of beam electrons $V_b \approx 5.8 \times 10^8\,\mathrm{cm/s}$, and in the collisionless limit, the CA gap time-of-flight of these electrons is $\tau \sim 0.2\,\mathrm{ns}$, which is comparable with $1/\omega$. Thus, the electrostatic plasma wave is induced by the electron beam propagating through the CA gap. This wave starts from the edge of the cathode sheath and propagates toward the anode with a group velocity equal to the beam velocity, i.e., $V_g \approx 5.8 \times 10^8\,\mathrm{cm/s}$.

For the energy of beam electrons $\varepsilon_e \sim 96\,\mathrm{eV}$, the ionization is due to the dominant electron-neutral collision. The elastic scattering cross section for these electrons is 2 times smaller, while the excitation cross section is ~5 orders of

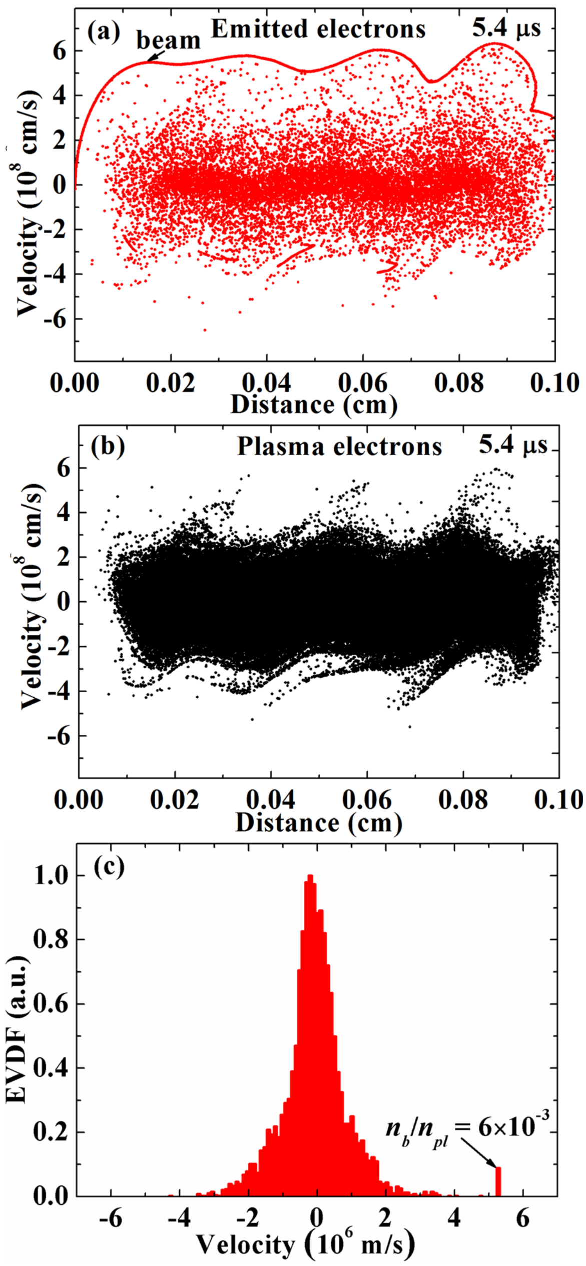


FIG. 3. (a) and (b) Phase space of emitted and plasma electrons, respectively, (c) electron velocity distribution function at 0.02 cm from the cathode.

magnitude smaller.[19] The collision frequency of beam electrons with the background gas is $\nu = n_g \cdot \sigma_{ion} \cdot V_{max} \sim 8 \times 10^8\,\mathrm{s}^{-1}$. This frequency is comparable with the wave frequency. Thus, the propagation of the wave is influenced by the inelastic collisions between beam electrons and neutrals.

The ratio between the densities of beam and plasma electrons is $\alpha \approx 0.006$. In accordance with Eq. (8), for this value of $\alpha$ and $\omega_{pe} \approx 10^{11}\,\mathrm{s}^{-1}$ the increment of instability is $\gamma \approx 10^{10}\,\mathrm{s}^{-1}$. Using $\nu = 8 \times 10^8\,\mathrm{s}^{-1}$ and $\tau_r \sim 2\pi/\omega \approx 6 \times 10^{-9}$ s in Eq. (9), one obtains $\gamma_{nl} \sim 2 \times 10^9\,\mathrm{s}^{-1}$, i.e., the nonlinear increment of instability is comparable with the collision frequency. Thus, one can conclude that the inelastic collisions between beam electrons and neutrals limit the growth of instability.

Fig. 2(a) shows that the largest potential difference obtained in the plasma bulk can reach 20 V, leading to the trapping of some fraction of plasma electrons by the propagating wave [see Fig. 3(b)]. These electrons acquire additional energy from the wave, which significantly exceeds the ionization threshold of $N_2$ gas. Thus, one obtains a local increase in the rate of gas ionization. However, the plasma electrons trapped in the potential well cannot cause the wave damping because the velocity of these electrons is much smaller than $V_g$. Therefore, the damping of the wave occurs because of beam electron-neutral collisions. Because of these collisions, beam electrons either lose their energy or change the direction of their propagation.

The sheath potential near the anode reaches ~40 V [see Fig. 2(a)]. Thus, the beam electrons experiencing ionization collisions with $\varepsilon_{eject} > I$ cannot overcome the potential barrier near the anode and become trapped in plasma. Such electrons populate the plasma electrons. This trapping also increases the temperature of plasma electrons, and as a consequence, increases the gas ionization rate.

The 1D profile of potential [see Fig. 2(a)] shows that its slope in the plasma sheath near the cathode remains almost constant, and Fig. 2(c) shows that the plasma density in the cathode sheath does not vary noticeably with time. Thus, the electric field at the cathode surface is also almost constant [see Fig. 1(b)], resulting in the constant density of thermionic current. However, small oscillations in the potential in the vicinity of the plasma sheath are obtained [see Fig. 2(a) at $x \approx 0.02$ cm]. These oscillations influence the energy and density of thermo-emitted electrons entering the plasma bulk.

A comparison of the densities of emitted electrons at two different times shows [Fig. 2(b)] the almost constant density of these electrons in the plasma bulk. However, one obtains a local increase in the electron density near the anode at time 5.4 $\mu$s, which can be explained by the plasma wave's [Fig. 2(a)] influence on the velocity of the beam electrons [Fig. 3(a)].

As discussed in Sec. II, another possible mechanism of energetic electrons generation is electrons generation in ionization events. The results of 1D PIC/MCC modeling showed the homogeneous population of the EVDF by ejected electrons. These electrons are generated mainly by beam electrons with $\varepsilon_e > 50$ eV, for which the ionization cross section is the largest and, accordingly, the mean free path is the smallest. In addition, the largest energy of ejected electrons is realized for $rnd \approx 1$ [see Eq. (6)] and the largest beam electron energy obtained in modeling is ~100 eV, resulting in the largest energy of ejected electrons being ~45 eV. Fig. 3(c) shows that there is no overpopulation of the EVDF at $\varepsilon_e \sim 45$ eV. One can therefore conclude that the obtained two-stream instability is excited by the thermo-emitted electrons.

The 1D PIC/MCC modeling was also conducted for a gas pressure of $10^3$ Pa (Fig. 4). The other parameters were not changed. The results of these simulations show the absence of the excitation of the two-stream instability. The latter can be explained by the insignificant ratio ($\alpha \sim 10^{-3}$) between the densities of the beam and plasma electrons. On the one hand, the increase in $p$ results in a decrease in the sheath thickness near the cathode. As a consequence, the

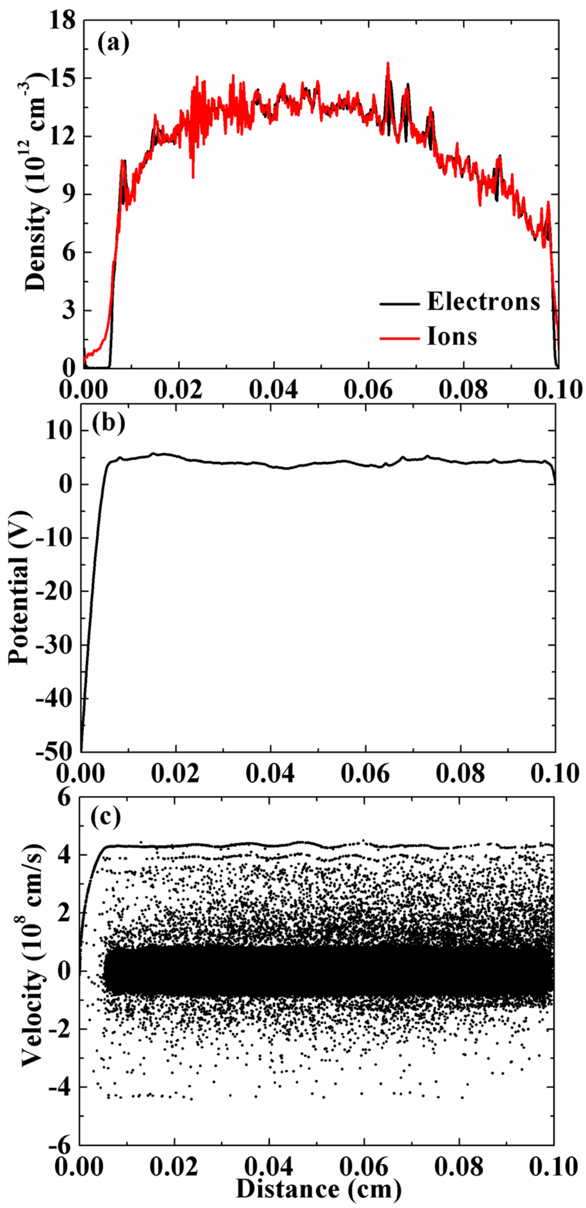


FIG. 4. (a) Electron and ion density, (b) potential, and (c) phase space of electrons obtained for $p = 10^3$ Pa.

thermionic current defined by Eq. (2) increases as compared with the case where $p = 10^2$ Pa. A comparison of Figs. 2(a) and 4(b) shows that for $p = 10^2$ Pa one obtains a larger voltage between the cathode and the plasma sheath boundary. That is, the velocity of beam electrons at the plasma sheath boundary is larger for $p = 10^2$ Pa, while the density is smaller. On the other hand, the increase in the gas pressure results in an increase in the rate of the plasma generation. This is caused by the increasing gas density and energy of beam electrons propagating through the plasma. The results of the modeling showed that $\alpha$ decreases when $p$ increases. Thus, there is no excitation of plasma waves at $p = 10^3$ Pa. Finally, from the results shown in Fig. 4 one can conclude that at $10^3$ Pa the dissipation of the energy of thermo-emitted electrons occurs mainly in inelastic collisions with neutrals. The heating of plasma is realized by the generation of energetic electrons in ionization events and is due to the transition of emitted electrons into the group of plasma electrons after the inelastic collisions.

## IV. CONCLUSIONS

The electron heating in low-pressure discharge driven by thermionic emission was studied using one-dimensional particle-in-cell Monte Carlo collisions modeling. The working gas was nitrogen with uniform density distribution at gas pressures of $10^2$ Pa and $10^3$ Pa. The dependence of the thermo-emitted current on plasma parameters was calculated self-consistently using the Richardson-Dushman law taking the Schottky effect into account.

The excitation of two-stream electron instability was obtained for the gas pressure of $10^2$ Pa. This instability was excited by the electron beam, which is emitted from the hot cathode and propagates through the plasma bulk. At these conditions, the heating of plasma electrons occurred because of their trapping in plasma wave, as a result of the generation of energetic electrons in ionization events and the scattering of thermo-emitted electrons and their consequent trapping in the plasma bulk.

The modeling conducted for the larger gas pressure ($10^3$ Pa) showed that the two-stream instability is not excited. The latter is explained by the insufficient ratio between the densities of the beam and plasma electrons. Therefore, the heating of plasma electrons is realized only because of the generation of energetic electrons in ionization events and the population of plasma electrons by scattered thermo-emitted electrons.

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