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DISPERSION OF A PLASMA CLUSTER OVER A LESS DENSE PLASMA BACKGROUND

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A number of theoretical reports have been devoted to a study of the question of the motion of plasma clusters over a plasma background (see [1-3], for example). The influence of finite conductivity (Coulomb and anomalous) on cluster motion and magnetic field growth is investigated in the present report. The problem of expansion of a plasma cylinder into a rarefied plasma is modeled numerically. Since in a first approximation we are interested not in the structure of the shock wave but in the process of plasma interaction with the magnetic field and the influence on this interaction of ion-sonic and beam instabilities, we chose the model of two-fluid hydrodynamics, which is simpler than the hybrid model.

We consider the nonsteady axisymmetric problem, when all the unknown functions are functions of the radius r and time t and the magnetic field has only a z component, directed along the axis of the cylindrical coordinate system: $\mathbf{H} = \{0, 0, H\}$. Then the system of equations of two-fluid magnetohydrodynamics [4] for the cylindrical case in Eulerian coordinates can be written in the form

$$\left(\frac{\partial}{\partial t} + u\frac{\partial}{\partial r}\right)n = -n\left(\frac{\partial u}{\partial r} + \frac{u}{r}\right), \quad (m_i + m_e)\,n\left(\frac{\partial}{\partial t} + u\frac{\partial}{\partial r}\right)u = \tag{1}$$

$$= -\frac{\partial}{\partial r}\left(nT_i + nT_e + \frac{H^2}{8\pi}\right) + \frac{c^2 m_i m_e}{16\pi^2 e^2 (m_e + m_i)\,n}\,\frac{1}{r}\left(\frac{\partial H}{\partial r}\right)^2,$$

$$\left(\frac{\partial}{\partial t} + u\frac{\partial}{\partial r}\right)H = -\left(\frac{\partial u}{\partial r} + \frac{u}{r}\right)H - \frac{c^2 m_e m_i}{4\pi e^2 (m_e + m_i)}\left(\frac{\partial}{\partial r} + \frac{1}{r}\right)\times$$

$$\times\left(\frac{\partial}{\partial t} + u\frac{\partial}{\partial r} + \frac{u}{r}\right)\left(\frac{1}{n}\,\frac{\partial H}{\partial r}\right) - \frac{c}{e}\left(\frac{\partial}{\partial r} + \frac{1}{r}\right)\frac{R_\varphi}{n},$$

$$\frac{3}{2}n\left(\frac{\partial}{\partial t} + u\frac{\partial}{\partial r}\right)T_e = -nT_e\left(\frac{\partial u}{\partial r} + \frac{u}{r}\right) - \left(\frac{\partial}{\partial r} + \frac{1}{r}\right)q_{er} + Q_e,$$

$$\frac{3}{2}n\left(\frac{\partial}{\partial t} + u\frac{\partial}{\partial r}\right)T_i = -nT_i\left(\frac{\partial u}{\partial r} + \frac{u}{r}\right) - \left(\frac{\partial}{\partial r} + \frac{1}{r}\right)q_{ir} + Q_i,$$

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where $R_\varphi = \left(\alpha_\perp \frac{c}{4\pi en} \frac{\partial H}{\partial r} + \beta_\Lambda^{uT} \frac{\partial T_e}{\partial r}\right)$ is a component of the frictional force between the ion and and electron components; $Q_i = (3nm_e\nu_{ei}/m_i)(T_e - T_i)$; $Q_e = -(Q_i + R_\varphi(v_{e\varphi} - v_{i\varphi}))$ (Q_i and Q_e are the heat release in the ion and electron gases); $q_{er} = -\left(\beta_\Lambda^{Tu} \frac{c}{4\pi en} \frac{\partial H}{\partial r} + \chi_\perp^e \frac{\partial T_e}{\partial r}\right)$ and $q_{ir} = -\chi_\perp^i \frac{\partial T_i}{\partial r}$ are the radial components of the electronic and ionic heat fluxes, respectively; $\alpha_\perp$, β_Λ^{uT}, β_Λ^{Tu}, $\chi_\perp^e$, and $\chi_\perp^i$ are defined just as in [4] for an arbitrary degree of plasma magnetization for $Z = 1$; $\mathbf{u} = (m_i\mathbf{v}_i + m_e\mathbf{v}_e)/(m_e + m_i)$ is the average-mass velocity. Here we should note that since $u_\varphi|_{t=0} = 0$ at the initial time (no rotation of the plasma as a whole), $u_\varphi = 0$ at any time.

In deriving the system of equations (1) we used the following relations:

$$\frac{\partial H}{\partial r} = -\frac{4\pi en}{c}(v_i - v_e), \quad \frac{\partial H}{\partial t} = -\frac{c}{r}\frac{\partial}{\partial r}(rE_\varphi).$$

We represent the effective collision frequency ν in the form of terms describing Coulomb and collective collisions. As an analysis shows [5], the latter are connected with the development of beam and ion-sonic instabilities: $\nu = \nu_{ei} + \nu_b + \nu_s$.

For the effective frequency connected with beam instability we obtain

$$\nu_b = \begin{cases} \left\{(v^2 - v_{Te}^2)\Big/\left[\left(\frac{m_e}{m_i}\right)^{1/3} v^2 + v_{Te}^2\right]\right\}\Omega_i & \text{for} \quad v \geqslant v_{Te}, \\ 0 & \text{for} \quad v < v_{Te}. \end{cases}$$

Here and later $v = \left|\frac{c}{4\pi en}\frac{\partial H}{\partial r}\right|$; $v_{Te} = \left(\frac{5}{3}\frac{T_e}{m_e}\right)^{1/2}$; $\Omega_i = \left(\frac{4\pi ne^2}{m_i}\right)^{1/2}$; $\nu_s = \nu_0(1 - K_1T_i/T_e)$ and $\left(1 - \frac{K_2(T_i/T_e)^{3/2}}{(v/c_s)}\right)$ for $T_e \geqslant K_1T_i$, $v > c_sK_2(T_i/T_e)^{3/2}$.

If these conditions are not satisfied, then $\nu_s = 0$, where $K_1 \approx 5$-7; $K_2 \approx 75$; $\nu_0 = (10^{-3}$-$10^{-2})T_e v\Omega_i/T_ic_s$; $c_s = (T_e/m_i)^{1/2}$ (see [5, 6]).

To complete the formulation of the problem, we must set up the initial and boundary conditions. At the initial time let the quiescent plasma fill a cylindrical chamber ($0 < r \leqslant r_0$), with the density of the plasma and its electron temperature being higher near the axis than at the periphery (cluster); the magnetic field is uniform everywhere and equal to H_0. Then the initial conditions can be written as follows:

$$u(r, 0) = 0, \tag{2}$$
$$T_e(r, 0)/T_e^0 = 1 + (A_T - 1)/(1 + \exp((\alpha/r^0)(r - r_1))),$$
$$n(r, 0)/n^0 = 1 + (A_\rho - 1)/(1 + \exp((\alpha/r^0)(r - r_1))).$$

The values of T_e°, n°, and r_0 will be determined below. The parameters A_ρ and A_T give the excess density and temperature in the cluster over the density and temperature in the background. The boundary conditions at the chamber axis are

$$u(0, t) = \partial H(0, t)/\partial r = \partial T_e(0, t)/\partial r = \partial T_i/\partial r = 0. \tag{3}$$

The boundary conditions at the plasma—metal interface are

$$u(r_0, t) = 0; \tag{4}$$

$$\partial T_e(r_0, t)/\partial r = \partial T_i(r_0, t)/\partial r = 0 \tag{5}$$

the absence of heat flux through the wall:

$$\partial H(r_0, t)/\partial r = 0. \tag{6}$$

The latter condition follows from the assumption of infinite conductivity of the chamber wall and the generalized Ohm's law.

Since it is convenient to analyze shock waves in a plasma in Lagrangian mass coordinates, the system (1) with initial (2) and boundary conditions (3)-(6) was written in these coordinates. For the numerical solution we used a scheme which is a generalization of the

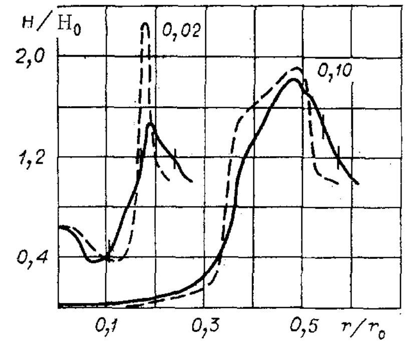

Fig. 1

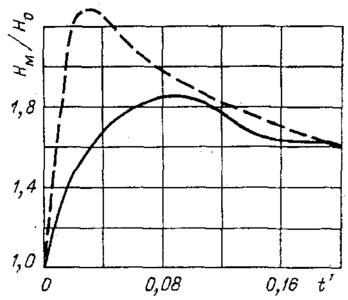

Fig. 2

"cross" gasdynamic scheme [7] and is described in detail for the plane case in [6] and for the cylindrical case in [8].

The values of the Eulerian coordinate and velocity were calculated at half-integral points ($r_{j-1/2}$, $u_{j-1/2}$) while all the remaining functions were calculated at integral points.

To increase the stability, all the dissipative terms were transposed to the top layer.

Let us examine the results of calculations of the stated problem with the following parameters: r_0 = 50 cm, T_e^0 = 100 eV, T_i^0 = 10 eV, $n^0 = 10^{12}$ cm^{-3}, H_0 = 30 Oe. In this case we will have $v_A^0 = 6 \cdot 10^6$ cm/sec and $t^0 = 8 \cdot 10^{-6}$ sec. At the time t = 0 the magnetic field was assigned as uniform everywhere, the ion temperature was constant, and $T_e(0, r)$ and $n(0, r)$ were given in accordance with Eq. (2) with the following coefficients: $A_\rho = 10$, $A_T = 5$, $\alpha = 100$, $r_1 = 0.1r_0$. The dimensionless pressure drop between the cluster and the background equals 220. The increased pressure at the center of the region causes plasma motion toward the periphery, which alters the initial magnetic field configuration.

The space—time dependence of the magnetic field is presented in Fig. 1. Here and in the other figures the solid lines correspond to an effective collision frequency ν, while the dashed lines are when only Coulomb collisions are taken into account. The numbers on the curves show the values of the time $t' = t/t^0$.

It is seen from Fig. 1 that the inclusion of anomalous conduction considerably decreases the amplitude of the magnetic field, especially at the initial time. The regions where ion-sonic instability develops are marked by vertical lines on the graph. In the course of time the region of development of ion-sonic instability decreases, while the field amplitude with anomalous conduction differs little from the field amplitude when only Coulomb conduction is taken into account, which is confirmed by Fig. 2, where the time dependence of the maximum value of H is given.

To monitor the accuracy of the solution we used the law of conservation of total energy, which is written for the system (1) in the form

$$W = \int_0^{r_0} \left\{ n(m_e + m_i)\frac{u^2}{2} + \left(\frac{\partial H}{\partial r}\frac{c}{4\pi e n}\right)^2 n m_e + \frac{H^2}{8\pi} + \frac{3}{2} n (T_e + T_i) \right\} r dr.$$

The value of W varied by no more than 0.2% in the calculations.

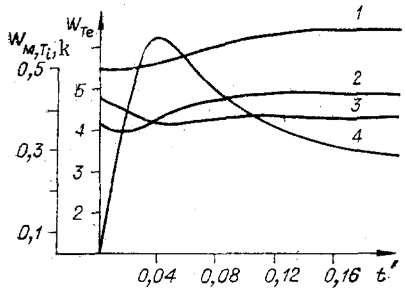

Fig. 3

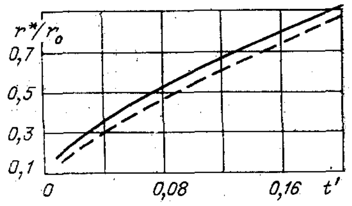

Fig. 4

The time dependence of the magnetic-field energy W_M, the thermal energies of the ion and electron gases W_{T_i} and W_{T_e}, and the kinetic energy W_k of the plasma is shown in Fig. 3 (curves 1-4, respectively). Here

$$W_M = \int_0^{r_0} [H^2/8\pi W^0]\, r dr, \quad W_{T_i} = \int_0^{r_0} (1{,}5 n T_i/W^0)\, r dr,$$

$$W_{T_e} = \int_0^{r_0} (1{,}5 n T_e/W^0)\, r dr,$$

$$W_k = \int_0^{r_0} \left\{ \left[n (m_e + m_i) \frac{u^2}{2} + \left(\frac{\partial H}{\partial r} \frac{c}{4\pi e n} \right)^2 n m_e \right] \Big/ W^0 \right\} r dr,$$

where $W^0 = H_0^2/8\pi$.

It is seen from Fig. 3 that the thermal energy of the electron gas decreases and is partially converted into an increase in magnetic-field energy and plasma kinetic energy, while the thermal energy of the ion gas also increases somewhat. Starting with $t \sim 0.16t^0$, the energy of each type remains about constant.

When only Coulomb conduction is taken into account, the form of variation of W_M, W_k, W_{T_e}, and W_{T_i} remains the same as in the preceding case, only greater heating of the ion gas is observed.

For both Coulomb and anomalous conduction the magnetic-field energy grows by about 20% compared with the initial value.

The time dependence of the cluster radius for Coulomb and anomalous plasma conduction is shown in Fig. 4. By the cluster radius we understand the position of the leading front of the density profile. It is seen that the velocity of the leading front decreases from about $t \approx 0.07t^0$.

As the calculations showed, the velocity of the leading front at the initial time is $\sim 7v_A^0$ for anomalous and $\sim 5v_A^0$ for Coulomb conduction. Starting with about $t \sim 0.1t^0$ the velocity of the leading front remains constant in both cases and comprises $\sim 3v_A^0$.

In conclusion, it must be mentioned that this work was initiated by the experiments described in [9, 10], and the authors consider it their duty to thank A. G. Ponomarenko and Yu. P. Zakharov for a discussion and stimulating debates.

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CURRENT SWITCHING FROM A FAST TRIP INTO A SHUNT WIRE

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In multistage current switching in a circuit containing an inductive energy store IES, the fast trip FT is shunted by one or more exploding wires [1-3]. To provide short times in the closing stage (on transferring the current to the load), it is desirable to perform the switching in the first stage at the maximum permissible current density in the shunting wire SW. We have to determine the maximum permissible current densities, since there are physical factors that prevent the transfer of a large current to a conductor of small cross section. So far, it has been considered that the main reason for failure at high current densities in the SW is breakdown in the FT at the stage when its electrical strength is recovering. In fact, as the density j_0 of the current switched into the wire increases, the time from the arc quenching in the FT to the explosion of the SW falls in proportion to j_0^{-2}, and therefore at high current densities the arc gap does not have time to recover its electrical strength by the instant of explosion, which leads to breakdown in the FT. However, we show here that the FT can fail also in an earlier stage, when the current transfer to the SW is not completed. The study is a theoretical consideration of the constraints in switching a current from an FT to SW arising from the rapid heating of the SW and the marked increase in resistance at high current densities.

1. The increase in resistance in the initial stages of electrical explosion in a conductor is [4, 5] determined in the main by the specific energy deposition q = Q/m, where $Q = \int_0^t i^2 R dt$ is the total deposited energy and m is the mass of the conductor. The increase in resistance is related to the energy input rate and is very marked at the stage of electrical explosion, but it does not play a large part in the initial stages (for $R/R_0 \leqslant 15$ for copper and for $R/R_0 \leqslant 11$ for aluminum). If the energy input rate to the conductor is small by comparison with the time for the phase transition from the solid state to the liquid one), then the dependence of the relative resistance on the specific energy deposition in

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