Plasma Physics Reports, Vol. 21, No. 5, 1995, pp. 410 - 417. Translated from Fizika Plazmy, Vol. 21, No. 5, 1995, pp. 434 - 441.

LOW-TEMPERATURE PLASMA

Microwave Discharge with Negative Ions at Moderate Pressures

I. D. Kaganovich

St. Petersburg State Technical University, Politekhnicheskaya ul. 29, St. Petersburg, 195251 Russia

Received June 6, 1994

Abstract – The properties of an electrode boundary layer in a microwave discharge at a moderate pressure are analyzed for electronegative gases. The structure of the electrode boundary layer depends on the ratio between the attachment frequency of the electrons and the time of ion motion through the layer. When the attachment frequency is very high or small, the electrode boundary layer is similar to the layer in a two-component plasma. However, for highly electronegative gases, the intermediate case is more typical. This case is characterized by the fact that the attachment frequency and the inverse time of the ion motion through the layer are comparable. This corresponds to small layer thicknesses in comparison with those in electropositive gases. In this case, the layer thickness grows proportionally to the current. The jumps or peaks in the ion density arise at the plasma–layer boundary depending on the ratio between the ambipolar diffusion and the ion velocity in the layer.

1. INTRODUCTION

Microwave discharges are widely utilized in plasma technology for deposition and etching of semiconductor coatings. The properties of microwave discharges in electropositive gases are well studied [1 - 3]. At the same time, there are only a few papers [4 - 7] devoted to the analysis of microwave discharges in electronegative gases, which are, as a rule, used in industrial applications. These papers are mainly devoted to numerical simulations. Here, we develop an analytical theory of microwave discharges. This theory permits us to determine the properties of the electrode boundary layer and the profiles of charged particle densities in a plasma. A microwave discharge with negative ions is described by many parameters and, consequently, a large variety of situations are to be studied. We investigate the most typical cases and qualitative effects. For this purpose, we use some assumptions to make our problem as simple as possible and to study the discharge peculiarities qualitatively.

The derived analytical solutions can be used for comparison the results of numerical simulations. This is especially important for such a complex system as a microwave discharge with negative ions. For example, Meyyappan and Govindan [8] obtained an erroneous stationary solution because the condition of the overall balance for ions was not satisfied [see footnote 1 to equation (4b)].

We will consider the discharge at moderate pressures when the electron distribution function is governed by the local electric field, and the Townsend approximation can be used for the ionization frequency

$$\nu_i(E) = AV_{dr}P\exp(-BP/E), \quad V_{dr} = b_eE, \qquad (1)$$

where A and B are constants, b_e is the electron mobility, E is the electric field, and P is the gas pressure. Usually, the energy threshold for attachment is considerably lower than the ionization potential. Consequently, the dependence of the attachment frequency ν_a on the electric field is less pronounced than that of the frequency ν_i. For simplicity, we assume the attachment frequency to be constant. We will consider that the discharge frequency ω satisfies the following conditions:

$$\tau_i^{-1} \ll \omega \ll \tau_e^{-1}, \qquad (2)$$

where $\tau_{i,e}^{-1} = 4\pi b_{i,e} e n_e$, b_i is the ion mobility, and n_e is the electron density. For simplicity, we will assume that the mobilities of ions are equal. The time interval τ_i characterizes the ion drift through the layer (with a thickness L).

Using the Poisson equation, we estimate τ_i:

$$\tau_i \sim \frac{b_iE}{L} \sim \frac{b_iE}{E/4\pi e n_e} = 4\pi e b_i n_e.$$

If the first inequality from (2) holds in the layer, then the densities of positive p and negative n ions in the layer are weakly modulated with respect to time, and we can use the time-averaged ion equations [2]. The second inequality permits us to neglect the displacement current in the plasma volume and keep only the electron conductivity current. Also, we suppose that the Debye radius is much less than the layer thickness or, in other words, the voltage across the layers is much more than $e^{-1}T_e$, where T_e is the electron temperature. This allows us to separate the layer into the following two regions: first, the plasma phase region, where the quasi-neutrality condition $p = n + n_e$ is valid and, second, the region of the ion volume charge, where $n_e = 0$. We shell ignore the influence of the ion density oscillations on the electron density. According to the quasi-neutrality condition, the electron density perturbations δn_e (neutralized by the oscillations of the ion volume

1063-780X/95/2105-0410$10.00 © 1995 МАИК Наука / Interperiodica Publishing

charge) can be comparable to the density itself n_e due to the large value of the ratio p/n_e

$$\frac{\delta n_e}{n_e} = \frac{\delta p - \delta n}{n_e} = \frac{\delta(p-n)}{p}\frac{p}{n_e}, \quad (3)$$

although the relative perturbation of the ion densities is small [6]. Below, we will see that, in most cases, the ion density perturbations can be ignored. Therefore, in the layer, the quantity $n_e = n_e(x) = p(x) - n(x)$ is a time-independent electron density in the plasma phase.

2. THE BASIC EQUATIONS

The mentioned assumptions substantially simplify the basic equations [2]. These equations are reduced to the time-averaged ion equations and the Poisson equation. The ion equations take the form

$$\begin{aligned}\frac{d\Gamma_p}{dx} &= \frac{d}{dx}\left(Vp - Dp\left\langle\frac{dn_e}{n_e dx}\right\rangle\right) \\ &= \langle \nu_i n_e\rangle - \beta np,\end{aligned} \quad (4a)$$

$$\begin{aligned}-\frac{d\Gamma_n}{dx} &= -\frac{d}{dx}\left(Vn - Dn\left\langle\frac{dn_e}{n_e dx}\right\rangle\right) \\ &= \langle \nu_i n_e\rangle - \beta np - \nu_d n,\end{aligned} \quad (4b)$$

where $V = b_i\langle E\rangle$ is the time-averaged velocity of ions in the volume-charge phase, $D = b_i T_e/e$ is the ambipolar diffusion coefficient, β is the ion-recombination coefficient, and ν_d is the frequency of the electron detachment from the negative ion. For the sake of simplicity, we consider moderately high pressures, when microwave diffusion can be neglected [2], and ignore the ion diffusion, because $T_e \gg T_i$. The ambipolar diffusion is important only in the plasma and in a small part of the layer next to the plasma [2]. The boundary conditions for equations (4) are determined by the direction of the ion motion. In the center of the discharge, the fluxes are zero $\Gamma_n = \Gamma_p = 0$ due to the symmetry of the considered problem. At the electrode, we have $\Gamma_n = 0$ because negative ions are drawn in the plasma region, and the surface plasma recombination at this electrode is high. From the boundary conditions, it follows that the integral of the right-hand side of equation (4b) should be zero over the entire discharge.[1]

In the volume-discharge phase, the current in the layer is transported by the displacement current and the arising field is shielded by the positive ion volume charge. If the current density depends on time as $j(t) = -j\sin\omega t$, then the field in the layer is determined as follows:

$$E(x,t) = \frac{4\pi j}{\omega}(\cos\omega t - \cos z(x)), \quad (5)$$

where $z(x)$ is the current phase corresponding to the motion of a sharp plasma–layer boundary. At a certain instant, the electric field at the electrode should be small in order to emit the electrons at the electrode and to neutralize the current of the positive ions (the current of negative ions at the electrode is zero). Consequently, the phase $z = \pi$ corresponds to the electrode (we assume that x increases toward the electrode) [2]. Using z, we can write the Poisson equation in the form

$$\sin z\frac{dz}{dx} = \frac{e\omega}{j}n_e(x), \quad n_e(x) = p - n. \quad (6)$$

Here, we take into account that the quasi-neutrality condition holds in the plasma phase. The boundary conditions for the Poisson equation are the following: $x = L$ and $z = \pi$ at the electrode and $x = 0$ and $z = 0$ at the layer boundary. Using the phase z, we can easily average equation (4) over time [2]:

$$\langle E\rangle(x) = \frac{4j}{\omega}(\sin z - z\cos z), \quad (7)$$

$$\begin{aligned}\langle I\rangle &= \frac{Aj}{\pi e}\left(\frac{2n_0}{n_e}\right)^{1/2}\exp\left(-\frac{n_e}{n_0}\right) \\ &\times \operatorname{erfc}\left[\left(z(x) - \frac{\pi}{2}\right)\left(\frac{n_e}{2n_0}\right)^{1/2}\right],\end{aligned} \quad (8)$$

$$\langle \nu_a n_e\rangle = \nu_a n_e(1 - z/\pi), \quad n_0 = j/eb_eBP. \quad (9)$$

Let us consider the case when the ion density p^c and n^c in the column[2] is much larger than the electron one n_e^c ($\nu_a \gg \beta n^c$, ν_d). The plasma region or the electrode boundary regions in which $n \gg n_e$, will be referred to as the ion–ion region (IIR) and the regions in which $n \le n_e$ will be referred to as the electron–ion region (EIR). In the electrode boundary layer, the negative ions drift from the electrode toward the plasma. Because the negative-ion density near the electrode is close to zero, there should exist a region with a small negative-ion density, specifically, an EIR. Below, we will see that the intermediate region, where $n \sim n_e$, is fairly narrow, and the jumps in n and p arise there. Consequently, the introduced definitions are convenient. For simplicity, we restrict ourselves to the case, when an interelectrode gap is so large that a uniform column can exist.

3. THE STRUCTURE OF THE ELECTRODE BOUNDARY LAYER

Depending on the discharge parameters, the electrode boundary layer consists of only EIR or of both regions, namely, IIR and EIR. This depends on the growth rate of the negative-ion density governed by the attachment frequency. The IIR arises in the layer if the

[1] The mentioned condition is not satisfied in [8]. This yields an erroneous result.

[2] We consider the column to be a discharge region, in which the particle densities are uniform and the right-hand sides of equations (4) vanish.

Table 1

	$\nu_a\tau_{i1} \ll 1$	$\tau_{i1}^{-1} \ll \nu_a \ll \tau_{i2}^{-1}$	$\nu_a\tau_{i2} \gg 1$
L	$\dfrac{2j}{e\omega n_e^c}$	$\dfrac{0.63\,(4\pi)\,jb_i}{\omega\nu_a}$	$\dfrac{2j}{e\omega n_e^c}$
p_{el}	n_e^c	$\nu_a/4\pi eb_i$	p_c
U	$\dfrac{8\pi j^2}{e\omega^2 n_e^c}$	$\dfrac{0.97\times 8\pi j^2}{e\omega^2 n_e^c\nu_a\tau_{i1}}$	$\dfrac{8\pi j^2}{e\omega^2 n_e^c}$
	$\tau_{i1} = (4\pi e n_e^c b_i)^{-1}$	$\tau_{i2} = (4\pi e p^c b_i)^{-1}$	

Here, L is the layer thickness, p_{el} is the value of the positive ion density at the electrode, U is the maximum voltage across the layer, and j is the amplitude of the current density.

time of ion motion through the EIR $\tau_{i1} = (4\pi eb_i n_e^c)^{-1}$ is larger than the time ν_a^{-1} required for the ion and electron density relaxation up to the establishment of the equilibrium densities p^c, n^c, and n_e^c in the column. At first glance, it is not evident that the electron-attachment time ν_a^{-1} governs the establishment of the ion profile. However, this follows from the quasi-neutrality condition because, at a given p, the variation of n_e yields the change in the density n (in accordance with n_e). As a result, the time ν_a^{-1} governs the characteristic time of the establishment of the ion profile. Below, we will see that the electron density in the layer can significantly vary from the value on the order of n_e^c to the value $\sim p^c$. As a result, the quantity τ_i also varies from $\tau_{i1} = (4\pi eb_i n_e^c)^{-1}$ to $\tau_{i2} = (4\pi eb_i p^c)^{-1}$. Hence, the structure of the layer depends on the values of the products $\nu_a\tau_{i1}$ and $\nu_a\tau_{i2}$.

3.1. The Case of a Low Attachment Frequency

$$\nu_a\tau_{i1} \ll 1. \tag{10}$$

In this case, the negative-ion density is low over the entire layer and the theory developed by Smirnov and Tsendin [2] can be applied to a two-component plasma. Because the quantity ν_i depends strongly on the electric field determined by the electron density in the plasma phase, this density varies slightly, that is, $n_e \approx n_e^c$. The quasi-neutrality condition yields $p \simeq n_e \simeq n_e^c$. The negative-ion density in the layer

$$n(x) = \frac{\int_x^L \langle\nu_a n_e\rangle dx}{V(x)} \simeq \nu_a\tau_{i1}n_e^c \tag{11}$$

is low almost everywhere except for the small values of x, because the ion velocity $V(x)$ decreases sharply with $x \longrightarrow (+0)$, and this density grows sharply when passing from the electrode boundary layer to the plasma. The layer properties turn out to be similar to those for the discharge in electropositive gases (see the table). For example, the thickness of the layer is $L \sim 2V_{dr}/\omega$, where $V_{dr} = b_e E^c$ is the drift velocity of the electron in the column. This situation was observed in numerical calculations carried out by Boeuf [5]. Figure 1 illustrates the results of this paper. The parameter $\nu_a\tau_{i1}$ is 0.08. The electron density in the layer is shown to be close to n_e^c. Correspondingly, for a 0.6-cm-thick layer, an estimate made according to the formula from the table gives fairly good agreement, specifically, 0.45 cm.

3.2. The Case of a Very High Attachment Frequency

$$\nu_a\tau_{i2} \gg 1 \tag{12}$$

is also simple. In this case, the IIR occupies the bulk of the layer. Below, we will prove that the negative-ion density in the bulk of the layer is close to the positive ion density p^c. In a small near-electrode layer with a thickness of about $V(L)/\nu_a \ll L$, the EIR arises in which the negative-ion density is low and the electron density is on the order of p^c (see Fig. 2). Because, in the plasma phase, the quasi-neutrality condition $p^c = n + n_e$ holds, the electron density decreases with the negative-ion density growth due to the attachment.

In the IIR, the electron density is close to n_e^c because, in this region, the ionization maintaining the density $n_e \approx n_e^c$ is important. Despite the large ion charge in the EIR (about p^c), an integral volume charge in this region $\int (p-n)\,dx$ is small (on the order of $p^cV/\nu_a \ll n_e L \sim j/\omega$) by virtue of inequality (12). Hence, in the EIR, the field is slightly shielded. Let us show that the positive ion density in the layer is nearly p^c. We have $n \sim p \gg (p-n)$ for the IIR. Because the ion diffusion is small, the absolute values of the ion fluxes satisfy the similar relationships

$$\Gamma_n \sim \Gamma_p \gg \Gamma_p - \Gamma_n. \tag{13}$$

If the ratio n/p varies slower than the quantities n and p themselves, then $\frac{d}{dx}(\Gamma_p - \Gamma_n) \ll \frac{d}{dx}\Gamma_{p,n}$. This means that the equation $\frac{d}{dx}(\Gamma_p - \Gamma_n) \simeq 0$ is valid in the IIR. Substituting the expressions for the derivatives of fluxes (4) in this equation, we find

$$\langle(\nu_a + \nu_i)n_e\rangle = 2\beta pn + \nu_d n. \tag{14}$$

Because ν_i depends strongly on n_e, the density n_e in the layer differs slightly from n_e^c. Consequently, the field profile in the layer is the same as in the Subsection 3.1. It is easier to determine the ion density profiles from the equation for negative ions (4b). We can neglect the left-

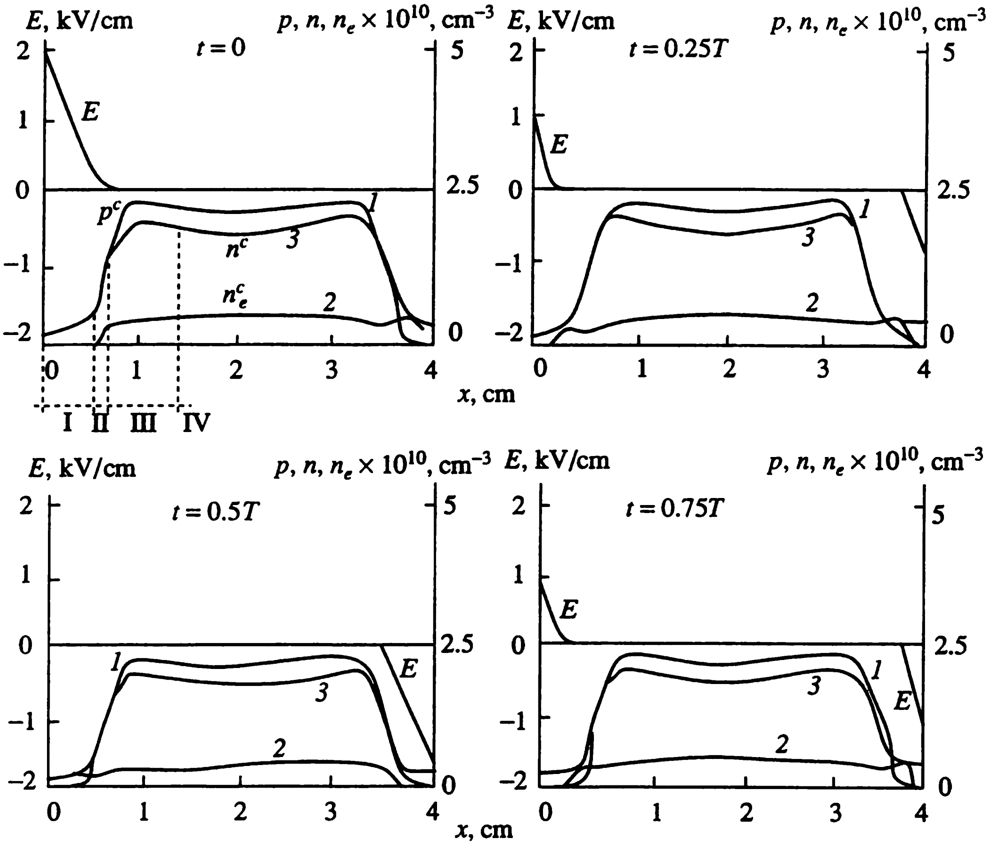

Fig. 1. The spatial profiles of the electric field and charged particle densities at four times of the microwave period for the parameters corresponding to Subsection 3.1 ($\nu_a\tau_{i1} = 0.08$), model electronegative gas based on helium [5], $f = 10$ MHz, and $U = 500$ V. (I) layer, (II) region of the jump-like transition from EIR to IIR (corresponding to the ion-density jump), (III) nonuniform plasma region, and (IV) column, (*1*) p, (*2*) n_e, (*3*) n.

hand side of this equation because $dV/dx \ll \beta p + \nu_d$ [this inequality is equivalent to (12)]. This circumference, along with expression (14), means that the quantities p and n are close to p^c and n^c. In the EIR, the positive ion density also changes insignificantly owing to the slow variation of both the ion flux and the ion velocity. The main layer parameters are listed in the table.

3.3. In the Intermediate Case, when the Inequalities

$$\tau_{i1}^{-1} \ll \nu_a \ll \tau_{i2}^{-1} \qquad (15)$$

hold, the ratio p^c/n_e^c can be very large (on the order of 10^2 - 10^3) for the strongly electronegative gases, for example, Cl_2 and SF_6. Hence, this case is typical for the discharges in such gases [4, 6 - 8]. In the cases described above, the EIR occupied either the entire layer (see Subsection 3.1) or a small part of the layer (see Subsection 3.2). Numerical simulations carried out by Shveigert [6] for the intermediate case show an almost exact coincidence of the EIR with the electrode boundary layer. The boundary between the EIR and IIR practically coincides with the plasma–layer one. The positive ion density near the electrode p_{el} and in the

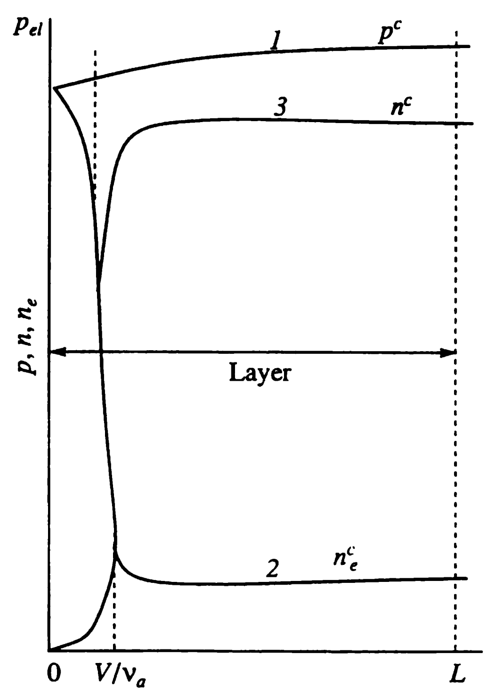

Fig. 2. The profiles of the charged-particle densities for $\nu_a\tau_{i2} \gg 1$.

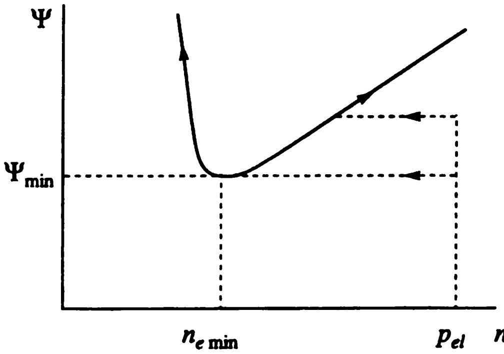

Fig. 3. The plot of the function $\Psi = (\nu_i(n_e) + \nu_a)n_e$.

layer turns out to be such that the time of the ion drift through the layer $\tau_i \sim (4\pi b_i e p_{el})^{-1}$ is comparable to ν_a^{-1} and the condition $p_{el} \sim \nu_a \tau_{i2} p^c \ll p^c$ is satisfied. This effect stems from the strong coupling of the Poison equation with the ion equations [9].

We will prove that the transition from the EIR to the IIR is really possible only near the plasma–layer boundary. Using *reductio ad absurdum*, we propose that this boundary lies somewhere inside the layer. The left inequality from (15) means that the EIR arises inside the layer. In the opposite case, the layer structure would be similar to that considered in Subsection 3.1. Because, in this case, the condition $n_e \approx p \approx n_e^c$ should be satisfied, the ion drift time (through the layer) determined by the quantity τ_{i1} is larger than the time $\nu_a^{-1} (\tau_{i1} \gg \nu_a^{-1})$ for which n becomes $\geq n_e$. This contradicts the basic assumption. On the other hand, the right inequality from (15) shows that the positive ion density p_{el} near the electrode and in the main layer region should be much less than p^c. In the EIR, we have $n_e = p \gg n_e^c$ (see below), and the ionization is exponentially low. Consequently, the positive-ion flux varies insignificantly. In the EIR with a thickness of $\sim V/\nu_a$, the average volume charge of the ions can be estimated as

$$4\pi e \int (p-n)\,dx \sim 4\pi e p_{el} V/\nu_a$$
$$= 4\pi e p_{el} b_i \langle E_{el} \rangle/\nu_a = 2\pi e p_{el} b_i E_{max}/\nu_a,$$

where E_{max} is the maximum value of the field at the electrode. The relationship between the charge and E_{max} yields

$$p_{el} < \tilde{p} \equiv p^c/\nu_a \tau_{i2} \ll p^c.$$

However, the quantity p_{el} cannot be much lower than $\tilde{p}$. Actually, when $p_{el} \ll \tilde{p}$, the average volume charge in the EIR is small, and the field in the layer should be shielded in the IIR (as in Subsection 3.2). Because $n_e \approx n_e^c \ll p_{el}$ in the IIR, the thickness of the IIR should be larger than that of the EIR, and the layer should mainly consist of the IIR (as in Subsection 3.2). Also, the positive ion density in the EIR should be almost uniform. In this case, summing the ion equations (4) yields

$$\frac{d}{dx} V n_e = \langle (\nu_i + \nu_a) n_e \rangle - 2\beta p_{el} n - n\nu_d. \qquad (16)$$

In the EIR, we have $n = p_{el} - n_e$ and equation (16) describes both the growth in n and the decrease in n_e (from the electrode toward the IIR) with a characteristic scale V/ν_a. In the IIR, the condition $p \approx n \gg n_e$ holds and equation (16) is reduced to equation (14) determining the dependence $n_e(p)$. This equation has the solution only when $p > p_{min}$. Actually, let us consider the function $\Psi(n_e) = \langle(\nu_i + \nu_a)n_e\rangle$ having two branches (Fig. 3).[3] For the large values $n_e \geq n_e^c$, we have $\nu_i \ll \nu_a$, and the function $\Psi \approx \nu_a n_e$ grows with n_e. For the small values of n_e, when $\nu_i \gg \nu_a$, the function Ψ decreases exponentially with the growth in n_e. From the form of Ψ, we find that the root of equation (14) exists only when $p > p_{min}$, where p_{min} is determined by the minimum of the function $\Psi - \Psi_{min}$. We have $\Psi_{min} \approx 1/2\nu_a n_e^c$ due to the strong dependence $\nu_i(n_e)$. When $\nu_d = 0$, this yields $p_{min} \approx p^c/\sqrt{2}$. In the general case, the quantity $p_{min} \sim p^c$.

An interesting peculiarity of equation (16) for $\nu_d = 0$ is that the point of transition from the EIR to the IIR coincides with $n_{e\,min}$ and p_{min}, where $\Psi(n_{e\,min}) = \Psi_{min}$. Actually, the left branch of the function $\Psi(n_e)$ corresponds to the IIR, in which the ionization is important. In the EIR, the quantity n_e should decrease from the large values on the order of p_{el} down to the values $n_e < n_{e\,min} \sim n_e^c$ at the left branch. It is impossible to pass from the right branch to the left one because, in the IIR, the density p increases toward the plasma and away from the left branch (see Fig. 3). Numerical simulations using the model velocity dependence on the coordinate confirm the fact that, in the region of transition from the EIR to the IIR, the positive ion density is $p^c/\sqrt{2}$.

Hence, it is possible to pass to the IIR only if the density p_{el} is on the order of p^c. This fact contradicts the initial assumption $p_{el} \ll p^c$. Consequently, the relationship $p_{el} \sim \tilde{p}$ should be satisfied and the layer consists mainly of the EIR. The velocity in the EIR varies significantly, and the transition to the IIR occurs at a low value of the ion velocity $V_{ii} \sim (p_{el}/p^c)/V(L) \ll V(L)$ near the plasma–layer boundary. In contrast to the cases considered in Subsections 3.1 and 3.2, in this case, the dependence of all the densities on x is important in the layer.

[3] According to (8) and (9), the quantity $\langle(\nu_i + \nu_a)n_e\rangle$ depends not only on n_e but also on the coordinate [by virtue of the dependence $z(x)$]. However, we can show that only the values $z > \pi/2$ can determine the boundary between the EIR and the IIR when this dependence is very weak.

From equation (16), we can find the exact value of p_{el} and the profiles of $p(x)$ and $n_e(x)$. Expressing the derivatives with respect to the coordinate through the derivatives with respect to z and using the Poisson equation, we obtain

$$\frac{dn_e V}{dz} = \frac{j\sin z}{e\omega n_e}[\langle(\nu_a+\nu_i)n_e\rangle - 2\beta pn - \nu_d n]. \quad (17)$$

In the EIR, the last three terms in equation (17) are small and the flux of positive ions to the electrode is determined by the integration of the first term in square brackets

$$\Gamma^* = p_{el}V(L) = \frac{j}{e\omega}\int_{z^*}^{\pi}\frac{\langle\nu_a n_e\rangle}{n_e}\sin z dz = \frac{\nu_a j}{e\omega}, \quad (18)$$

where $z^* \ll 1$ corresponds to the boundary between the EIR and the IIR. Note that the same expression for the negative-ion flux at the EIR boundary has been obtained in Subsection 3.1. Using (7), we substitute the expression for the velocity at the electrode in formula (18) and find the quantity p_{el}:

$$p_{el} = \nu_a/4\pi e b_i. \quad (19)$$

The system of equations (17) and (18) allows us to determine all the layer parameters. The table lists the results of integration. Figure 4 shows the results of numerical simulations carried out by Shveigert [6], which correspond to Subsection 3.3. The parameter values are $\nu_a\tau_{i1} = 9$ and $\nu_a\tau_{i1} = 0.03$. The value of p_{el} calculated according to formula (19) is 3×10^9 cm^{-3}, whereas simulations [6] give the value 2.6×10^9 cm^{-3}. In [6], the quantity ν_i was calculated as a function of the average energy. To incorporate the nonlocal effects, the equation for the average energy with allowance for the electron thermal conductivity was solved. However, the theory is also valid for these effects because, in deriving equation (18), we use only the condition of a strong dependence of ν_i on n_e.

In the considered case, the layer thickness is proportional to the current and is inversely proportional to the attachment frequency. This explains the small layer thickness in strongly electronegative gases observed in numerical calculations [4, 6, 7]. For low currents and large values of ν_a, we must take into account the finiteness of the Debye radius (as it was done in [6]), because the layer thickness can become comparable to this radius.

4. THE ION AND ELECTRON DENSITY PROFILES IN THE IIR

The negative ion flux induced by the electron attachment in the EIR should fall off to zero when passing to the column. Consequently, a nonuniform plasma region, in which the charge particle densities approach the equilibrium values in the column, should exist. When the attachment velocity is very high (see Subsection 3.2), this transition occurs at the layer near the electrode. In the cases considered in Subsections 3.1 and 3.3, the transition region can arise in the plasma or in the layer near the plasma–layer boundary, in which the velocity $V = b_i\langle E\rangle$ [see also (5) and (7)] is low.

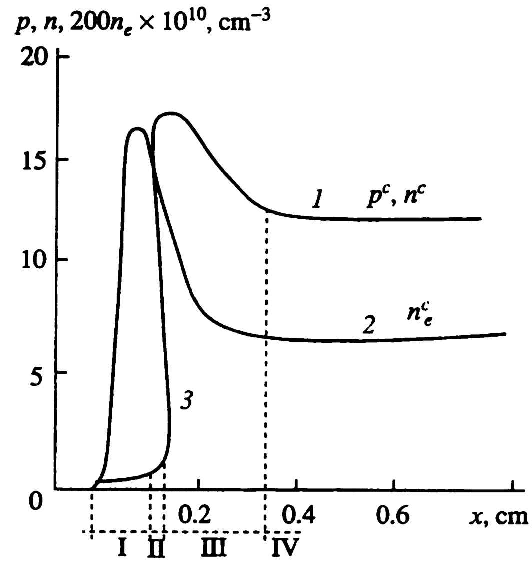

Fig. 4. The profiles of the charged-particle densities for $\tau_{i1}^{-1} \ll \nu_a \le \tau_{i2}$. The results of modeling a discharge in SF_6 [6]: $P = 0.13$ torr, $j = 2$ mA/cm^2, $f = 13.6$ MHz, $\nu_a\tau_{i1} = 9.0$, and $\nu_a\tau_{i2} = 0.03$; (*1*) positive ions, (*2*) time-averaged electron density $\langle n_e\rangle$, the dash-dotted curve corresponds to the negative-ion density. (I) Layer, (II) region of the jump-like transition from the EIR to the IIR (corresponding to the ion-density jump), (III) nonuniform plasma region, and (IV) column.

If the densities approach the equilibrium values in the plasma, then only the diffusive terms on the left-hand sides of equation (4) are nonzero. The same situation arises when the transition to the column occurs in the layer near the plasma–layer boundary, in which the diffusive terms in (4) are larger than the convective ones (in [2], this layer is referred to as the δ-layer). In this case, the velocity v_{ii} at the boundary between the EIR and the IIR is governed only by diffusion, specifically, by the flux, which is proportional to the gradient of n_e in (4). At this boundary, the diffusive flux is larger than the convective one if

$$D > \frac{\hat{\nu}_i j V(L)}{e\omega n_e}(\nu_a\tau_{i2})^{5/3}, \quad (20)$$

where $\hat{\nu}_i = \dfrac{d\ln\nu_i(n_e)}{d\ln n_e}$. In this case, the negative-ion flux decreases in the plasma and the ion-density jumps arise near the plasma–layer boundary (see the results of computations carried out in [6, 7, 10]). The nature of these jumps, which has been analyzed in [11 - 13], is related to the strong nonlinearity of the equation for negative ions when $n \gg n_e$. In the region of a jump, the parti-

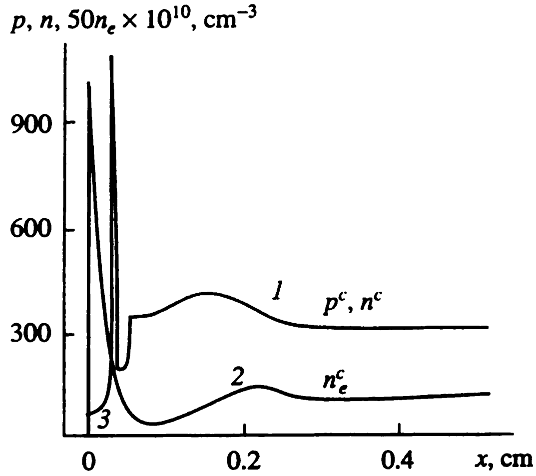

Fig. 5. The formation of the ion density peaks. The conditions are the same as in Fig. 4 except $P = 1.33$ torr and $j = 100$ mA/cm^2.

cle fluxes are practically constant, and the ion densities increase sharply toward the column (region II in Figs. 1 and 4 corresponds to the location of the jumps) and reach values comparable to those in the column. Relationship (14), coupling n and n_e, holds to the right of the jumps because $n^c \gg n_e^c$. Because n_e depends weakly on p, the equation for negative ions in this region can be simplified by neglecting the small terms on the order of $1/\hat{\nu}_i$:

$$\frac{d}{dx}\left(\frac{Dp}{n_e^c}\frac{dn_e}{dx}\right) = \nu_a n_e^c - \beta p^2 - \nu_d p. \qquad (21)$$

Equations (14) and (21) together with the negative-ion flux (18) at the plasma–layer boundary govern the properties of a nonuniform IIR, in which the negative ion flux arising in the layer falls off to zero. The negative-ion flux should fall off toward the column; hence, the right-hand side of equation (21) should be negative. The ion density should be higher than that in the column (see Figs. 1, 4, and 5) because, in the IIR, $n_e \approx n_e^c$. Formula (21) implies that the thickness of the nonuniform ion–ion plasma region is $l \sim [D/(2\beta p + \nu_d)\hat{\nu}_i]^{1/2}$, and the excess value of the ion density to the right of the jump is $\Delta p = \Gamma^*[\hat{\nu}_i/D(2\beta p + \nu_d)]^{1/2}$. For example, for the conditions of Fig. 4: $l = 5$ mm and $\Delta p = 5 \times 10^{10}$ cm^{-3}, numerical simulations [6] give $l = 2$ mm and $\Delta p = 5.2 \times 10^{10}$ cm^{-3}.

For moderate pressures (when $\hat{\nu}_i \gg 1$), the jump arises at the plasma–layer boundary because the ion velocity in this region is low and the flux Γ_n is sufficiently large and the ion density grows sharply when passing from the layer to the plasma. The case of low pressures is to be analyzed separately because the ion velocity in the plasma can be close to the velocity in a layer due to a large value of the diffusion coefficient, and the flux Γ_n is small due to the small value of ν_a. Also, the nonlocality of the electron distribution becomes important. Roughly speaking, this corresponds to the case when ν_i does not depend on n_e [14]. Accordingly, the jump can arise in the plasma at a distance on the order of $(D/\nu_a)^{1/2}$ from the plasma–layer boundary. This distance is inversely proportional to the pressure and can be large at low pressures. At this distance, the negative ion velocity falls off abruptly as in the positive column of the constant current discharge [11]. This situation was observed in numerical simulations carried out in [10].

The structure of the jump can be found by analogy with [11, 13]. Because this region is very narrow, we can ignore the variation of n_e in comparison with dn_e/dx. Also, we can neglect recombination and detachment and keep only attachment and ionization because $p, n < p^c, n^c$. The growth in the ion densities is related to the decrease in the ion velocity $V = -(D/n_e)dn_e/dx$. Using formula (16) to express the change in the velocity, we estimate the jump width:

$$\delta l \sim V\left(\frac{dV}{dx}\right)^{-1} = \frac{\Gamma^*}{n}(\nu_i + \nu_a)^{-1} = \frac{\nu_a}{\nu_i + \nu_a}\frac{p_{el}}{n}L < \frac{p_{el}}{n}L \sim \frac{p_{el}}{p^c}L. \qquad (22)$$

The jump width (22) is much less than the layer thickness because the ratio p_{el}/p^c is small. For example, under the conditions corresponding to Fig. 4, the ratio is $p_{el}/p^c \sim 0.03$ and the transition is practically jump-like. The width of this jump can be increased by the ion diffusion; one can find the corresponding expressions in [11].

If the inequality opposite to (20) holds, then all the negative ion flux should recombine in a small region of the electrode boundary layer near the plasma–layer boundary because, in the absence of diffusion, the velocity at the plasma–layer boundary is zero. The IIR arises at small values of $z = z^*$, specifically, $V(z^*) = \Gamma^*/n^c = V(z = \pi)(\nu_a \tau_{i2})$. Similarly to the previous case, the ion density in this region is higher than the equilibrium one. The growth rate of the ion density determined by the derivative of the velocity dV/dx [see also (6) and (7)] turns out to be higher than the value of $(\beta n + \nu_d)$. This gives rise to the formation of narrow peaks in the ion density. Analyzing the ion equation for $\nu_d = 0$, we obtain the expressions for the density maximum and for the peak width L. Because $dn/dx = 0$ at a maximum, we have $dV/dx \sim \beta p_{max}$ from system (4). This yields

$$p_{max} = p^c\left(\left(\frac{3}{4}\right)^{1/2}\frac{1}{\pi}\nu_a \tau_{i2}\right)^{-1/2}. \qquad (23)$$

Because the negative ion flux from the layer should recombine in the peak, we have $\beta p_{max}^2 \tilde{L} \sim \Gamma^*$ and

$$\tilde{L} = \frac{j(\nu_a \tau_{i2})}{e\omega n_e^c}. \tag{24}$$

The formation of peaks was found by Shveigert [6] in modeling a discharge. Under the conditions of Fig. 5, expression (23) gives $p_{max} = 7p^c$, whereas the calculations carried out in [6] give $p_{max} = 3.3p^c$. This difference is likely to be connected with the incorporation of both the ion diffusion and the influence of the ion oscillations on the average field.

As mentioned, the ion oscillations under the action of a microwave field can significantly affect the average field, primarily, in the region of a strong nonuniformity of the ion density, because the electron density perturbations (3) due to these oscillations can be high. The analysis of the nonuniformity region shows that the ion density oscillations arise in the vicinity of the plasma–layer boundary. The excitation of these oscillations can yield some interesting effects. For example, the numerical calculations [6] show that the harmonics with frequencies $2\omega/3$ and $\omega/3$ can arise under certain conditions. However, these oscillations are unlikely to have an important impact on the average parameters of the discharge because they are localized in a narrow region near the plasma–layer boundary.

ACKNOWLEDGMENTS

I am grateful to L.D. Tsendin for fruitful discussions and to R.M. Meier and V.A. Shveigert for providing some results of numerical calculations. This work was partially supported by the Russian Foundation for Basic Research, project no. 93-02-16873.

REFERENCES

1. Raizer, Yu.P., *Fizika Gazovogo Razryada* (Gas-Discharge Physics), Moscow: Nauka, 1987, p. 592.
2. Smirnov, A.S. and Tsendin, L.D., *IEEE Trans. Plasma Sci.*, 1991, vol. 19, p. 130.
3. Sommerer, T.J., *Modeling of Electron Kinetics in Nonuniform Fields*, Invited Lectures, XI Escampig, St. Petersburg, 1992, p. 45.
4. Gogolides, E. and Sawin, H.H., *J. Appl. Phys.*, 1992, vol. 72, p. 3971.
5. Boeuf, J.P., *Phys. Rev. A*, 1987, vol. 36, p. 2782.
6. Shveigert, V.A., *Fiz. Plazmy*, 1991, vol. 17, p. 844.
7. Oh, Y.H., Choi, N.H., and Choi, D.I., *J. Appl. Phys.*, 1990, vol. 67, p. 3264.
8. Meyyappan, M. and Govindan, T.R., *IEEE Trans. Plasma Sci.*, 1992, vol. 19, p. 122.
9. Kaganovich, I.D. and Tsendin, L.D., *Proc. XXI Int. Conf. on Phenomena in Ionized Gases*, Bochum, 1993, vol. 1, p. 63.
10. Passchier, J.P.P. and Goedheer, W.J., *J. Appl. Phys.*, 1993, vol. 73, p. 1073.
11. Tsendin, L.D., *Zh. Tekh. Fiz.*, 1989, vol. 59, p. 21.
12. Kaganovich, I.D. and Tsendin, L.D., *Proc. XXI Int. Conf. on Phenomena in Ionized Gases*, St. Petersburg, 1992, p. 340.
13. Kaganovich, I.D. and Tsendin, L.D., *Fiz. Plazmy*, 1993, vol. 19, p. 1229.
14. Kaganovich, I.D. and Tsendin, L.D., *IEEE Trans. Plasma Sci.*, 1992, vol. 20, pp. 66, 86.