Electron trapping by electric field reversal and Fermi mechanism

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We investigate the existence of electric field reversal in the negative glow of a dc discharge, its location, the width of the well trapping the electrons, the slow electrons scattering time, and the trapping time. Based on a stress-energy tensor analysis we show the inherent instability of the well. We suggest that the Fermi mechanism is a possible process for pumping out electrons from the trough, interrelated with electrostatic plasma instabilities. A power-law distribution function for trapped electrons is also obtained. Analytical expressions are derived which can be used to calculate these characteristics from geometrical dimensions and the operational parameters of the discharge.

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I. INTRODUCTION

The phenomenon of field reversal of the axial electric field in the negative glow of a dc discharge is of great importance in the physics of gas discharges, since the fraction of ions returning to the cathode depends on its existence and location and is related to plasma instabilities. Technological application of gas discharges, particularly to plasma display panels and plasma processing, needs better knowledge of the processes involved. The study of nonlocal phenomena in electron kinetics of collisional gas discharge plasma has shown that in the presence of field reversals the bulk electrons in the cathode plasma are clearly separated into two groups of slow electrons: trapped and free electrons [1]. Trapped electrons give no contribution to the current but represent the majority of the electron population.

Kolobov and Tsendin [1] have shown that the first field reversal is located near the end of the negative glow (NG), near the position (although located slightly to the cathode side) where ions density attains its greatest magnitude [2]. If the discharge length has enough extension and the pressure decrease to lower values, a second field reversal appears on the boundary between the Faraday dark space and the positive column (PC) [1,2].

Moreover, Kolobov and Tsendin explained how ions produced to the left of the first reversal location move to the cathode by ambipolar diffusion—helping to maintain the glow by secondary electron emission—and ions generated to the right of this location drift to the anode.

More recent work [3,4] presenting a comparison of experimental data and the predictions of a hybrid fluid–Monte Carlo model also supports the view that the point where the field is extrapolated to zero is practically coincident with the maximum of the emission (even when j/p^2 scaling is no longer valid). Those characteristics were experimentally observed by laser optogalvanic spectroscopy [2]. For a detailed review see also [5].

Boeuf and Pitchford [6] with a simple fluid model gave an analytical expression of the field reversal location showing its dependence solely on the cathode sheath length, the gap length, and the ionization relaxation length. They obtained as well the fraction of ions arriving at the cathode and the magnitude of the plasma maximum density.

In the present work we introduce a quite simple dielectriclike model of a plasma-sheath system. This approach has been addressed by other authors [7,8] to explain how the electrical field inversion occurs at the interface between the plasma sheath and the beginning of the negative glow. Our aim is to obtain more information about the fundamental properties related to field inversion phenomena in the frame of a dielectric model. A simple analytical dependence is obtained of the axial location where field reversal occurs in terms of macroscopic parameters. In addition, the magnitude of the minimum electric field inside the trough, the trapped well length, and the trapping time of the slow electrons into the well are obtained. Our model emphasizes in particular the description of the dielectric behavior and does not contemplate plasma chemistry and plasma-surface interactions.

The analytical results hereby obtained could be useful for hybrid fluid-particle models (e.g., Fiala *et al.* [9], Bogaerts *et al.* [10], Marić *et al.* [3,4], and Kolobov and Arslanbekov [11]), since simple criteria can be applied to accurately remove electrons from the simulations.

On the grounds of stress-energy tensor considerations the intrinsic instability of the field reversal sheath is shown. The slow electrons (carrying most of the current to the anode) distribution function is obtained assuming the Fermi [12] mechanism responsible for their acceleration from the trapping well.

II. THEORETICAL MODEL

Lets consider a plasma formed between two parallel-plate electrodes due to an applied dc electric field. We assume a planar geometry, but extension to cylindrical geometry is straightforward. The fields are calculated for a unidimensional system, being perpendicular to the electrodes and hence neglecting end effects. The applied voltage is V_a and we assume the cathode fall length is l and the negative glow+eventually the positive column extends over the length l_0, such that the total length is $L=l+l_0$. We have

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$$-V_a = lE_s + l_0 E_p, \quad (1)$$

where E_s and E_p are, respectively, the electric fields in the sheath and NG (possibly including the positive column). We assumed a constant electric field within each region, which is clearly a rough assumption. Indeed, experimental diagnostics of the electric field in the cathode fall (CF) region show a linear behavior (Aston's law) [13]. However, taking into account a linear dependence of the electric field in the cathode sheath results in a decrease of the maximum amplitude of the electric field by a factor of 2 and does not alter the location where field reversal takes place.

At the end of the cathode sheath the following boundary condition by the displacement field $\mathbf{D}$ must be verified:

$$\mathbf{n} \cdot (\mathbf{D}_p - \mathbf{D}_s) = \sigma, \quad (2)$$

where σ is the surface charge density accumulated at the boundary surface and $\mathbf{n}$ is the normal to the surface. In more explicit form,

$$\varepsilon_p E_p - \varepsilon_s E_s = \sigma. \quad (3)$$

Here, ε_s and ε_p are, respectively, the electrical permittivity of the sheath and the positive column. We have to solve the following algebraic system of equations:

$$\begin{aligned} l_0 E_p + lE_s &= -V_a, \\ \varepsilon_p E_p - \varepsilon_s E_s &= \sigma. \end{aligned} \quad (4)$$

They give the electric field strength in each region:

$$E_s = -\frac{V_a}{L}\left(1 - \alpha + \frac{l_0 \sigma}{V_a \varepsilon_s}\right)\frac{1}{1 - l\alpha/L},$$

$$E_p = -\frac{V_a}{L}\left(1 - \frac{l\sigma}{V_a \varepsilon_s}\right)\frac{1}{1 - l\alpha/L}. \quad (5)$$

Here, we introduced $\alpha = 1 - \varepsilon_p/\varepsilon_s = \omega_p^2/\nu_{en}^2$. Recall that in the dc case the absolute permittivity is given by $\varepsilon_p = \varepsilon_0(1 - \omega_p^2/\nu_{en}^2)$ and $\varepsilon_s = \varepsilon_0$, with ω_p denoting the plasma frequency and ν_{en} the electron-neutral collision frequency in the negative glow region. In fact, our assumption $\varepsilon_s = \varepsilon_0$ is plainly justified, since experiments have shown the occurrence of a significant gas heating and a corresponding gas density reduction in the cathode fall region, mainly due to symmetric charge exchanges processes which lead to an efficient conversion of electrical energy to heavy-particle kinetic energy and thus to heating [13]. Moreover, the cathode sheath is relatively empty of electrons.

In particular, notice that Eqs. (5) show the possibility to sustain a steady-state resonant discharge. Whenever $\alpha = 1 + l_0\sigma/V_a\varepsilon_0$ the field can be entirely applied to the plasma (inductive region) while being zero in the sheath (capacitive region). On the contrary, if $l\sigma/V_a\varepsilon_0 = 1$, the inverse is true (see also [14]).

Two extreme cases can be considered: (i) if $\omega_p > \nu_{en}$, implying $\varepsilon_p < 0$, meaning that $\tau_{coll} > \tau_{plasma}$; i.e., the noncollisional regime prevails; (ii) $\omega_p < \nu_{en}$, $\varepsilon_p > 0$, and then $\tau_{coll} > \tau_{plasma}$; i.e., the collisional regime dominates.

From the above Eqs. (5) we estimate that the field inversion should occur for the condition $1 - l\alpha/L < 0$, enabling us to obtain the position on the axis where field inversion should occur:

$$\frac{l}{L} = \frac{\nu_{en}^2}{\omega_p^2}. \quad (6)$$

Examining again Eqs. (5) we realize that its underlying reasoning can be generalized introducing a third region (e.g., the positive column) and thus giving evidence that at least another second field reversal region must be present the discharge, this one located at the NG-PC boundary. In fact, whenever two regions with different characteristic energies are in contact a double layer should be formed (see [2] and references therein).

Equation (6) provides a criterion for field reversal: it only occurs in the noncollisional regime; on the contrary, in the collisional regime—and to the extent of validity of this simple model—no field reversal will occurs, since the slow electron scattering time inside the well is higher than the well lifetime, and collisions (in particular, coulombian collisions) and trapping become competitive processes. Our result is consistent with other findings. In fact, a similar condition was obtained in [15] when studying the effect of electron trapping in ion-wave instability. Likewise, a self-consistent analytic model [1] has shown that at sufficiently high pressure (that is, for $l_D\mu_i/l\mu_e$, where μ_e and μ_i are, respectively, the electron and ion mobilities, l_D is a diffusion length, and l is the distance between the cathode fall distance and the maximum electron density position), field reversal is absent. As noted above, the initial assumption of a plasma formed by two regions in our dielectriclike model does not allow a direct comparison with Boeuf and Pitchford [6]. In order to be able to predict the second field reversal, we should introduce at least a third region—in fact, the positive column. Moreover, as the electron energy degrades from the sheath to the bulk plasma, the strongest field reversals should be expected to be located at the boundary sheath-nearby region, necessarily the second field reversal being associated with a weaker electric field reversal. We postpone to a future work this study. Nevertheless, there is agreement with these theories in the sense that for higher pressures and therefore a shorter ratio of the energy relaxation lengths of fast electrons to the distance between the sheath-plasma boundary and anode, the field reversal position tends to coincide with the end of the cathode sheath.

Due to the accumulation of slow electrons after a distance $\xi_c = L - l_0$, real charges accumulate on a surface separating the cathode fall region from the negative glow. Naturally, polarization charges appear on each side of this surface and a double layer is formed with a surface charge $-\sigma_1' < 0$ on the cathode side and σ_2' on the anode side. But $\sigma' = (\mathbf{P} \cdot \mathbf{n})$, $\mathbf{P} = \varepsilon_0 \chi_e \mathbf{E}$ with $\varepsilon = \varepsilon_0(1 + \chi_e)$, χ_e denoting the dimensionless quantity called electric susceptibility. As the electric displacement is the same everywhere, we have $\mathbf{D}_0 = \mathbf{D}_1 = \mathbf{D}_2$. Thus, the residual (true) surface charge in between is given by

$$\sigma = -\sigma_1' + \sigma_2'. \quad (7)$$

After a straightforward but lengthy algebraic operation we obtain

$$\sigma = \varepsilon_p V_a \frac{B}{A}, \quad (8)$$

where

$$A = L\left(-1 + \frac{\varepsilon_0 - \varepsilon_s}{\varepsilon_p}\right) + l\left(-\frac{\varepsilon_p}{\varepsilon_s} + \frac{\varepsilon_s}{\varepsilon_p}\right) \quad (9)$$

and

$$B = \frac{\varepsilon_0(\varepsilon_s - \varepsilon_p)}{\varepsilon_s \varepsilon_p}. \quad (10)$$

We can verify that σ must be equal to

$$\sigma = \alpha \frac{V_a \varepsilon_0}{2 l_0}. \quad (11)$$

Considering that $\sigma = \varepsilon_0 \chi_e E$, we obtain the minimum value of the electric field at the reversal point:

$$E_m = \frac{\omega_p^2}{\nu_{en}^2} \frac{V_a}{2 l_0 \chi_e}. \quad (12)$$

Here, $\chi_e = \varepsilon_{rw} - 1$, with ε_{rw} designating the relative permittivity of the plasma trapped in the well. From the above equation we can obtain a more practical expression for the electrical field at its minimum strength:

$$E_m = -\frac{n_{ep}}{n_{ew}} \frac{\nu_{enw}^2}{\nu_{en}^2} \frac{V_a}{2 l_0} \approx -\frac{n_{ep}}{n_{ew}} \frac{T_{ew}}{T_{ep}} \frac{V_a}{2 l_0}. \quad (13)$$

The magnitude of the reversed electric field depends on the applied voltage and the length of the negative glow l_0. This also means that without NG there is no place for field reversal, and the bigger its length, the minor the magnitude of the reversed electric field. Moreover, the magnitude of the electric field at this point depends on the density and temperature of trapped electrons. This is consistent with earlier investigations [16].

The length of the negative glow can be estimated by the free path length l_0 of the fastest electrons possessing an energy equal to the cathode potential fall value eV_a:

$$l_0 = \int_0^{eV_a} \frac{dw}{NF(w)}. \quad (14)$$

Here, w is the electron kinetic energy and $NF(w)$ is the stopping power. For example, for He, $pl_0 = 0.02 eV_a$ is estimated [1] (in cm Torr units, with V_a in volts). We denote by n_{ew} the density of trapped electrons and by T_{ew} their respective temperature. Altogether, n_{ep} and T_{ep} are, respectively, the electron density and electron temperature in the negative glow region.

By other side, we can estimate the true surface charge density accumulated on the interface of the two regions through the expression

$$\sigma = \frac{Q}{A} = -\frac{n_{ep} e A \Delta\xi}{A}. \quad (15)$$

Here, Q is the total charge over the cross sectional area where the current flows and $\Delta\xi$ is the full width of the potential well.

The accumulation of charges at the referred interface creates a double layer, and we can realize that due to the repulsive character of Coulombian interactions, negative charges accumulate at the cathode side and positive charges at the anode side. It is well known that double layers happen when there is a sudden jump in the plasma potential [16]. This is probably the physical mechanism explaining the occurrence of field reversal in a glow discharge. Moreover, this could explain why a sudden rise in the concentration of positive ions is observed immediately after the field inversion location and as well why positive ions flowing from the anode to the cathode are stopped by it. This also explains recent results obtained through self-consistent hybrid particle-fluid simulations [3,4] showing evidence of the coincidence of the point where the field is extrapolated to zero, in excellent agreement with the maximum of the plasma emissions. These calculations are consistent with experimental observation from which a laser optogalvanic signal shows a maximum, suggesting the electric field attains a maximum [2].

Our simplified assumption of one-dimensional lines of the electric field imposes necessarily a plane surface for the double layers. This assumption is good as long as the plasma potential and wall potential are well balanced. In fact, hot electrons usually diffuse rapidly to the walls of the chamber, covering them with a surface charge of negative electrons and thus distorting the field lines. The initial assumption could be relaxed, in which case the electric field lines should also acquire a transversal component. But the basic mechanism remains the same and only the magnitudes of the electric field in each region should change and not the condition for occurrence of field reversals.

A. Instability and width of the potential well

From Eqs. (11) and (15) the trapping well width is easily obtained:

$$\Delta\xi = -\frac{eV_a}{2 m l_0 \nu_{enw}^2}. \quad (16)$$

It is expected that the potential trough should have a characteristic width of the order in between the electron Debye length ($\lambda_{De} = \sqrt{\varepsilon_0 k T_e / n_e e^2}$) and the mean scattering length. Using Eq. (16), in a He plasma, and assuming $V_a = 1$ kV, $l_0 = 1$ m, and $\nu_{en} = 1.85 \times 10^9$ s^{-1} (with $T_e = 0.03$ eV) at 1 Torr ($n = 3.22 \times 10^{16}$ cm^{-3}), we estimate $\Delta\xi \approx 2.6 \times 10^{-3}$ cm, while the Debye length is $\lambda_{De} = 2.4 \times 10^{-3}$ cm. So our Eq. (16) gives a good order of magnitude for the potential width.

Table I presents the set of parameters used to obtain our estimations. We give in Table II the estimated minimum electric field inside the well. The first field reversal located at $\xi_c \approx l_{NG}$ corresponds to the maximum density $n_{ew} \gg n_{ep}$

TABLE I. Data used for $E/p=100$ V/cm/Torr. Cross sections and electron temperatures are taken from Siglo Data base, CPAT and Kinema Software, http://www.Siglo-Kinema.com

Gas	T_e (eV)	σ (10^{-16} cm^2)
Ar	8	4.0
He	35	2.0
O_2	6	4.5
N_2	4	9.0
H_2	8	6.0

[6,17]. Thus, the assumed values for the ratio of electron temperatures and densities of the trapped electrons and electrons on the NG are typical estimates.

To examine the consistency of our model, we compare with Eq. (56) from Ref. [1] (value E_m^1 presented in Table II). However, it should be remarked that Kolobov and Tsendin do not give an explicit expression for the field reversal thin region and so we have to use ours, together with the data $\mu_e=9.6\times 10^3$ cm^2 V^{-1} s^{-1}, $\mu_i=10$ cm^2 V^{-1} s^{-1}, and $T_e^{(0)}=0.1$ eV.

Turner and Hopkins [18] explain the field reversal effect in a low-pressure rf discharge as due to the collisional drag force on the electrons advancing into the sheath. They obtained a simple analytical formula with which we also compare our result (E_m^2 in Table II). Although their specific assumptions are well adequate for rf-driven plasma, we calculated the magnitude of the maximum reversed field using the same set of data for the electron-neutral collision frequency as used before and assuming the frequency of the applied voltage, $f_{rf}=13.56$ MHz. Their formula gives a large amplitude of the electric field (in reasonable agreement with space- and time-resolved electric field measurements in He and H_2 [19]). This is possibly related to their assumption of a steplike electron density leading edge with length equal to the sheath width, instead of the smaller thin region of length $\Delta\xi$ where indeed in our model the anomaly of field reversal takes place. Nevertheless, our theory is consistent with these referred works.

It can be shown that there is no finite configuration of fields and plasma that can be in equilibrium without some external stress [20]. Hence this trough is forced to be unstable and to burst electrons periodically (or in a chaotic process), releasing the trapped electrons to the main plasma. This phenomenon produces a local perturbation in the ionization rate and the electric field, giving rise to ionization waves (striations). In fact, double layers are known to be at the onset of instabilities in a plasma [16]. In the next section, we will calculate the time of trapping with a simple Brownian model.

TABLE II. Minimum electric field at reversal point and well width. Conditions: He gas, $p=1$ Torr, $l_0=1$ cm, $V_a=1$ kV, $T_{ew}/T_{ep}=0.1$, $n_{ew}/n_{ep}=10$. E_m, present work; E_m^1, from Ref. [1]; E_m^2, from Ref. [18].

E_m (V cm^{-1})	$\Delta\xi$ (cm)	E_m^1 (V cm^{-1})	E_m^2 (V cm^{-1})
$\lesssim -5.0$	2.6×10^{-1}	-0.31	-44.8

TABLE III. Comparison between theoretical and experimental cathode fall distances at $p=1$ Torr, $E/p=100$ V/cm/Torr. Experimental data are collected from Ref. [21] (with kind of cathode material).

Gas	ξ_c^{theor} (cm)	ξ_c^{expt} (cm)
Ar	7.40	0.29 (Al)
He	1.32	1.32 (Al)
H_2	0.80	0.80 (Cu)
N_2	0.45	0.31 (Al)
Ne	0.80	0.64 (Al)
O_2	0.30	0.24 (Al)

From Eq. (6) we calculate the cathode fall length for some gases. For this purpose we took He and H_2 data as reference for atomic and molecular gases, respectively. The agreement with experimental data collected from [21] is good, with the exception of Ar. Due to the Ramsauer effect, direct comparison is difficult.

In Table III we summarized the comparison of the experimental cathode fall distances to the theoretical prediction, as given by Eq. (16). Taking into account the limitations of this model these estimates are well consistent with experimental data [21].

Of course, besides theoretical models advancing an explanation for this phenomenon, more recently certain numerical models have succeeded in describing field reversal phenomena and almost a full investigation of the CF-NG boundary [3,4] since the need to incorporate essential changes with a self-consistent-field calculation coupling the behavior of all charged species through the Poisson equation was clear (e.g., [13]). The success relies mostly on the way the nonlocality of electron kinetics is handled, in particular if primary electrons kinetics (containing electrons with high energies, from $10-10^3$ eV) which are responsible for the excitation and ionization processes sustaining the plasma are properly inserted into the model.

B. Lifetime of a slow electron in the potential well

The trapped electrons most probably diffuse inside the well with a characteristic time much shorter than the lifetime of the trough. Trapping can be avoided by Coulomb collisions [15] or by the ion-wave instability, both probably one outcome of the stress-energy unbalance, as previously mentioned. We consider a simple Brownian motion model for the slow electrons in order to obtain the scattering time τ and the lifetime T of the well. A Fermi-like model [12] will allow us to obtain the slow electron energy distribution function.

Considering the slow electron jiggling within the well, the estimated scattering time is

$$\tau=\frac{(\Delta\xi)^2}{\mathcal{D}_e}. \tag{17}$$

Here, $\mathcal{D}_e$ is the electron diffusion coefficient at thermal velocities.

TABLE IV. Scattering time and trapping time in the well. The parameters are $E/N=100$ Td, $T_g=300$ K, $V_a=1$ kV, and $l_0=0.1$ m.

Gas	$\mathcal{D}_e$ (cm^2 s^{-1})[a]	ν_{enw} (s^{-1})[b]	$\Delta\xi$ (cm)	τ (s)	T (s)
Ar	2.52×10^6	8.10×10^9	1.34×10^{-3}	7.10×10^{-13}	3.97×10^{-5}
He	5.99×10^6	2.39×10^9	1.54×10^{-2}	3.95×10^{-11}	1.70×10^{-5}
N_2	6.11×10^5	6.15×10^9	2.32×10^{-3}	8.81×10^{-12}	1.64×10^{-4}
CO_2	1.70×10^6	3.60×10^9	6.78×10^{-3}	2.70×10^{-11}	5.90×10^{-5}

[a]Data obtained through resolution of the homogeneous electron Boltzmann equation with two term expansion of the distribution function in spherical harmonics [23].
[b]Same remark as in footnote a.

The fluctuations arising in the plasma are due to the breaking of the well and we can estimate the amplitude of the fluctuating field by means of Eq. (13). We obtain

$$\delta E_m = \frac{n_{ep}}{n_{ew}}\frac{\nu_{enw}^2}{\nu_{en}^2}\frac{V_a}{el_0^2}\Delta\xi. \tag{18}$$

Meanwhile, we introduce for convenience the following adimensional variable:

$$\mathcal{E}_c = \frac{\delta E_m}{E_m} = \frac{\Delta\xi}{l_0}. \tag{19}$$

In the next section, the concept of the Fermi mechanism will be incorporated into the present problem, allowing the exact treatment of trapped electrons kinetics.

C. Power-law slow electron distribution function

As slow electrons are confined by field reversal effects, some process must be at work to pull them out from the well. We attempt to explain this phenomenon suggesting that fluctuations of the electric field in the plasma (with order of magnitude of $\mathcal{E}_c$) act over electrons, giving energy to the slow ones, which collide with those irregularities in the same manner as with heavy particles. From this mechanism a gain of energy results as well a loss. This model was first advanced by Fermi [12] when developing a theory of the origin of cosmic radiation. We shall focus now on the rate at which energy is acquired.

The average energy gain per collision by the trapped electrons (in order of magnitude) is given by

$$\Delta w = \overline{U}w(t), \tag{20}$$

with $\overline{U}\cong\mathcal{E}_c^2$ and where w is their kinetic energy. After N collisions the electrons energy will be

$$w(t) = \varepsilon_t \exp\left(\frac{\overline{U}t}{\tau}\right), \tag{21}$$

with ε_t being their thermal energy, typical of slow electrons.

The time between scattering collisions is denoted by τ. Assuming a Poisson distribution $P(t)$ for electrons escaping from the trapping well, then we state

$$P(t) = \exp(-t/\tau)dt/T. \tag{22}$$

The probability distribution of the energy gained is a function of one random variable (the energy), such as

$$f_w(w)dw = P\{w < \overline{w} < w + dw\}. \tag{23}$$

This density $f_w(w)$ can be determined in terms of the density $P(t)$. Denoting by $t_1=T$ the real root of the equation $w=w(t_1=T)$, it can be readily shown that slow electrons obey in fact the power-law distribution function

$$f_w(w)dw = \frac{\tau}{\overline{U}T}\varepsilon_t^{\tau/\overline{U}T}\frac{dw}{w^{1+\tau/\overline{U}T}}. \tag{24}$$

Like many man-made and naturally occurring phenomena (e.g., earthquake magnitude, distribution of income), it is expected that the trapped electron distribution function is a power law [see Eq. (24)], and hence $1+\tau/\mathcal{E}_c^2T=n$, with $n=2-4$ as a reasonable guess. Therefore, we estimate the trapping time to be of the order

$$T \approx \frac{\tau}{\mathcal{E}_c^2 n}. \tag{25}$$

In Table IV we summarize scattering and trapping times for a few gases.

Figure 1 shows the slow electron distribution function pumped out from the well applied for two cases: Ar (solid curve) and N_2 (dashed curve). A power exponent $n=2$ was chosen. Those distributions show that the higher confining time is associated with fewer slow electrons present in the well. When the width of the well increases (from the solid to dashed curve) the scattering time becomes longer and as well the confining time—due to a decrease of the relative number of slow electrons per given energy. This mechanism of pumping slow (trapped) electrons out of from the well can possibly explain the generation of electrostatic plasma instabilities.

Note that the trapping time is, in fact, proportional to the length of the NG and inversely proportional to the electrons diffusion coefficient at thermal energies:

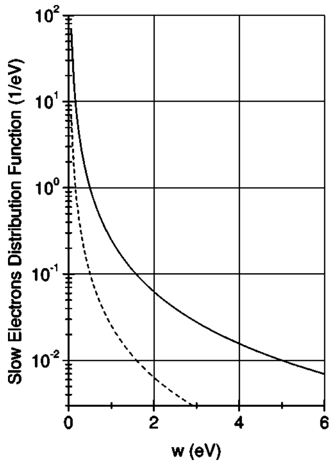

FIG. 1. Slow electrons distribution function vs energy, for the same conditions as presented in Table IV. Solid curve: Ar, dashed curve: N_2.

$$T \approx \frac{l_0^2}{\mathcal{D}_e}. \qquad (26)$$

The survival frequency of trapped electrons is $\nu_t = 1/T$. As the electron diffusion coefficient is typically higher in atomic gases, it is natural to expect plasma instabilities and waves with higher frequencies in atomic gases. This result is in agreement with a kinetic analysis of instabilities in microwave discharges (see, for example, [22]). In addition, the length of the NG will influence the magnitude of the frequencies registered by the instabilities, since wavelengths have more or less space to build up. To our best knowledge these findings are knew. Table IV summarizes the previous results for some atomic and molecular gases. The transport parameters used therefore were calculated by solving the electron Boltzmann equation, under the two-term approximation, in a steady-state Townsend discharge [23].

III. CONCLUSION

In summary, we have shown in the framework of a simple dielectric model that the magnitude of the minimum electric field (on the edge of the negative glow) depends directly on the applied voltage and is inversely proportional to the NG length.

The width of the well trapping the slow electrons is directly dependent on the applied electric field and is inversely proportional to the square of the electron-neutral collision frequency for slow electrons. It is, as well, inversely proportional to the NG length and has typically the extension of a Debye length. We state that for typical conditions of a low-pressure glow discharge, field reversal occurs whenever $\omega_p > \nu_{en}$, due to a lack of collisions necessary to pump out electrons from the well. The essence of this approach leads to the conclusion that field reversal occurs naturally at the boundary of two different regions of a glow discharge. Hence, whenever conditions to occur are verified, it should be present in a glow discharge at least two field reversals. Furthermore, the analytical expressions obtained for the scattering and trapping time of the slow electrons could be useful in self-consistent hybrid fluid-particle plasma modeling. A shortcoming of the above-expounded theory is its inability to give the exact location of the second field reversal. But this inability is compensated by a deeper understanding of the physical mechanism.

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[1] V. I. Kolobov and L. D. Tsendin, Phys. Rev. A **46**, 7837 (1992).
[2] R. A. Gottscho, A. Mitchell, G. R. Scheller, Y.-Y. Chan, and D. B. Graves, Phys. Rev. A **40**, 6407 (1989).
[3] D. Marić, K. Kutasi, G. Malović, Z. Donkó, and Z. Lj. Petrović, Eur. Phys. J. D **21**, 73 (2002).
[4] D. Marić, P. Hartmann, G. Malović, Z. Donkó, and Z. Lj. Petrović, J. Phys. D **36**, 2639 (2003).
[5] V. I. Kolobov and V. A. Godyak, IEEE Trans. Plasma Sci. **23**, 503 (1995), and references therein.
[6] J. P. Boeuf and L. C. Pitchford, J. Phys. D **28**, 2083 (1995).
[7] J. Taillet, Am. J. Phys. **37**, 423 (1969).
[8] G. S. Harmon, Am. J. Phys. **44**, 869 (1976).
[9] A. Fiala, L. C. Pitchford, and J. P. Boeuf, Phys. Rev. E **49**, 5607 (1994).
[10] A. Bogaerts and R. Gijbels, J. Anal. At. Spectrom. **13**, 721726 (1998); A. Bogaerts, A. Okhrimovskyy, and R. Gijbels, *ibid.* **17**, 10761082 (2002).
[11] V. Kolobov and R. Arslanbekov, Microelectron. Eng. **69**, 606615 (2003); A. Fedoseyev, V. Kolobov, R. Arslanbekov, and A. Przekwas, *ibid.* **69**, 577 (2003).
[12] E. Fermi, Phys. Rev. **75**, 1169 (1949).
[13] E. A. Den Hartog, D. A. Doughty, and J. E. Lawler, Phys. Rev. A **38**, 2471 (1988); D. A. Doughty, E. A. Den Hartog, and J. E. Lawler, Phys. Rev. Lett. **58**, 2668 (1987); T. J. Sommerer, W. N. G. Hitchon, and J. E. Lawler, Phys. Rev. A **39**, 6356 (1989); M. Surendra, D. B. Graves, and G. M. Jellum, *ibid.* **41**, 1112 (1990).
[14] J. Bacri and C. Saint-Martin, Nuovo Cimento B **62**, 81 (1969); S. Vacquie, J. Bacri, D. Blanc, and F. Clavery, *ibid.* **62** 69 (1969).
[15] K. Nishikawa and C.-S. Wu, Phys. Rev. Lett. **23**, 1020 (1969).
[16] P. Leung, A. Y. Wong, and B. H. Quon, Phys. Fluids **23**, 992 (1980).
[17] A. A. Kudryatsev and L. D. Tsendin, Tech. Phys. Lett. **27**, 284 (2001).
[18] M. M. Turner and M. B. Hopkins, Phys. Rev. Lett. **69**, 3511 (1992).
[19] U. Czarnetzki, D. Luggenhölscher, and H. F. Döbele, Plasma Sources Sci. Technol. **8**, 230 (1999).

[20] C. L. Longmire, *Elementary Plasma*, *Physics* (Wiley, New York, 1963), Sec. 3.7.
[21] S. C. Brown, *Basic Data of Plasma Physics* (MIT Press, Cambridge, MA, 1959).
[22] A. Shivarova, E. Tatarova, and V. Angelova, J. Phys. D **21**, 1605 (1988).
[23] M. J. Pinheiro and J. Loureiro, J. Phys. D **35**, 3077 (2002).