EXPANSION OF THE REGION OF MICROARC DISCHARGES ON THE SURFACE OF AN ELECTRODE ACCOMPANYING FORMATION OF TURBULENCE IN THE BOUNDARY LAYER IN A SUPERSONIC PLASMA FLOW

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It is well known that nonstationary, rapidly mixing microarcs, supported on a dividing cell of the cathode spot, whose number is proportional to the total discharge current, form on the surface of electrodes coated with an oxide film under conditions of vacuum arcs and low pressures of the surrounding gas [1]. In addition, the cells of the spot repel one another, which is caused by the characteristic magnetic fields of the cells and the spot [1, 2], and propagate over the entire surface of the electrode. When the pressure of the surrounding gas is raised the mutual repulsion of the cells becomes weaker and then vanishes, and the microarcs are attracted to one another by ampere forces acting between currents flowing in parallel directions, which gives a compact arrangement of the cells of the spot and gives rise to the formation of virtually stationary current attachment on the surface of the electrode.

Dispersal of the dividing microarcs on the surface of a mesh electrode with a circular cross section 1 cm in diameter, placed on one wall of the dielectric section of a shock tube, was observed under conditions of supersonic flow of an argon plasma with strong shock waves in a shock tube with flow velocities of $3\cdot10^5$ cm$\cdot$sec^{-1} at a temperature of ~8000°K and under a pressure of ~0.15 MPa. The other electrode (the anode) was located on the opposite wall of the section. The electrodes were connected in the circuit with a precharged capacitor bank, whose discharge through the interelectrode gap was initiated by the leading edge of the ionizing shock wave at the location of the electrodes.

Figure 1 shows a typical photoscan of the luminescence from the cells of the cathode spot on the surface of the wall electrode and the corresponding phased oscillogram of the discharge current. According to the photoscan, the dispersal of the cells starts deep in the region of the shock-compressed flow of argon plasma. The start of the dispersal (branching of the tracks) of the cells, as established with the help of a film temperature sensor, corresponds to the moment (~110 μsec) at which turbulence appears in the nonstationary laminar boundary layer behind the shock wave at the location of the electrode on the wall of the measuring section.

The result obtained, however, is not consistent with the well-known fact that for gas pressures of the order of one atmosphere there is no mutual repulsion of the cells of the cathode spot [2]. In this case, the dispersal of the microarcs is probably caused not by the mutual repulsion of the cells in the characteristic magnetic field of the cathode spot, but rather it is determined by the action of chaotic turbulent pulsations of the gas velocity in the boundary layer, which in spite of the attractive ampere forces acting between the microarcs, smear the compact group of cathode spots, forming under the conditions of a laminar boundary layer preceding turbulence; which is what explains the character of the observed photoscan after 110 μsec.

It is conjectured that the dispersal of the microarcs under the action of the turbulent pulsations can be described by diffusion processes analogously to [3] under the conditions of mutual attraction of microarcs by ampere forces. To simplify the analysis it is assumed that the current in a separate microarc remains constant and the number of microarcs remains constant during the process of dispersal, so that the microarcs are regarded as stationary formations.

We shall write the equation of diffusion of microarcs, having the form of the Fokker–Planck equation, in the cylindrical centrosymmetric case as

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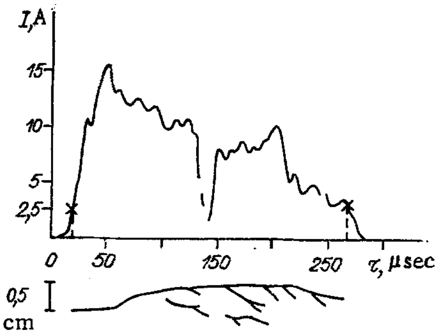

Fig. 1

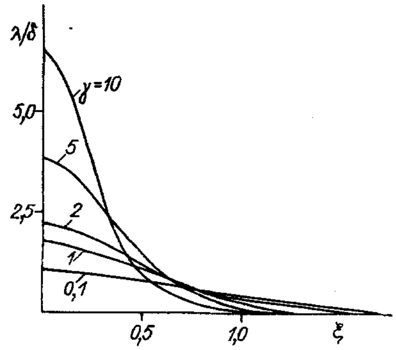

Fig. 2

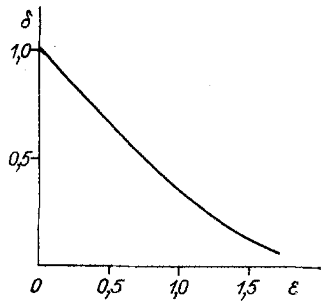

Fig. 3

$$\frac{\partial n}{\partial t} = \frac{1}{r}\frac{\partial}{\partial r}\left[r\left(D\frac{\partial n}{\partial r} - nbf\right)\right] \tag{1}$$

where n is the density of microarcs, D and b are the turbulent analogs of the coefficient of diffusion and mobility, f is the ampere force, r is the radius, and t is the time. The functions D and b, as usual, are related by the Einstein relation

$$D/b = T \tag{2}$$

(T is the turbulent analog of the temperature). The ampere force acting on the straight electric current I of length ℓ in the magnetic field H has the form f = −IHℓ/c (the minus sign indicates attraction). From the Maxwell's equation $\operatorname{rot}\mathbf{H} = \frac{4\pi}{c}\mathbf{j}$ we find

$$H = \frac{2}{cr}\int j\,dS.$$

Here j is the current density; the integration extends over a circle with radius r. If I is the current in one cell of the spot, then assuming that I = const, which agrees with the experimental data of [1] we obtain

$$f = -4\pi\left(\frac{I}{c}\right)^2 l\frac{1}{r}\int\limits_0^r n\rho\,d\rho. \tag{3}$$

We write Eq. (1), substituting (2) and (3), as

$$\frac{\partial n}{\partial t} = \frac{1}{r}\frac{\partial}{\partial r}\left[rD\left(\frac{\partial n}{\partial r} + \varkappa\frac{n}{r}\int\limits_0^r n\rho\,d\rho\right)\right], \tag{4}$$

where for T = const $\varkappa = \frac{4\pi\,(I/c)^2 l}{T} = \mathrm{const} > 0.$

In what follows, to simplify the analysis, aside from the condition $\varkappa$ = const it is assumed that D = const.

Equation (4) must be solved for n(r, t) under the conditions

$$n(r,0) = n_0(r),\ \partial n/\partial r(0,t) = 0,\ n(\infty,t) = 0;$$
$$\int\limits_0^\infty n2\pi r\,dr = \mathrm{const} = N \tag{5}$$

(N is the total number of microarcs).

We seek the self-similar solution of (4) in the form

$$n = \varphi(t)\nu(\eta),\ r = \psi(t)\eta \tag{6}$$

(ν and η are the dimensionless values of the function and the argument). Substitution of (6) into (4) gives

$$-\frac{\psi\psi'}{D}\left(\eta^2\frac{d\nu}{d\eta} - \frac{\varphi'/\varphi}{\psi'/\psi}\eta\nu\right) = \frac{d}{d\eta}\left[\eta\left(\frac{d\nu}{d\eta} + \varkappa\varphi\psi^2\frac{\nu}{\eta}\int\limits_0^\eta \nu\xi\,d\xi\right)\right] \tag{7}$$

(the prime indicates differentiation with respect to t). The condition (5) implies the relation

$$-\frac{\varphi'/\varphi}{\psi'/\psi} = \text{const} = 2, \tag{8}$$

while the first integral of (7) is

$$\eta\left(\frac{d\nu}{d\eta} + \varkappa\varphi\psi^2\frac{\nu}{\eta}\int_0^\eta \nu\xi d\xi\right) + \frac{\psi\psi'}{D}\eta^2\nu = C$$

(C is the integration constant).

For $n(\infty, t) = 0$ $C = 0$, whence we obtain the equation

$$\frac{\partial\nu}{\partial\eta} + \left(\frac{\psi\psi'}{D}\eta + \varkappa\varphi\psi^2\frac{1}{\eta}\int_0^\eta \nu\xi d\xi\right)\nu = 0. \tag{9}$$

The conditions for self-similarity follows from (8) and (9): $\varphi\psi^2 = \text{const} = \alpha > 0$, $\psi\psi'/D = \text{const} = \beta > 0$, the latter condition gives the usual result $\psi = \sqrt{2\beta Dt}$ if the integration constant equals zero.

Introducing $\zeta = \frac{\beta\eta^2}{2} = \left(\frac{r}{2\sqrt{Dt}}\right)^2$ and the functions $\lambda = \frac{\nu}{\nu(0)}$, $\mu = \int_0^\zeta \lambda d\xi$, we obtain

$$\frac{d}{d\zeta}\left(\frac{\zeta d\lambda}{\lambda d\zeta}\right) + (1 + \gamma\lambda) = 0; \tag{10}$$

$$\frac{d^2\mu}{d\zeta^2} + \left(1 + \gamma\frac{\mu}{\zeta}\right)\frac{d\mu}{d\zeta} = 0, \tag{11}$$

where $\gamma = \varkappa\frac{\alpha}{2\beta}\nu(0) > 0$. The integral representation

$$\lambda = \exp\left[-\zeta - \gamma\int_0^\zeta \lambda(\xi)\ln\frac{\zeta}{\xi}d\xi\right] \tag{12}$$

follows from (11). By definition $\lambda(0) = 1$, $\mu(0) = 0$, and from (12), since $\lambda \geq 0$, we have $\lambda(\infty) = 0$. The undetermined constant γ is determined from the condition (5), which has the form

$$\mu(\infty) = \int_0^\infty \lambda d\zeta \equiv \delta = \frac{\varepsilon}{\gamma} \quad (\varepsilon = \varkappa N/4\pi). \tag{13}$$

The fact that the solutions of Eqs. (10) and (11) do indeed exist can be seen from the iteration procedure according to (12):

$$0 < \lambda < 1,\ 0 < \exp[-(1+\gamma)\zeta] < \lambda < \exp(-\zeta) < 1,$$

$$0 < \exp[-(1+\gamma)\zeta] < \exp\left[-\zeta - \gamma\int_0^\zeta \exp(-\xi)\ln\frac{\zeta}{\xi}d\xi\right] <$$

$$< \lambda < \exp\left\{-\zeta - \gamma\int_0^\zeta \exp[-(1+\gamma)\xi]\ln\frac{\zeta}{\xi}d\xi\right\} < \exp(-\zeta) < 1$$

etc. The solutions of these equations in quadratures, by virtue of their nonlinearity and apparently the absence of group properties of the transformations (except the substitution $\lambda \to \gamma\lambda$, $\mu \to \gamma\mu$, eliminating γ), present a problem. We note that the solution of (10) can be constructed in the form of a power series

$$\lambda = \sum_{k=0}^\infty a_k\zeta^k/k! \quad (a_0 = 1),$$
$$a_{k+1} = -a_k - \gamma\sum_{l=0}^k \frac{1}{l+1}c_k^l a_{k-l}a_l \quad (k \geq 0). \tag{14}$$

Numerical calculations of λ in the function ζ were performed by solving Eq. (10) by the Runge–Kutta method. Calculations of λ using (14) gave identical results. Figure 2 shows the dependence of the normalized density λ/δ on the dimensionless coordinate $\xi = r/(2\sqrt{Dt})$ for the parameter γ; Fig. 3 shows the dependence of the constant δ on ε. For the absent ampere force ($\varepsilon = 0$), evidently, $\lambda = \exp(-\xi^2)$ and $\delta = 1$. Taking into account ampere forces causes the source function λ/δ to peak at the origin of coordinates, but for any intensity of the ampere interaction the diffusive dispersal of the ensemble of microarcs is not blocked by this interaction, and a limiting stationary distribution density of microarcs does not exist. This is a consequence of the fact that the ampere force, passing through a maximum, approaches zero at infinity.

Thus it has been demonstrated that the diffusion mechanism for dispersal of a compact distribution of microarcs by turbulent pulsations on the surface of an electrode is in principle possible.

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DISCHARGE ACCOMPANYING LEAKAGE OF MAGNETIC FLUX FROM PLASMA INTO AN INSULATOR

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In many problems, such as the confinement of plasma with a magnetic field by walls, compression of a magnetized plasma with liners, etc., losses of magnetic flux and plasma owing to diffusion of the field and heat conduction to the wall must be taken into account. The role of the discharge arising in the plasma as magnetic flux leaks out of it must be especially significant for a hydrogen plasma, whose conductivity, owing to the weak effect of radiative processes, can be large compared with the conductivity of the plasma in a magnetically compressed discharge [1] arising on the surface of the wall. In this case, if the plasma density is too high, the resistance of the discharge will be determined by the discharge along the hydrogen plasma.

We shall study the development of this discharge in the case of a hydrogen plasma with a magnetic field bounded by a rigid nonconducting insulating wall. This problem was solved qualitatively in [2, 3], and as a result the effective diffusion coefficient for a plasma with $\beta \ll 1$ ($\beta = 16\pi N_0T_0/H_0^2$ is the ratio of the thermal pressure of the plasma to the magnetic pressure, and N_0, T_0, H_0 are the density, temperature, and magnetic field in the plasma far from the discharge zone) $D \sim cH_0/4\pi eN_0$, and for $\beta \gg 1$, $D \sim cT_0/10eH_0$.

In this paper the structure of the current layer near the wall is studied quantitatively and the boundary condition with whose help the effect of this discharge on the motion of the plasma in the entire volume can be described is formulated.

Let all quantities depend on the coordinate X and the time t, let the magnetic H and electric E fields be perpendicular to one another and the X axis, and let the characteristic times be long compared with the gas-dynamic times, so that there is enough time for the total pressure in the system to be equalized:

$$2NT + H^2/8\pi = p_0 \equiv 2N_0T_0 + H_0^2/8\pi. \tag{1}$$

The plasma density in the main volume is assumed to be low compared with the density of the discharge zone near the wall. In this case, as shown in [2], the problem is quasistationary, i.e., the time derivatives in the magnetic and electric field equations and the equation of

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