Plasma-resonance discharge

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Microwave discharges have been studied experimentally and theoretically at pressures at which the rate of electron-molecule collisions is low in comparison with the field frequency, and the plasma density is above the critical value. It has been found that under these conditions the sharp enhancement of the field in the plasma-resonance region governs the discharge dynamics. A simplified theory for the dynamics of this discharge is derived. Nonlinear electrostrictive effects in the plasma-resonance layer are taken into account. The theoretical predictions agree qualitatively with experimental results.

The conditions for discharges to occur at low pressures and the mechanisms for their propagation have recently attracted considerable experimental and theoretical interest. At low pressures the rate (ν) of electron-molecule collisions is low in comparison with the frequency (ω) of the incident radiation. The propagation mechanisms have been studied quite thoroughly in cases in which the density (n) of the plasma which is produced is low in comparison with the critical density $n_{cr} = m\omega^2/4\pi e^2$ (Refs. 1 and 2, for example). In many problems of practical importance, on the other hand, the density of the plasma which is produced exceeds the critical value. This is the situation when radiation interests with solid targets in laser fusion (see, for example, the review by Afanas'ev *et al.*[3]), in the "detection" of intense microwave radiation,[4] in radiation transport through vacuum microwave lines,[5] etc. In all these cases, effects resulting from the field intensification in the plasma-resonance region appear to have a dominant effect on the dynamics of the discharge, which we will call the "plasma-resonance discharge."

In this paper we report an experimental and theoretical study of the characteristics of a microwave plasma-resonance discharge over the pressure range $3 \cdot 10^{-2} \leqslant p \leqslant 2 \cdot 10^{-1}$ torr, in which collisionless collective effects in the discharge plasma are apparently unimportant.

1. EXPERIMENTAL CONDITIONS

The experimental apparatus, described in detail in Ref. 6, is shown schematically in Fig. 1. A converging beam of millimeter-wavelength electromagnetic waves propagates in a vacuum chamber and has an approximately axisymmetric Gaussian field distribution in the focal region, with characteristic dimensions of 4 and 1 cm along and across the beam, respectively. The power of the microwave source is sufficient to produce a field $E_{max} \lesssim 7$ kV/cm at the center of the focus. The microwave source operates in pulses with a pulse length $\tau = 4 \cdot 10^{-5}$ s. The experiments are carried out at an air pressure $p = 3 \cdot 10^{-2}$–$2 \cdot 10^{-1}$ torr at the center of the induced discharge. The discharge is initiated in one of two ways:either through a preionization of the gas by a dc glow discharge or through a field intensification at sharp metal points in the focal region.

In the experiments we measure several characteristics of the plasma and the radiation, as follows:

1. A waveguide (1 in Fig. 1) beyond the focal region measures the intensity of the signal transmitted by the plasma.

2. The plasma density distributions in longitudinal and transverse cross sections with respect to the beam axis are measured by a double probe, 3. The plasma density measurements are calibrated on the basis of the cutoff of a 3-cm probing signal which is transmitted through the plasma along a bifilar line with conductors separated by 1 mm. The bifilar line consists of the conductors of a double probe.[6]

3. A qualitative study of the microwave field distribution in the discharge is carried out with a short coaxial antenna 4, which is 3 cm long and 0.5 mm in diameter, with an adapter section joining to the waveguide. We found no change in the transmitted signal which could be attributed to a perturbing effect of this device on the beam field or the discharge.

4. To study the polarization characteristics of the field we use a narrow-slot antenna, which launches the fundamental mode in a rectangular waveguide.

5. The discharge propagation velocity is studied with an FÉR-7 "photochronograph" (streak camera) and a photomultiplier with a multislit collimator. The photomultiplier detects the instant at which the plasmoid passes by the slits of the collimator, at various distances from the focal point. An absorbing screen 2 beyond the beam focus is used to attenuate the microwave signal reflected from the chamber wall.

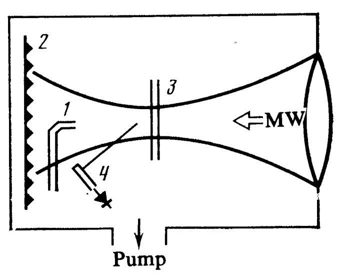

FIG. 1. Experimental arrangment.

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2. BASIC RESULTS

Regardless of the method used to initiate the breakdown, an amorphous plasma cloud initially appears at the focus, with a plasma density below the critical value. In this stage, the discharge evolves as described in Refs. 1 and 2.

When the radiation power reaches a sufficiently high level, however, the electron density reaches the critical value, and the discharge detaches from the initiating device and moves toward the microwave source. The detached discharge becomes stretched out along the field of the incident wave, acquiring an elliptical cross section with an axis ratio of about 1:4. Figure 2 shows time-integrated photographs of the discharge taken along and across the direction of the field of the incident wave. The threshold for detachment of the discharge from the initiating device is determined from the minimum incident power at which the discharge acquires the characteristic elliptical shape during the pulse. From these measurements we constructed the dependence of the field in the focal region ($E_{\rm thr}$) corresponding to the threshold value of the incident power on the air pressure in the chamber (Fig. 3). It should be noted that the detached discharge propagates in fields well below the breakdown level (the condition $E_{\rm thr} \leqslant E_{\rm br}/7$ holds under the present experimental conditions; equality is reached at $p \approx 2 \cdot 10^{-1}$ torr).

The plasmoid which is produced moves opposite the radiation source at a velocity $u = 10^5$–10^6 cm/s. This velocity varies slightly with time, despite the fact that the discharge moves into a region of weaker field. The maximum distance which the discharge traverses during a pulse is 10 cm. Under these conditions, the field falls off by a factor of five from its vacuum level along the path of the discharge. In certain cases, a secondary discharge occurs in the focal region after a delay. It moves behind the first discharge, 1–2 cm away from it. Figure 4 is a space-time diagram of the discharge.

As we mentioned earlier, the glowing region in the discharge is stretched along the field of the incident wave, with a dimension exceeding the characteristic beam width. Since the luminosity of the gas in a nonequilibrium discharge is determined by the electron temperature, which depends in turn on the applied field, we assumed that the spatial distribution of the field in the presence of the plasma was substantially different from the vacuum distribution. A study of the field by means of a coaxial antenna oriented along the incident field showed that when the plasma goes past the antenna a strong signal is induced in the antenna (Fig. 5a). This effect is observed only at the leading edge of the glowing region.[1)] The field distribution along the leading edge of the discharge is nearly uniform. Figure 5b also shows oscilloscope traces of the double-probe current. The transverse field distribution in the focal region, found with a coaxial antenna in the presence of a plasma or with a waveguide in the absence of a plasma, is shown in Fig. 6.

Study of the plasma density distribution in the discharge shows that the transverse dimension of the plasmoid exceeds the size of the wave beam and corresponds to the width of the glowing region. Comparison of the oscilloscope traces of the signals from the Langmuir and microwave probes (Fig. 5) shows that the region of intensified field is smaller by a factor of at least ten than the size of the plasmoid.[2)] Figure 7 shows the longitudinal plasma distribution found with the help of the double probe. Clearly, the width of the plasmoids is about 4 cm, and their leading and trailing edges differ in steepness. Comparison of the oscilloscope traces of the probe current with those of the signal from an rf probe shows that the field enhancement is observed only at the leading edge of the plasmoid (Fig. 5).

To study the polarization characteristics of the field in the discharge we use single-mode waveguides with a narrow

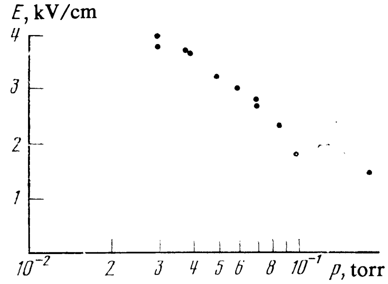

FIG. 3. Threshold field for the plasma-resonance discharge versus the pressure.

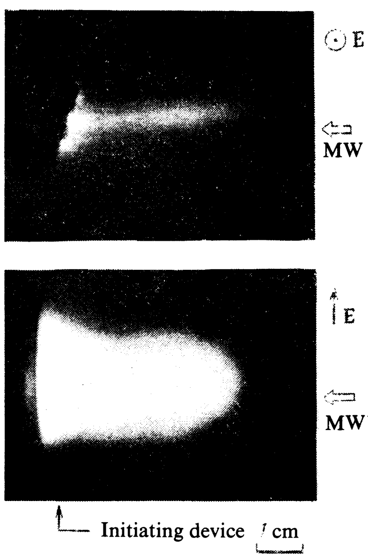

FIG. 2. Time-integrated photographs of the discharge.

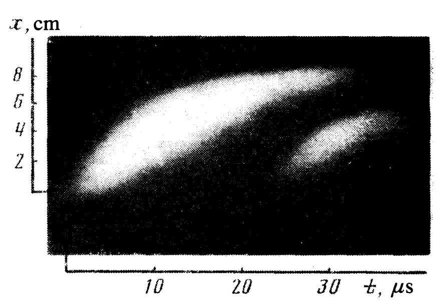

FIG. 4. "Photochronograms" (streak photos) of the discharge.

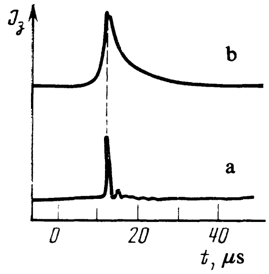

FIG. 5. Oscilloscope traces of the signals from the antenna (a) and from the double probe (b).

slot at their end. These experiments were carried out with waveguide antennas in two modifications:

1) a short-circuited waveguide of rectangular cross section, 0.2×0.9 cm, with a slot in the wide wall, oriented perpendicular to the axis of the waveguide;

2) a short-circuited waveguide of rectangular cross section, 0.2×0.9 cm, with a lot in the narrow wall, oriented perpendicular to the waveguide axis.

The waveguide antenna is oriented so that the slot faces the microwave source.

These experiments showed that if the antenna does not touch the discharge plasma no signal is detected in it, regardless of the slot orientation. In particular, the poor sensitivity of these antennas made it impossible to detect the field in vacuum even in the focal region. If the antenna was instead inserted into the discharge region, the waveguide was oriented perpendicular to the wave propagation direction (perpendicular to the vector **k**), and the slot was cut in the narrow wall, then a signal was induced in the antenna when the leading edge passed by. If the slot was instead in the wide wall, a signal was detected when the slot was perpendicular to both the field **E** and the vector **k** of the incident wave. It apparently follows from these experimental results that the field at the plasma boundary is oriented parallel to the density gradient.

Measurements at various microwave power levels and pressures in the chamber showed that the discharge propagation velocity depends only slightly on the extent to which

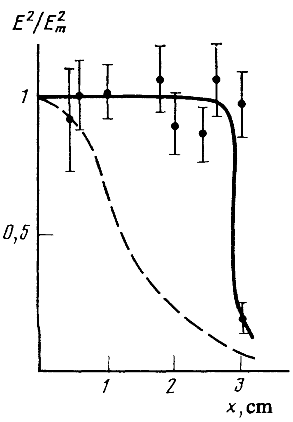

FIG. 6. Transverse field distribution in the focal region measured in the presence of the plasma (solid line) and in the absence of the plasma (dashed line).

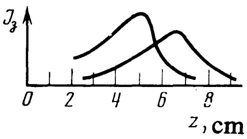

FIG. 7. Plasma profile in the longitudinal direction.

the microwave field exceeds the threshold, $E^2/E^2_{\rm thr}$, and decreases with increasing pressure.

At a pressure $p \approx 2 \cdot 10^{-1}$ torr, the discharge becomes unstable, and the plasma density distribution over the cross section of the discharge becomes extremely nonuniform.

Let us summarize the experimental results:

1) The discharge as described above occurs at pressures in the range $p = 3 \cdot 10^{-2}$–$2 \cdot 10^{-1}$ torr.

2) The threshold field for detachment of the discharge from the initiating device is well below the breakdown field; the conditions $E_{\rm thr} \lesssim E_{\rm br}/7$ holds over the entire pressure range.

3) The plasma density in the moving discharge exceeds the critical value.

4) The discharge propagation velocity reaches 10^6 cm/s and depends only weakly on the pressure and the power of the incident radiation well above the threshold.

5) The scale size of the enhanced-field region in the plasma is substantially smaller than the dimension of the plasmoid.

6) The enhanced field is oriented along the plasma density gradient.

7) The scale dimension of the plasma density distribution in the discharge along the direction of the field of the incident wave is substantially greater than (a) the transverse dimension of the unperturbed wave beam at the focus and (b) the dimension of the plasmoid in the direction perpendicular to the field.

3. SIMPLIFIED MODEL OF THE PLASMA-RESONANCE DISCHARGE

The experiments showed that the occurence and dynamics of the discharge initiated at low pressures are intimately related to the presence of a plasma-resonance region ($\varepsilon = 0$) in the plasma. In this region, the absorbed energy increases because of the increase in the normal component of the electric field ($E_n \parallel \nabla\varepsilon$). Under our conditions this effect is seen most clearly in the result that the discharge detaches from the initiating device at fields well below the breakdown value only if the plasma density reaches the critical value.

These experimental results can be described qualitatively by the following simplified model. A linearly polarized electromagnetic wave beam is incident on a plasmoid produced by the initiating device. If the plasma density is initially low ($n < n_{\rm cr}$), there is no plasma-resonance region, and the absorption is weak if the plasma slab, with scale size a, is thin in comparison with the absorption length, $\sim\lambda\omega/\nu$:

$$\frac{P_{\rm abs}}{P_{\rm inc}} \sim \frac{n}{n_{\rm cr}} \frac{a\nu}{\lambda\omega}.$$

If, instead, the condition $n > n_{\rm cr}$ holds, then the loss in the plasma-resonance region becomes dominant (we will give some specific estimates below). Since the resonant slab is thin, we can use the thin-source approximation in analyzing the electron temperature distribution. In this case the size of the heated plasma region and thus of the ionization region is determined by the scale length (L_T) of the electron thermal conductivity. The "knife-blade" (two-dimensional) structure of the discharge which is observed results from the polarization and isotropy of the field in the incident beam, since the absorption of radiation energy is localized in regions where the critical-slab surface intersects the polarization plane of the electric field. Taking the ambipolar mechanism for plasma diffusion into account, we find that the discharge will accordingly become stretched out along the incident field, and it will move in the direction opposite the incident radiation at a velocity determined by the field component (E_x) along the wave propagation direction. Under these experimental conditions (with a converging beam and with a focus size on the order of the wavelength), the field component E_x is estimated to reach values of about $0.3E_{\max}$. The scale size of the discharge in the direction perpendicular to the polarization plane of the incident field is $L_\perp \sim \max\{a, L_T\}$.

To simplify the calculations in the analysis of both the temperature profile and the density profile, we adopt a one-dimensional model along the x direction, which is the direction along which the beam and the discharge propagate. The derivatives with respect to the transverse coordinates are taken into account by means of corresponding scaling factors. We use this model to seek a solution in the form of a steady-state wave. We first determine the specific power absorbed in the plasma-resonance region, corrected for nonlinear electrostrictive effects. We then find the temperature profile; and finally, we solve the boundary-value problem for the plasma density profile. The discharge propagation velocity is determined as a function of the field and the gas pressure from the condition for the existence of a solution of this boundary-value problem. We should also point out that we ignore heating of the gas everywhere since the observed discharge propagation velocities are substantially higher than the sound velocity in the gas.

Let us estimate the quantitative characteristics of the discharge. The power (Q) of the source which heats the electrons is determined by the normal component of the electric displacement (D) in the plasma-resonance region, $D = H \sin\theta$ (θ is the angle of incidence) and by electrostrictive effects.

According to Ref. 7, the loss in the plasma-resonance region depends on the ratio of two parameters in this case. If the collision rate is low, the determing factor is the loss (Q_δ) due to the excitation of an outgoing plasma wave:

$$Q_\delta=0{,}6Q_0\frac{1}{\eta_\delta^{1/2}},\quad Q_0=\frac{1}{8}\omega l H^2\sin^2\theta,\qquad \eta_\delta=\frac{3l^2}{r_D^2}\frac{H^2\sin^2\theta}{16\pi n_{\rm cr}T},\tag{1}$$

where r_D is the electron Debye length, $l = n(dn/dx)^{-1}|_{n=n_{\rm cr}}$ is the scale of the plasma density profile in the plasma-resonance region, and the notation is otherwise standard. If instead collisions are important, the loss Q_ν is determined by

$$Q_\nu=1{,}7Q_0(\eta_\nu)^{-1/6},\qquad \eta_\nu=\left(\frac{\omega}{\nu}\right)^3\frac{H^2\sin^2\theta}{16\pi n_{\rm cr}T}.\tag{2}$$

The parameter $g = Q_\nu/Q_\delta$, which determines the relative importance of these loss mechanisms, can be written as follows according to (1) and (2):

$$g=1{,}3\frac{l}{\Lambda}\left(\frac{\varepsilon_\sim}{T}\right)^{1/6},\qquad \varepsilon_\sim=\frac{e^2H^2\sin^2\theta}{m\omega^2}.\tag{3}$$

Here $\Lambda = V_{Te}/\nu$ is the electron mean free path with respect to collisions with molecules, and V_{Te} is the electron thermal velocity. For air we would have Λ [cm] = $1/(30p)$ [torr]. Under our experimental conditions we have $l \gtrsim 1$ cm (Fig. 7), and since the field dependence is comparatively weak, according to (3), it is clear that the condition $g > 1$ can hold only at sufficiently high pressure. In particular, at the lower boundary of the pressure range studied, $p \approx 3 \cdot 10^{-2}$ torr, we have $g \gtrsim 1$. We will therefore assume that the loss is determined by collisions, and we will use expression (2).

Knowing the power of the heating source, we can easily find the electron temperature profile $T(x)$. For this purpose we use the heat-conduction equation, which we simplify by ignoring the coordinate dependence of the coefficients:

$$\frac{\partial T}{\partial t}=\varkappa_e\frac{\partial^2T}{\partial x^2}-(\delta\nu+\nu_\perp)T+\frac{Q}{l}\delta(n(x)-n_{\rm cr}).\tag{4}$$

Here $\alpha_e = V_{Te}^2/\nu$ is the electron thermal conductivity; $\delta \approx 10^{-1}$–10^{-2} is the fraction of the energy which is lost in the collision of electrons with molecules; and $\nu_\perp \approx \varkappa_e/L_\perp^2$ is the effective transverse loss rate caused by the finite transverse dimensions of the discharge. We will assume that the transverse loss is described by the term with $\delta\nu$. This approach is formally legitimate, since the condition $L_\perp^2 \sim L_T^2 = \varkappa_e/\delta\nu = \Lambda^2/\delta$ holds under these experimental conditions. We seek a solution of Eq. (4) in the form of a steady-state wave $T(x - ut)$. In particular, for $u = 0$ we have

$$T(x)=T_m\exp\left\{-\frac{|x|}{L_T}\right\},\qquad L_T^2=\frac{\Lambda}{\delta},\qquad T_m=\frac{Q_\nu}{2n_{\rm cr}L_T\delta\nu}.\tag{5}$$

Simple estimates show that when

$$u\ll L_T\delta\nu\approx V_{Te}\delta^{1/2},\tag{6}$$

which holds under our experimental conditions, the effect of drift on the temperature distribution can be ignored. For the discussion below it is convenient to express T_m in terms of the parameters of the problem: H, θ, ω, and Λ. Using $\varkappa_e = V_{Te}^2/\nu$, $\nu = V_{Te}/\Lambda$, and $V_{Te}^2 = T/m$, and recalling that the condition $T = T_m$ holds in the plasma resonance region, we find from (2) and (3)

$$T_m = 3.1\,(l/\delta^{1/2}\Lambda)^{2/3}\varepsilon_\sim. \quad (7)$$

Let us examine the particle number balance. We use the diffusion equation

$$\frac{\partial n}{\partial t} = D\frac{\partial^2 n}{\partial x^2} + \nu_i n - \frac{D}{L_\perp^2}n - \alpha n^2, \qquad \nu_i = 5\cdot10^4 p\left(\frac{T}{T_0}\right)^{8/3}. \quad (8)$$

Here $D = V_s^2/_{im}$ is the ambipolar diffusion coefficient (under our conditions, the diffusion becomes ambipolar at a density as low as $n \geqslant 10^6$ cm^{-3}); ν_i is the ionization rate ($T_0 \sim$ 1–2 eV; the gas pressure p is expressed in torr)[8]; V_s is the ion acoustic velocity; ν_{im} is the rate of ion-molecule collisions; α is the recombination coefficient; and the coefficient D/L^2 describes the transverse loss of electrons through ambipolar diffusion. Estimates show that the recombination is inconsequential under our conditions ($\alpha < 10^{-7}$ cm^3/s).

Seeking a solution of (7) in the form of a steady-state wave $n(x - ut)$, we find a boundary-value problem for $n(\xi)$, where $\xi = x - ut$, with the boundary conditions $n(x) \to 0$ and $x \to \infty$. Since we have $n = n_{cr}$ at $\xi = 0$, we have, by definition,

$$l = 2D/u. \quad (9)$$

We introduce a new function and some new dimensionless variables:

$$n = n_{dl}\ \exp\ (u\xi/2D),\ \xi_{dl} = \xi/L_T. \quad (10)$$

For air we can then write (dropping the subscript "dl")

$$\frac{d^2n}{d\xi^2} + \bar{\nu} n \exp\left(-\frac{8}{3}|\xi|\right) = \gamma n, \qquad \gamma = \left(\frac{\Lambda}{\delta^{1/2}L_T}\right)^2 + \left(\frac{\Lambda}{\delta^{1/2}l}\right)^2,$$

$$\bar{\nu} = \frac{5\cdot10^{-3}}{\delta}\left(\frac{T}{T_0}\right)^{5/3}. \quad (11)$$

Here $n(\xi)$ is the symmetric (fundamental) mode of Eq. (11). Since

$$4\,\mathrm{ch}^{-2}\left(\frac{4}{3}\xi\right) \leqslant \exp\left(-\frac{8}{3}|\xi|\right) < \cosh^{-2}\left(\frac{4}{3}|\xi|\right), \quad (12)$$

and since Eq. (11) can be solved exactly with the "potential" $A\cosh^{-2}\xi$, we find a uniform estimate for γ. This estimate takes a particularly simple form under the condition $\gamma \gg (\Lambda/\delta^{1/2}L_T)^2$ (above the threshold):

$$\gamma \sim \bar{\nu}. \quad (13)$$

Using (71), (8), and (11), we then find the expression which we are seeking for the velocity at a microwave power well above the threshold: With $\gamma = (\Lambda/\delta^{1/2}l)^2 \sim \bar{\nu}$,

$$u \approx \frac{3\cdot10^6}{\delta^{0,6}}\left(\frac{\varepsilon_\sim}{T_0}\right)^{0,8} \approx 5\cdot10^7\left(\frac{\varepsilon_\sim}{T_0}\right)^{0,8}\Bigg|_{\delta=10^{-2}}. \quad (14)$$

The temperature in the discharge, like the velocity, is independent of the pressure:

$$T_m[\mathrm{eV}] \approx \frac{10}{\delta^{1/3}}(\varepsilon_\sim T_0)^{1/2}, \qquad \frac{l}{\Lambda} \approx 2\delta^{0,3}\left(\frac{T_\nu}{\varepsilon_\sim}\right)^{0,4}. \quad (15)$$

We see from (15) that, as expected, l decreases with increasing field.

As we have already mentioned, relations (14) and (15) hold only above the threshold, when the discharge plasma is well above the critical density. The formal reason is that in the limit $u \to 0$ we cannot use Eq. (9) for l in estimating the absorbed power Q [according to (9), we have $l \to 0$ and $Q \to \infty$ in the limit $u \to 0$]. Actually, of course, the absorption is finite. At the threshold, with a plasma slightly above the critical density,

$$\varepsilon = (n_{cr} - n_{max})/n_{cr}, \qquad |\varepsilon| < \theta \ll 1, \quad (16)$$

field screening effects become important (the normal component of the electric displacement decreases).[9] Estimates show that we have $l_{max} \approx L_T$ under these conditions. Since we have $\gamma = 1$ at the threshold, we also have $\nu \approx 1$, and, according to (7) and (11), we have

$$\varepsilon_\sim|_{n=n_{cr}} \approx 0{,}3\delta^{4/3}T_0. \quad (17)$$

It is a simple matter to express the electron oscillation energy in the plasma-resonance region in terms of the field of the incident wave if we note that at the threshold the condition $l \gg \lambda$ holds, where λ is the length of the electromagnetic wave. Here we have

$$\varepsilon_\sim|_{n=n_{cr}} = \varepsilon_{\sim\,thr}\left(\frac{L_T}{\lambda}\right)^{1/3} \approx \varepsilon_{thr}\left(\frac{\Lambda}{\lambda\delta^{1/2}}\right)^{1/3}. \quad (18)$$

Comparing (17) and (18), we find

$$\varepsilon_{\sim\,thr} \sim 0{,}3\delta^{5/3}\left(\frac{\Lambda}{\lambda}\right)^{1/3}T_0 \approx 0{,}1\delta^{5/3}(p\cdot\lambda\ [\mathrm{torr\cdot cm}])^{-1/3}. \quad (19)$$

It can be seen from this relation that the threshold field decreases with increasing pressure, in qualitative agreement with the experimental results (Fig. 3).

4. CONCLUSION

It has thus been shown experimentally and theoretically that the threshold characteristics and dynamics of the low-pressure discharge are governed by a plasma resonance. The qualitative agreement of the simple model outlined above with the experimental results [the threshold dependence in Fig. 3 and (19); the weak dependence of the discharge velocity on the pressure far from the threshold; the theoretically predicted and experimentally observed steepness of the leading edge (Fig. 7)] is not simply a fortuitous result, since the discharge temperature is low, according to (15), and our assumption of a collisional mechanism for absorption is clearly justified at pressures which are not too low, $3 \cdot 10^{-2} < p < 2 \cdot 10^{-1}$ torr. It should be noted that the discharge velocity, determined for a microwave power above the threshold by (14) in this theory, also agrees well with the experimental data when we note that the angle of incidence (an adjustable parameter in the one-dimensional model) is on the order of the convergence angle of the incident beam. In a realistic model, the discharge profile would not be planar, and the effective angle of incidence would have to be determined in a self-consistent multidimensional model.

One of the assumptions underlying this model is that the loss in the plasma-resonance region is determined by collisions. Under this condition, the parameter g in (3) would have to be greater than unity. From (5) we have $g \approx 1.1(\delta l^3/\Lambda^3)^{1/8}$. We thus see that the condition $g > 1$ holds if $l > \Lambda\delta^{1/3}$. Under our experimental conditions this condition does hold, although not by a wide margin.

The arguments which we have presented here do not, of course, describe discharges at lower pressures, where the spatial dispersion of the waves excited in the plasma-resonance region becomes a governing factor.[7] In that case, in contrast with our own situation, fast electrons can apparently play an important role in the discharge energy balance. Until this question is resolved, we do not believe it is appropriate to write the relations corresponding to a collisionless loss Q_δ in (2). It should be noted that the mechanism proposed here could apparently be important in gas breakdown in the IR range, e.g., in a CO_2 laser beam, where plasma densities on the order of the critical density are also observed.[10]

We wish to thank V. B. Gil'denburg, V. E. Semenov, and A. V. Kim for useful discussions.

[1)]The field enhancement cannot be attributed to the occurrence of resonant conditions at the probe, since the probe was covered by a relatively thick insulating film.

[2)]More accurate measurements of the field would be possible only if the relation between the width of the slot in the waveguide and the width of the resonance region were known.

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