SOVIET PHYSICS JETP VOLUME 29, NUMBER 5 NOVEMBER, 1969

PLASMA FOCUS AND THE NEUTRON EMISSION MECHANISM IN A Z-PINCH

V. P. DYACHENKO and V. S. IMSHENNIK

Institute of Applied Mathematics, USSR Academy of Sciences

Submitted December 20, 1968

Zh. Eksp. Teor. Fiz. 56, 1766–1777 (May, 1969)

The two-dimensional nonstationary magnetohydrodynamic problem employed for the theoretical interpretation of experiments with a noncylindrical z-pinch is considered. For the sake of simplicity the only dissipative process taken into account in the equations is viscosity which yields a rough description of the thickness of the shock front. The main results of a numerical solution of the problem by the large particle method are, on the one hand, elucidation of the process of formation of a plasma focus and, on the other hand, determination of the parameters of the axial plasma jet. The calculations show that for complete determination of the characteristics of the plasma focus it is necessary to include in the two-dimensional problem effects of finiteness of the plasma electric conductivity and of the sharp increase of inductive resistance at the moment of formation of the focus. Neutron escape from the jet region is estimated on the basis of the quantitative parameters of the axial plasma jet; this value is comparable to that of the general neutron yield in the discharge. From the viewpoint of neutron emission the axial plasma jet plays the role of a moving thermonuclear reactor which satisfactorily simulates the properties of the experimental neutron radiation. The practically motionless plasma focus may also be an important contribution to the neutron yield. The general problem of the gas dynamic accelerating mechanism in pinches is discussed.

As a result of successful experimental investigations by Filippov et al.[1,2], interest in z-pinches has greatly increased in recent years. A large number of foreign experimental papers have been published[3-7], which while differing to a greater[3-5,7] or lesser[6] degree from the investigations described in[1,2], are essentially devoted to the study of the same phenomenon—a plasma focus. A plasma focus is an unusually dense high-temperature miniature plasma formation, produced near the geometrical axis of a system as a result of the cumulative effect.

Of great importance to the understanding of the physics of the phenomenon are the researches of Peacock et al.[5], who, in particular, determined spectroscopically (in the soft x-ray region) the duration of the existence of the plasma focus. In experiments by Maisonier et al.[6], the best focusing of the cumulation process was reached by precise axial symmetrization of the discharge. Record neutron yields in deuterium gas (up to 10^{11} per pulse) were obtained in the latest experiments of an Italian group[6] and by Mather and co-workers[7]. In the latter series of experiments, Filippov et al.[8] have demonstrated the contraction at the anode of the total current through the plasma to very small radii (less than 1 cm), and also the absence of a direct coupling between the hard x-rays and the neutron radiation.

A theoretical interpretation of the z-pinch with formation of a plasma focus was previously undertaken by the authors of[9] using a one-dimensional magnetohydrodynamic approach. The most significant comparisons with experiments in a deuterium plasma[1,2] were made in the same reference[9]. The results have made it clear that in the course of the experiments there occurs a large outflow of mass (not less than 90%) from the z-pinch region where the plasma focus is produced. Only under this assumption can the calculated radial velocities of plasma motion ($\sim 2.5 \times 10^7$ cm/sec) and temperatures (~ 1 keV) in the plasma focus be satisfactorily reconciled with the observations. However, in the one-dimensional theory with effective ejection of the mass, there still remain appreciable deviations from the experimental values of the radius of the plasma focus and the time of its existence. The theoretical value of the radius of the focus (~ 1 cm) was much larger than the observed radius (~ 1 mm), and the tentative time of existence of the high parameters ($\sim 5 \times 10^{-8}$ sec), to the contrary, was underestimated in comparison with the experimental value (~ 1–2×10^{-7} sec). It is quite clear that in fact the ejection of the mass proceeds continuously, in the form of a leakage, and the aggregate of the experimental facts indicates that this leakage develops in the direction of the axis of the system. Simultaneously, the z-pinch is characterized by a high degree of axial symmetry.

The foregoing circumstances have led to the formulation of a two-dimensional nonstationary magnetohydrodynamic problem, in which it is necessary to investigate the dependence of all the quantities on the cylindrical coordinates r and z. The present paper is devoted to a solution of the two-dimensional problem and to certain physical conclusions of this solution. The first stage of this computationally complicated problem was in essence the solution of the hydrodynamic problem with corresponding external boundary conditions. The electric conductivity of the plasma was assumed to be infinitely large, and the only dissipation process taken into account was the plasma viscosity, and furthermore in a simplified variant intended to blur the shockwave front over the local mean free path of the charged particles in the plasma. The relative role and significance of the different dissipation processes were clarified in the one-dimensional theory[9], and these data should be borne in mind in the two-dimensional case.

1. In the case of axial symmetry ($\partial/\partial\varphi = 0$) the equations of hydrodynamics of an ideal plasma ($E = (3/2)p/\rho$) with allowance for the viscous terms are written in a cylindrical coordinate system:

$$\rho\frac{du}{dt}+\frac{\partial}{\partial r}(p-q)=0, \qquad (1)$$

$$\rho\frac{dv}{dt}+\frac{\partial}{\partial z}(p-q)=0, \qquad (2)$$

$$\frac{dp}{dt}-\frac{5p-2q}{3\rho}\frac{d\rho}{dt}=0, \qquad (3)$$

$$\frac{d\rho}{dt}+\rho\left(\frac{\partial u}{\partial r}+\frac{\partial v}{\partial z}+\frac{u}{r}\right)=0, \qquad (4)$$

where the substantial derivative is

$$\frac{d}{dt}=\frac{\partial}{\partial t}+u\frac{\partial}{\partial r}+v\frac{\partial}{\partial z}.$$

These equations contain the radial and axial components of the velocity u and v, the scalar pressure p, the density ρ, and the diagonal element of the second-viscosity tensor q. The main role of the viscosity is to ensure a finite width of the shock-wave front, which is never less than the mean free path of the charged particles in the plasma—the ions and electrons. It is well known that in a fully ionized ideal plasma the coefficient of second viscosity is equal to zero[10]. However, for a local plane shock front the contribution of the first viscosity practically coincides quantitatively with the contribution of the second viscosity, if the corresponding coefficients are equal. With the exception of the thickness of the shock front, the introduced viscosity did not play an essential role. The quantity q was taken in the form

$$q=\alpha\left(\frac{p}{\rho}\right)^{5/2}\left(\frac{\partial u}{\partial r}+\frac{\partial v}{\partial z}+\frac{u}{r}\right)-\beta\rho\left(\frac{\partial u}{\partial r}+\frac{\partial v}{\partial z}+\frac{u}{r}\right)^{2}, \qquad (5)$$

but for the discharge conditions ($\partial u/\partial r + \partial v/\partial z + u/r \geq 0$) it was assumed that $q \equiv 0$. In the case when the mean free path turned out to be smaller than the spatial interval of the difference analog of Eqs. (1)–(4), the second term of (5) ensures blurring of the shock front over several intervals. Relation (5) contains the constant coefficients α and β, which will be defined later.

The common temperature of the ions and electrons of a deuterium plasma is

$$T = p/2\rho. \qquad (6)$$

For all quantities contained in the equations (1)–(6), the main measurement units are: the length R_0, the pressure p_0, and the density ρ_0. Then the velocity unit is $v_0 = \sqrt{p_0/\rho_0}$, the time unit is $t_0 = R_0/v_0$, the temperature unit is $T_0 = (m_d/k)p_0/\rho_0 = (2m_p/k)p_0/\rho_0$ (m_p—proton mass) and the electric current unit is $I_0 = c(2\pi R_0^2 p_0)^{1/2}$ (c—velocity of light).

The constant coefficient α in formula (5) is determined, in accordance with the foregoing, by means of the following relation for the second-viscosity coefficient:

$$\zeta=\rho c_c l_c=\rho\sqrt{\frac{3}{2}\frac{p}{\rho}}\,\frac{2m_p(kT)^2}{\rho e^4 L}=l_0\left(\frac{p}{\rho}\right)^{5/2}, \quad l_0=\sqrt{\frac{3}{2}}\,\frac{2m_p^3}{e^4 L}, \qquad (7)$$

where all the quantities are written in dimensional form, and L is the Coulomb logarithm, subsequently set equal to L = 20. As a result of a change over to dimensionless variables, we then obtain

$$\alpha=\frac{p_0^2}{\rho_0^3}\frac{l_0}{R_0}. \qquad (8)$$

Let us proceed to formulate the initial and boundary conditions of the problem in connection with a real discharge chamber[1)]. The numerical solution of the problem was obtained in the region (see Fig. 1):

$$0 \leqslant z \leqslant 0.545, \quad 0 \leqslant r \leqslant R(z,t),$$

where R(z, t) is the radius of the current sheath on which the boundary condition stipulates equality of the momentum flux in the plasma with a ponderomotive force[2)]:

$$p - q = I^2/R^2, \qquad (9)$$

where the dimensionless current approximating the experimental oscillogram is specified in the form of a function of the time:

$$I = 1 - 0.42t. \qquad (10)$$

On the remaining boundaries of the region, it was specified that the normal component of the velocity vanish:

$$r=0,\quad u=0;\quad z=0,\quad z=0.545,\quad v=0. \qquad (11)$$

It is easily seen that this boundary condition is sufficient also when second-viscosity terms are present in the equations. The surfaces z = 0 and z = 0.545 correspond to the electrodes of the chamber.

The initial conditions were chosen in the simplest manner:

$$t=0,\quad u=v=p=0,\quad \rho=1. \qquad (12)$$

The initial position of the current sheath R(z, 0) is shown in Fig. 1.

The foregoing formulation of the initial and boundary conditions of the problem has a simple physical meaning. On the basis of the experimental data (magnetic probe measurements), the starting point of the calcula-

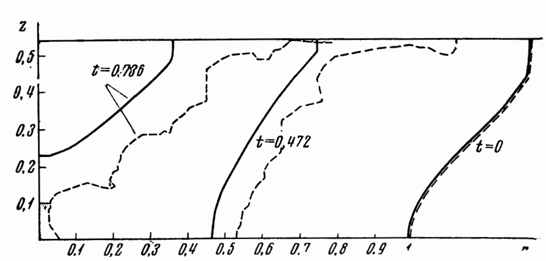

FIG. 1. General picture of two-dimensional plasma motion. For three instants of time, the solid lines represent the shock-wave front, and the dashed lines represent the current sheath.

1) V. D. Ivanov took part in the concrete choice of these conditions. The theoretical analysis and the experiment are briefly described in a joint paper [8].

2) More accurately speaking, it is necessary to equate separately two components of the momentum flux (axial and radial) to the corresponding projections of the ponderomotive force, directed at each point normally to the boundary. For an isotropic and diagonal viscosity tensor, this obviously reduces to the condition (9) given in the text.

tions was chosen to be a certain instant of time with a known configuration of the current sheath. This instant was sufficiently remote from the instants of time at which cumulation and formation of the plasma focus take place, so as to be able to neglect the plasma energy at t = 0 compared with its values in the succeeding instants of time.

The main task of our problem is to study the character of the continuous process of axial outflow of plasma from the region of plasma-focus formation, due to the experimentally observed non-cylindrical configuration of the current sheath for a certain intermediate instant of time. It is perfectly clear that many features of the phenomenon of a non-cylindrical z-pinch, including some very important ones, are not taken into account in such a simple formulation of the problem. But in a further refinement of the physical formulation of the problem it is natural to start from the results of the obtained solution, as is demonstrated in the concluding section of this paper.

We do not describe here the numerical solution method. The gist of the method of large particles employed here is contained in a paper by one of the authors[11].

2. The results of the numerical calculations are shown in Fig. 1—5. We assumed here the following measurement units: $R_0 = 13$ cm, $p_0 = 0.94 \times 10^7$ dyne/cm^2, $\rho_0 = 2.4 \times 10^{-7}$ g/cm^3, $v_0 = 0.625 \times 10^7$ cm/sec, $t_0 = 2.08 \times 10^{-6}$ sec, $T_0 = 82$ eV, and $I_0 = 10^6$ A. We note that this problem has a similarity property. The concrete choice of the measurement units affects only the value of the constant coefficient α, given by the relation (8): $\alpha = 5.25 \times 10^{-3}$, and also the rate of decrease of the total current in formula (10) with time. The other constant coefficient in (5) was set equal to $\beta = 5.76 \times 10^{-4}$. The discussion that follows is in terms of dimensionless units.

Figure 1 shows the general picture of plasma motion. We call attention to the almost periodic non-smoothness of the current sheath, which is particularly noticeable in the last instant of time. It is remarkable that, unlike the sheath, the shock-wave front retains its smooth character. The current sheath comes closest to the axis ($r \lesssim 0.02$) in the vicinity of z = 0.06 (Fig. 2). The location of the chains of plasma gas particles (they are calculated points) reveals that the plasma slips away upward from this region. The sections of the current sheath with $z \gtrsim 0.05$, against which an axial jet is produced, gradually overtake the narrow section with z = 0.05, which originally was ahead of them, and the region of minimum radius of the sheath is itself displaced somewhat upward (to z = 0.06).

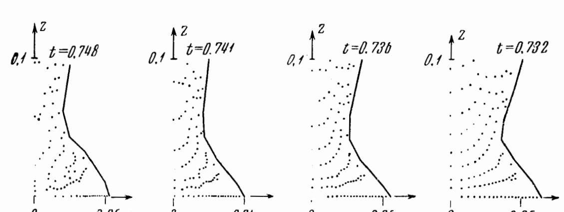

FIG. 2. Process of formation of plasma focus. The solid line represents the current sheath. The points denote gas particles.

Figure 3 shows the cross sections of the pressure p and the axial velocity v along the system axis (r = 0) for several instants of time. The instant t = 0.723 corresponds approximately to the emergence of the shock front to the symmetry axis. This occurs first at z = 0. Soon afterwards, a noticeable velocity $v \approx 1$ appears. Then, with initial intensity of heating and additional compression of the near-axis region ($t > 0.732$), the maximum of the pressure moves away from the anode in the interval z = 0.04—0.08 and reaches $p_{max} = 1600$ at the instant t = 0.741 at velocities v = 2—3. Far from the anode ($z \approx 0.2$), the axial velocity increases to much higher values, v = 10—11 (we note that its maximum is on the axis). The front of the upward-propagating shock wave is quite widely blurred. It can be assumed conditionally that its rear boundary (as determined from a limiting compression by a factor of four and from the contribution of the viscous terms to the momentum flux), for example for the instant t = 0.748, is located at z = 0.29, whereas the front boundary is located at z = 0.40. We note that in accordance with this the dimensionless mean free path is $l_K/R_0 = 4.3 \times 10^{-3} p^2/\rho^3 = 0.043$ at z = 0.29.

An analysis of the behavior of the quantities along the section z = 0.06, corresponding to the closest approach of the sheath to the system axis, demonstrates the following (see Fig. 4). Starting with the instant t = 0.732, when the pressure on the axis p_0 begins to exceed approximately double the pressure on the boundary p_C, a sharp deceleration of the radial compression begins (from t = 0.732 to t = 0.741 the

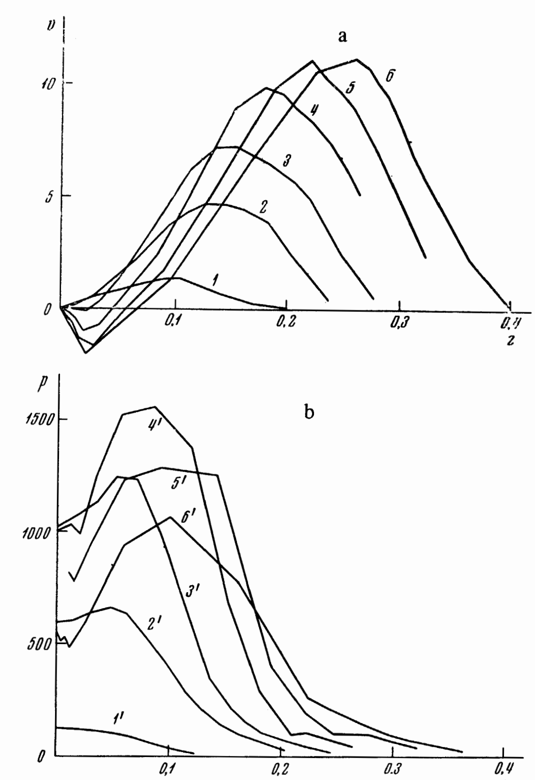

FIG. 3. Profiles of axial velocity v (a) and of the pressure p (b) as functions of the coordinate z on the symmetry axis (r = 0). 1, 1′ – t = 0.723; 2, 2′ – t = 0.732; 3, 3′ – t = 0.736; 4, 4′ – t = 0.741; 5, 5′ – t = 0.745; 6, 6′ – t = 0.748.

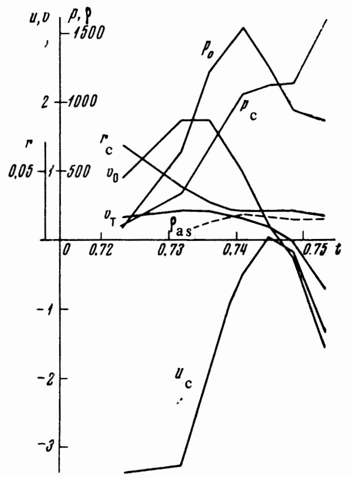

FIG. 4. Time dependence of the following quantities (in the section z = 0.06 where the plasma focus is produced): pressure and axial velocity on the axis (p_0, v_0); pressure, radial and axial velocity in the current sheath (p_c, u_c, v_c); the sheath radius (r_c); the density averaged over the cross section (ρ_{av}).

velocity of the boundary changes from $u_c = -3.3$ to $u_c = -0.5$). Owing to the considerable radial gradient of the component of the velocity u:

$$\left|\frac{\partial u}{\partial r}+\frac{u}{r}\right| > \frac{\partial v}{\partial z},$$

the density, and with it the plasma pressure, continues to increase in the near-axis region in this time interval. After the stopping of the sheath ($t \geq 0.745$), this gradient disappears completely, and the pressure inside the plasma, including also on the axis (p_0), begins to decrease, owing to the presence of a finite axial gradient of the velocity component $v(\partial v/\partial z > 0)$. By the instant $t = 0.748$, it becomes smaller than the limiting value p_c, which remains practically unchanged in the interval $0.741 < t < 0.748$. After equalization of the pressure ($t > 0.750$), a second compression of the plasma begins in which much larger parameters are reached ($p > 10^4$, $r_c < 0.01$). It must be emphasized that the described secondary compression of the plasma has a one-dimensional analog not in the adiabatic additional compression of the plasma described in[9], but in the subsequent second pulsation of the pinch (see, for example[12]). However, in the considered two-dimensional case the secondary compression follows practically immediately the first one (according to Fig. 4, where a plot of $r_c(t)$ is shown, the time interval between them is only 0.01), whereas in the one-dimensional theory the period of the pinch pulsations, as is well known[12] is more or less close to the time of the primary compression ~1.0.

The results of the calculation at $t > 0.750$ and the quantitative characteristics of the secondary compression of the plasma should be approached with caution for the following reasons. First, owing to the strong outflow of the plasma from the region of maximum compression of interest to us, too few calculated points remain in the region, and this affects adversely the accuracy of the calculation. Second, the solution of the problem was in fact stopped when the unceasing compression of the plasma led to a discontinuity and to an unbounded growth of the pressure at the location of the discontinuity. Although the accuracy of the calculation in this region is utterly insufficient, the very fact of unceasing secondary compression is worthy of attention. Further analysis shows that without taking into account the new physical factors, the secondary compression can occur without stopping up to a complete rupture of the plasma at a certain point on the axis. This shows clearly the important role of these new physical factors in general in the determination of the characteristics of the secondary compression. For these two reasons, we do not present here the results of the calculation for instants of time $t > 0.750$.

Let us list the most important results of the calculations (this time in dimensional units), without touching upon secondary compression, for the reasons indicated above. The plasma focus (the first maximum of the pressure on the axis) occurred 1.54 μsec after the start of the sheath from a radius of 13 cm on the anode. By that time the current dropped to 0.69×10^6 A. The maximum pressure was 1.5×10^{10} dyne/cm^2 = 1.5×10^4 atm at a compression of 180 times (averaged over the cross section), at an average temperature 0.4 keV, and an axial velocity 1.5×10^7 cm/sec. The dimensions and the position of the focus were: radius—2.5 mm, height above the anode—10 mm, longitudinal dimension—10 mm. The velocity of the plasma jet along the axis at a height 2.5 cm reaches 7×10^7 cm/sec, and its width is 2.5 mm (see Fig. 5).

The values of the temperature on the axis greatly exceed the temperature averaged over the cross section. For example, for $t = 0.741$ the maximum temperature at $z = 0.06$ is 1.25 keV, and in subsequent instants of time a second temperature maximum develops in a region adjacent to the rear boundary of the front of the axial shock wave. For $t = 0.748$ at $z = 0.27$, the temperature equals 0.98 keV, whereas on the axis of the plasma focus at $z = 0.07$ the temperature is 0.94 keV. A temperature minimum, 0.62 keV, is located between them at $z = 0.23$. However, the absolute values of the temperature on the plasma-focus axis may be overestimated. On the basis of the one-dimensional theory[9] it can be assumed that allowance for the ionic thermal conductivity would decrease the temperature on the axis of the focus by an approximate factor of three. Thus, in the estimate of the temperature in the plasma focus it is necessary to start, most probably, from the already mentioned mean tempera-

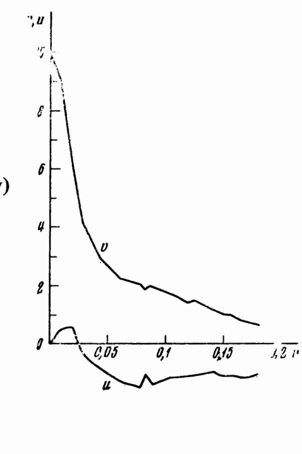

FIG. 5. Profiles of axial velocity (v) and radial velocity (u) as functions of the radius r for the sections of the axial plasma jet z = 0.18 at t = 0.741.

ture ~0.5 keV. These considerations do not apply to the temperature of the second axial maximum, since its magnitude must satisfy the Hugoniot condition, and the ionic specific heat could only decrease the dimension of the hot region.

3. From the described results of the two-dimensional hydrodynamic calculations we can draw certain physical conclusions. First, the calculations have illustrated the process of formation of the plasma focus, which obviously is the region where the plasma has its maximal parameters. Such a plasma focus occurs opposite the location where the current sheath comes closest to the symmetry axis. Although the maximal plasma parameters turned out to be much lower than those observed in the experiment (for example, a temperature 0.4 keV as against 1–2 keV), the secondary compression seems to be highly promising. It is important to note, that it lags the first relatively-weak compression very little in time ($\Delta t \approx 2.0 \times 10^{-8}$ sec). As noted by Filippov, the observed plasma focus is apparently characterized also by the occurrence of a more powerful secondary compression following an equally small time interval (see[2]).

Additional calculations of the plasma-focus parameters call for an improvement in the physical formulation of the problem. First, it is necessary to take into account the sharp increase of the inductive reactance (approximately two times) at the instant of formation of the focus. Such an affect can be naturally described by including in the calculations the electrotechnical equation, the solution of which gives the dependence of the total discharge current on the time, in lieu of relation (10) of the present calculation. It is clear that this effect will prevent an unceasing compression of the plasma. Second, account must be taken of the finite electric conductivity of the plasma. The existence and the structure of the plasma focus may depend on the Joule heating. We shall estimate below the influence of this heating on the dynamics of the focus.

We can assume on the basis of the calculations that ultimately the longitudinal dimension of the focus becomes much larger than its radius. Physically this is probably connected with the known stability of a plane shock front against short-wave perturbations of its shape[13], which cannot fail to be reflected in the development of the perturbations of the current sheath. For further estimates we assume a power-law variation of the average plasma density at the focus, $\overline{\rho} \sim R^{-\gamma}$. When $\gamma < 2$, this variation takes into account the longitudinal outflow of the plasma, and the smaller the exponent γ, the more intense this outflow. Without Joule heating, the average plasma pressure increases adiabatically, $\overline{p} \sim \overline{\rho}^{5/3} \sim R^{-5\gamma/3}$. It is obvious that this cannot overcome the boundary pressure, $p_c \sim R^{-2}$, when $\gamma < 6/5$, at which unceasing compression takes place. Let us estimate further the Joule heating. The characteristic thickness of the skin layer, for a fully ionized plasma with conductivity σ, is proportional to

$$d \approx \sqrt{D_m t_g} = \sqrt{\frac{c^2 R u_c^{-1}}{4\pi\sigma}} \sim \sqrt{\sigma^{-1} R}, \qquad D_m = c^2 / 4\pi\sigma, \tag{13}$$

since the characteristic growth time t_g in formula (13) must be chosen to be the growth time of the magnetic field on the boundary, which is proportional to the radius R in the case of an almost constant radial velocity u_c. The Joule heating per unit volume, in the case of strong skin-layer concentration, is then proportional to

$$Q \approx \frac{j^2}{\sigma} \sim \frac{1}{R^2 d^2 \sigma} \sim R^{-3}. \tag{14}$$

The dependence of the electron temperature on the sheath radius is determined from the energy equation

$$\rho \frac{dT_e}{dt} \sim R^{-\gamma} \frac{dT_e}{dR} \sim Q \sim R^{-3}$$

from which we get inside the skin layer (without allowance for the adiabatic heating)

$$T_e \sim R^{\gamma-2}, \quad p \sim \rho(T_i + T_e) \sim \rho T_e \sim R^{-2}, \tag{15}$$

which is sufficient to counteract the ponderomotive forces.

To complete these estimates, it is necessary, in retrospect, to justify the following two previously made assumptions: 1) the predominance of the Joule heating over the adiabatic heating when $\gamma < 6/5$; 2) the strong skin-layer concentration in the entire region of variation of the exponent γ. In the case of adiabatic heating, the temperature of the plasma changes like $\overline{T} \sim \overline{T}_e \sim \overline{T}_i \sim \overline{\rho}^{2/3} \sim R^{-2\gamma/3}$. Indeed, according to (15), the Joule heating predominates over the adiabatic heating precisely when $\gamma < 6/5$ $(\gamma - 2 < -2\gamma/3)$. Further, strong skin concentration takes place if $t_{\lim} \approx R^2 D_m^{-1} \sim R^2\sigma \gg t_g \sim R$. In other words, we should have $\beta < 0$ if $t_{\lim}/t_g \sim R^{\beta}$. It is easy to verify directly that the inequality $\beta < 0$ is satisfied for all values of γ.

In addition, in the estimate (15) it was tacitly assumed that the electron temperature of the skin layer does not become equalized during the course of the process with the ion temperature radially, in other words, the energy exchange of the electrons with the ions, and the electron thermal conductivity, are not effective. This assumption was justified in the one-dimensional calculations. Nonetheless, it can be readily shown that in the opposite case, that of complete equalization of the electron temperature in the radial direction, the plasma pressure cannot counteract the ponderomotive forces, although both auxiliary assumptions are valid here, too.

Thus, our estimates of the Joule heating indicate that in the case of a strong outflow of the plasma, precisely when the adiabatic heating does not ensure that the ponderomotive forces counteract the compression, the Joule heating can produce in principle, in conditions of strong skin concentration, a sufficient counter-reaction. In conjunction with the sharp decrease of the total current as a result of the increased inductance, conditions are created facilitating the stoppage of the plasma-focus compression. In contrast to the calculation made in the two-dimensional problem, it is necessary to take into account the finite electric conductivity of the plasma and to include the electrotechnical equation in the calculation. The thermal conductivity and the viscosity of the plasma also play a definite role in the dynamics of the plasma focus, but this role can hardly be as significant as the role of the indicated factors.

Although the problem of the plasma focus calls for further theoretical research, including new two-dimensional calculations, considerable interest attaches to data on the axial plasma jet. This phenomenon was observed during the first stage of the experiments with a non-cylindrical z-pinch[14]. It was determined there spectroscopically that the velocity in the jet exceeds by several times the velocity of the radial compression. In the present investigation we used the results of the calculation to develop further the hypothesis advanced in[14] concerning the role of the axial plasma jet in the production of the neutron radiation from the discharge. We have already pointed out two properties of the jet, namely the large velocity (several times larger than that of the radial compression) and the temperature maximum behind the shock front. If we take into account the fact that the deuteron velocity distribution inside the thickness of the shock front differs greatly from Maxwellian, then we can imagine, with some exaggeration, that the thickness of the shock front serves as a certain gas target for the fast deuterons moving at the macroscopic velocity of the plasma behind the front[3)]. We can then readily estimate the neutron yield:

$$W = n_1 v_1 \pi r_1^2 \sigma(v_1) n_0 \lambda_1 \tau, \qquad (16)$$

where n_1 is the density of the jet, v_1 the velocity of the plasma behind the front, r_1 the radius of the jet, $\sigma(v_1)$ the effective cross section of the dd reaction, n_0 the density of the plasma ahead of the front,

$$\lambda_1 = l_K(E_1) \approx \frac{E_1^2}{n_0 e^4 L} = \frac{(m_d v_1^2)^2}{4 n_0 e^4 L}$$

the mean free path of the fast deuterons in the stationary plasma, identified with the thickness of the shock front, and τ the lifetime of the jet. We assume on the basis of the calculation (see Fig. 5) $n_1 = 2.5 \times 10^{17}$ cm^{-3}, $n_0 = 6.3 \times 10^{16}$ cm^{-3}, $v_1 = 7 \times 10^7$ cm/sec, $r_1 = 0.25$ cm, and $\tau = 10^{-7}$ sec. Then $E_1 = 5.2$ keV, $\sigma(E_1) \simeq 10^{-31}$ cm^2 [17]. Substitution of all these numbers in formula (16) yields $W = 2 \times 10^6$ neutrons, which, of course, is very small compared with the observed yields[1-8]. However, it suffices to increase the deuteron velocity to 10^8 cm/sec, which is perfectly admissible, to increase the neutron yield W to 10^9 neutrons, which is already practically comparable with the observed yield.

The notion that the shock front is the target must, of course, not be taken too literally. The difference becomes quite clear, if it is recognized that immediately behind the front there is produced a plasma in thermodynamic equilibrium, with a dimensionless temperature, in accordance with the Hugoniot condition, $T = v^2/6$, i.e., approximately 2 keV[4)]. In such a plasma, a thermonuclear reaction should take place with a yield that increases with the width of this region. The corresponding neutron source will have the properties of the "moving boiler"[2]. Within the framework of these estimates, it is difficult to obtain any definite conclusion concerning the ratio and connection of the sources of the neutrons from the thermonuclear and target mechanisms. One can only indicate that the total yield of the neutrons in the plasma jet turns out to be the same order as the observed yield, the time of neutron emission is quite large, since it coincides with the lifetime of the jet ($\tau \approx 10^{-7}$ sec), the source grows in dimension along the axis to practically the entire distance between the electrodes, although its radius is very small. The neutron spectrum has a mixed character and reflects the combined contribution of the moving thermonuclear boiler and the target.

Simultaneously, a thermonuclear reaction takes place in the plasma focus itself. The plasma in this focus may have lower temperatures compared with the hot region of the jet, but on the other hand much larger densities of matter are attained in the focus (the degree of compression apparently reaches 10^3). Quantitative estimates of the neutron yield were made in the one-dimensional theory[9]. They can be regarded as proof of the appreciable contribution of the thermonuclear reactions in the region of the focus to the total observed neutron yield. Experimental material with measurement of the neutron time of flight demonstrates the existence of two different sources: one of the type of the moving thermonuclear boiler, and the other practically stationary; in this sense, the experimental material confirms the considerations advanced on the basis of the calculation[8]. In all probability, the neutron-measurement paradox noted in[7] can be resolved on the same basis.

It is quite possible that these mechanisms of the neutron emission were realized also in linear z-pinches. The problem of the use of the accelerating mechanism to explain the neutron emission in electric fields has not been completely solved. In the analysis of this question, given in the book[17], it was concluded that a major role is played in the formation of the electric fields by plasma-pinch constrictions, which are essentially these very plasma foci. On the basis of the foregoing, it is natural to advance an alternate hypothesis, that of gas dynamic acceleration of the deuterons in the plasma jets that accompany the constrictions.

In conclusion, we are sincerely grateful to N. V. Filippov for valuable advice and remarks, V. D. Ivanov for help in reduction of the experimental data, and L. A. Artsimovich and M. A. Leontovich for useful discussions of this work.

3)This representation corresponds literally to the method first proposed by Mott-Smith [15] for solving the problem of the structure of the shock front. However, this does not contradict in essence the description of the front structure based on equations of the Navier-Stokes type [16].

4)The calculated temperature turned out to be almost half as large. This inaccuracy is probably due to the sharp character of the temperature peak ($\Delta z = 0.05 - 0.1$).

[1] N. V. Filippov, T. I. Filippova, and V. P. Vinogradov, Nucl. Fusion Suppl., pt. II (1962).

[2] N. V. Filippov and T. I. Filippova, Paper at Second Internat. Conf. on Controlled Thermnuclear Reactions, Culham, England, CN-21/250, 1965.

[3] J. W. Mather, Phys. of Fluids, 8, 366 (1965).

[4] J. W. Mather and P. Bottoms, Phys. of Fluids 11, 611 (1968).

[5] P. D. Morgan, H. J. Peacock, R. J. Speer, and P. D. Wilcock, Third Conf. on Research in Plasma Physics and Controlled Thermonuclear Reactions, Novosibirsk, CN-24/G-4, 1968.

[6] Ch. Maisonnier, C. Gourlan, M. Haegi, J. G. Linhart, and M. Samuelli, ibid., CN-24/G-6, 1968.

[7] P. J. Bottoms, J. W. Mather, A. H. Williams, J. P. Carpenter, and K. D. Ware, ibid., CN-24/G-5, 1968.

[8] V. I. Agafonov, V. P. Vinogradov, G. V. Golub, L. G. Golubchikov, V. G. D'yachenko, V. D. Ivanov, V. S. Imshennik, Yu. A. Kolesnikov, É. B. Svirskiĭ, N. V. Filippov, and T. I. Filippova, ibid., CN-24/G-2, 1968.

[9] V. F. D'yachenko and V. S. Imshennik, Voprosy teorii plazmy (Problems of Plasma Theory) 5, 1967, p. 394.

[10] S. Chapman and T. G. Cowling, Mathematical Theory of Nonuniform Gases, Cambridge, 1962.

[11] V. F. D'yachenko, ZhVM i MF (J. of Comput. Math. and Math. Phys.) 5, 680 (1965).

[12] S. I. Braginskiĭ, I. M. Gel'fand, and R. P. Fedorenko, Fizika plazmy i problema upravlyaemykh termoyadernykh reaktsiĭ (Plasma Physics and Problem of Controlled Thermonuclear Reactions) 4, 1958, p. 201.

[13] S. P. D'yakov, Zh. Eksp. Teor. Fiz. 27, 288 (1954).

[14] D. P. Petrov. N. V. Filippov, T. I. Filippova, and V. A. Khrabrov. Fizika plazmy i problema upravlyaemykh termoyadernykh reaktsiĭ (Plasma Physics and Problem of Controlled Thermonuclear Reactions) 4, 1958, p. 170.

[15] H. M. Mott-Smith, Phys. Rev. 82, 885 (1951).

[16] V. S. Imshennik, PMTF No. 1, 15 (1968).

[17] L. A. Artsimovich, Upravlyaemye termoyadernye reaktsii (Controlled Thermonuclear Reactions), Fizmatgiz, 1961.

Translated by J. G. Adashko
202