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.

We thank P. P. Lazarev for assistance in performing the calculations.

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DISCHARGE ACCOMPANYING LEAKAGE OF MAGNETIC FLUX FROM PLASMA INTO AN INSULATOR

S. F. Garanin

UDC 533.95

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

Moscow. Translated from Zhurnal Prikladnoi Mekhaniki i Tekhnicheskoi Fiziki, No. 6, pp. 13-16, November-December, 1987. Original article submitted August 26, 1986.

 0021-8944/87/2806-0816$12.50

thermal balance of the plasma in the discharge zone can be neglected, while the electric field and the energy flux can be assumed to be constant. Then these equations have the form

$$E=-\frac{c}{4\pi\sigma}\frac{\partial H}{\partial X}-\frac{b_\Lambda}{e}\frac{\partial T}{\partial X},\quad Q=-\chi\frac{\partial T}{\partial X}-\frac{cT}{4\pi e}b_\Lambda\frac{\partial H}{\partial X}+\frac{c}{4\pi}EH, \tag{2}$$

where σ, χ, and b_Λ are the transverse conductivity, thermal conductivity, and thermoelectric coefficient; Q is the energy flux. The mass of the plasma accumulated in the layer near the wall is given by

$$a=\int_0^{X_0} N dX \tag{3}$$

(X_0 is the boundary of the discharge zone). The plane X = 0 is assumed to be the wall, and the plasma occupies the region X > 0. Then the boundary conditions to Eqs. (1) and (2) are as follows:

$$T(0)=0,\ H(0)=H_1,\ N(X_0)=0 \tag{4}$$

(H_1 is the magnetic field in the insulator). The energy flux flowing into the discharge zone from the plasma $Q=\frac{c}{4\pi}EH_0+5N_0T_0v$ (v is the velocity of the inflowing plasma). Because the magnetic field is frozen into the plasma, far from the insulator

$$v=cE/H_0 \tag{5}$$

and hence $Q=\frac{c}{4\pi}EH_0\left(1+\frac{5}{4}\beta\right)$.

On the boundary of the plasma with the insulator the plasma is not magnetized because of (4). In this region, as the temperature increases the coefficients χ, b_Λ increase away from the wall; the characteristic dimensions X, corresponding to the temperature T, increases according to (2) and accumulation of the discharge plasma occurs. When the plasma is magnetized the coefficients χ, b_Λ decrease, and therefore the characteristic dimensions X as well as the density decrease. For this reason the region in which the degree of magnetization of the electrons $\omega_e\tau_e \sim 1$ will make the main contribution to the mass of the plasma accumulated in the discharge. For the unit of measurement of the pressure in the problem it is natural to choose p_0, while the unit of measurement of the temperature [T] and the density [N] are chosen so that for T = [T], N = [N] $p = p_0$ and $\omega_e\tau_e \sim 1$ (see [1]). From these conditions we obtain

$$[T]=(\sqrt{p_0}c\lambda e^3\sqrt{m})^{2/5},\quad [N]=p_0^{4/5}/(c\lambda e^3\sqrt{m})^{2/5} \tag{6}$$

(m is the electron mass and λ is the Coulomb logarithm).

It is convenient to choose the unit of measurement of the coordinate X starting from Eq. (2), substituting into them for the temperature and density [T] and [N] from (6). Then, introducing the dimensionless coordinate $x=-\frac{E}{e^{0,2}m^{0,2}c^{0,4}\lambda^{0,4}p_0^{0,2}}X$ (the electric field E is negative) and the dimensionless variables

$$\theta(x)=T/[T],\ n(x)=N/[N],\ h(x)=H/\sqrt{8\pi p_0},\ \xi=-eEa/p_0, \tag{7}$$

the system of equations (1)-(3) can be rewritten in the form

$$\begin{gathered}2n\theta+h^2=1,\quad \frac{4}{3}\frac{\alpha}{\theta^{3/2}}h'+b\theta'=1,\\ b\theta h'+\frac{3}{4}\left(\gamma+\frac{\gamma_i}{\sqrt{912A}}\right)\theta^{5/2}\theta'=\frac{1+1{,}25\beta}{\sqrt{1+\beta}}-h,\\ \xi=\int_0^{x_0} n dx,\end{gathered} \tag{8}$$

where α, b, γ, γ_i depend on the degree of magnetization $y=\omega_e\tau_e=\frac{3}{2}\frac{h}{n}\theta^{3/2}$ and are determined by the approximate formulas [4]

$$\alpha=1-\frac{6{,}42y^2+1{,}86}{\Delta},\quad b=\frac{y(1{,}5y^2+3{,}05)}{\Delta},\quad \gamma=\frac{4{,}66y^2+12{,}1}{\Delta},$$
$$\gamma_i=\frac{2y_i^2+2{,}64}{y_i^4+2{,}7y_i^2+0{,}677},\quad \Delta=3{,}77+14{,}8y^2+y^4,\quad y_i=y/\sqrt{912A}$$

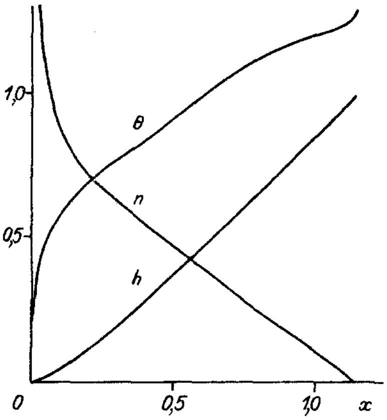

Fig. 1

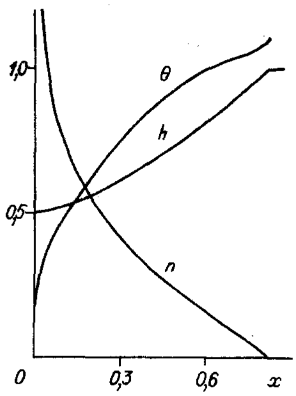

Fig. 2

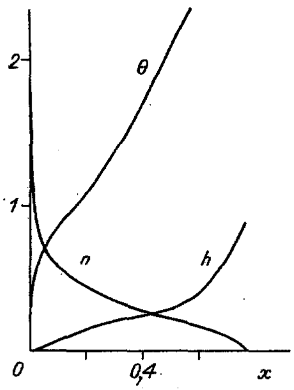

Fig. 3

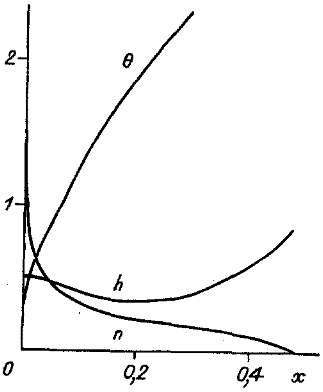

Fig. 4

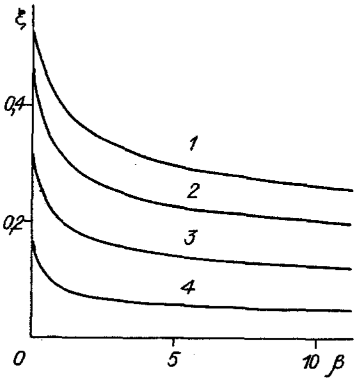

Fig. 5

(A is the atomic weight). Here, unlike [1], the term γ_i, representing the ionic thermal conductivity, which is a small correction, is taken into account. Taking into account the ionic thermal conductivity, however, changes the behavior of the solution in the region of strong magnetization $x \simeq x_0$. The boundary conditions (4) in the new notation are written as

$$h(0) = h_1,\ \theta(0) = 0,\ n(x_0) = 0,\ h_1 = H_1/\sqrt{8\pi p_0}. \tag{9}$$

Let us evaluate the orders of magnitude of the quantities characterizing the discharge zone for large β. In this case, it follows from Eqs. (8) that the thermoelectric transport of magnetic flux makes the main contribution to the electric field, while heat conduction makes the main contribution to the energy flux. Then, taking into account the fact that in the discharge zone $y \sim 1$, we obtain $h \sim 1/\sqrt{\beta}$, $T \sim \beta^{0.2}$, $n \sim \beta^{-0.2}$, $x \sim \beta^{0.2}$, $\xi \sim 1$. Thus the parameter ξ, characterizing the accumulated mass, is virtually independent of β and changes only for $\beta \sim 1$.

The results of numerical calculations of the system (8) with the boundary conditions (9) for $h_1 = 0;\ 0.5;\ 0;\ 0.5$, $\beta = 0;\ 0;\ 10;\ 10$, $A = 2$ are presented in Figs. 1-4, while the dependence of ξ on β and h_1 is presented in Fig. 5 ($h_1 = 0;\ 0.25;\ 0.5;\ 0.75$, lines 1-4).

The calculations showed that the thermoelectric phenomena which play the main role in the transport of magnetic flux for $\beta \gg 1$ and lead to leakage of magnetic flux even into an insulator with a stronger magnetic field than in the plasma (Fig. 4) are numerically not very significant for $\beta \sim 1$. Thus if in Eqs. (8) b is set equal to zero, then for $\beta = 0$ and $h_1 = 0$ the value of ξ decreases by only 13%. An appreciable decrease in ξ (approximately by 30%) accompanying cutoff of the thermoelectric fluxes occurs only for $\beta \simeq 10$. The ionic thermal conductivity, which is a correction $\sim 1/\sqrt{912A}$, makes a very small contribution to the accumulated mass ξ, switching it off for $\beta = 0$, $h_1 = 0$ decreases ξ by 2%.

Let us examine the dynamics of plasma deposition. The rate of accumulation of mass $da/dt = N_0 v$, according to (5) and (7), is determined from the differential equation $a\frac{da}{dt} = \xi(\beta, h_1)\frac{c}{e}\frac{p_0 N_0}{H_0}$, which for the entire plasma volume can be regarded as a boundary condition, describing the loss of plasma and magnetic flux. For $H_0(t)$ = const, $N_0(t)$ = const, $p_0(t)$ = const, $H_1(t)$ = const, $\beta(t)$ = const the deposition of the plasma occurs according to the diffusion law

$$a = \sqrt{2\xi \frac{c}{e}\frac{p_0 N_0}{H_0} t}, \qquad\qquad (10)$$
$$E = \sqrt{\frac{\xi}{2ec}\frac{p_0 H_0}{N_0 t}}, \quad [X] = \frac{e^{0,7} m^{0,2} c^{0,9} \lambda^{0,4}}{p_0^{0,3}} \sqrt{\frac{2N_0 t}{\xi H_0}}.$$

The effective coefficient of diffusion in this case $D \sim 2\xi c p_0/eN_0H_0$, which for $\beta \ll 1$, when $\xi \approx 0.5$ (Fig. 5), gives $D \sim cH_0/8\pi e N_0$ and approximately corresponds to the estimate of [2], while for $\beta \gg 1$, when $\xi \approx 0.25$ (Fig. 5), $D \sim cT_0/eH_0$ and exceeds the estimate of [3] and Bohm's thermal conductivity by approximately an order of magnitude.

We shall examine the conditions of applicability of the solution presented for the problem of a discharge near a wall. We confine our attention to megagauss magnetic fields and $\beta \sim 1$. The plasma density N_0 must be much less than the plasma density in the discharge zone, i.e., in accordance with [1] $N_0(\mathrm{cm}^{-3}) < 3\cdot 10^{20} H^{1.6}$ (MG). In this case the problem can be regarded as quasistationary and Eqs. (2) and the boundary condition $N(X_0) = 0$ can be employed. For $N_0(\mathrm{cm}^{-3}) \gtrsim 3\cdot 10^{20} H^{1.6}$ (MG) the plasma density in the discharge zone is of the order of N_0 and the calculations performed are not applicable. The condition for being able to neglect the radiation losses in the discharge zone gives

$$N_0(\mathrm{cm}^{-3}) t(\mu\mathrm{sec}) < 0.8\cdot 10^{20} H^{0.2}\,(\mathrm{MG}). \qquad (11)$$

In addition, the role of the magnetically compressed discharge in the insulator was neglected, i.e., it was assumed that the magnetic flux flowing from the discharge along the plasma outside the insulator is too weak for a magnetically compressed discharge to form in the insulator. This means that the electric fields (10) must be less than the fields found in [1].* For an insulator consisting of organic glass $H_8C_5O_2$ we obtain from this condition for $t \times (\mu\mathrm{sec}) < 0.03/H^{1.12}$ (MG)

$$N_0(\mathrm{cm}^{-3}) t(\mu\mathrm{sec}) > 9\cdot 10^{15} H^{0.6}\,(\mathrm{MG}); \qquad (12)$$

and for $t(\mu\mathrm{sec}) > 0.03/H^{1.12}$(MG)

$$N_0(\mathrm{cm}^{-3}) > 2\cdot 10^{17} H^{1.72}\,(\mathrm{MG}). \qquad (13)$$

When the condition (11) does not hold, i.e., for sufficiently long times, the discharge enters a stationary stage, described in [1]. When the conditions (12) and (13) no longer hold, i.e., for sufficiently small densities of the hydrogen plasma, the magnetically compressed discharge studied in [1] appears on the surface of the insulator.

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