COLLISIONLESS DISPERSION OF AN IONIZED CLOUD INTO A HOMOGENEOUS MAGNETIZED PLASMA

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Introduction

A plasma cloud expanding with a velocity $u_0 \approx 10^7$-10^8 cm/sec has been formed in a number of laboratory experiments [1-3] and active experiments in space [4, 5]. The surrounding rarefied medium is in a uniform magnetic field H_0 and is ionized by the x-ray and ultraviolet radiation escaping from the cloud.

The problem of the retardation of a plasma cloud during expansion into an empty space in which there is a magnetic field was analyzed in [6-9]; it was shown that the characteristic radius R_0 of retardation of the cloud is determined from the equation $R_0 = (Nm_1u_0^2/H_0^2)^{1/3}$ (N is the total number of ions in the cloud and m_1 is the mass of the cloud ions).

When an external plasma is present, one must allow for its influence on the motion of the cloud ions. Simple energy estimates shows [10] that if the dispersing ions interact intensely with the surrounding plasma then the characteristic distance in which the cloud is retarded is $R_* = (3N/4\pi n_*)^{1/3}$ (n_* is the ion concentration of the surrounding plasma). If $R_0 \ll R_*$, then the influence of the external plasma on the motion of the cloud is negligibly small, so that the data of [6-9] can be used in the interpretation of experimental results. The condition $R_0 \ll R_*$ is satisfied if the Alfvén Mach number is $M_A = u_0/v_A \ll 1$ ($v_A = H_0/\sqrt{4\pi n_* m_2}$ is the Alfvén speed in the external plasma and m_2 is the mass of ions of the external plasma).

The problem of a cylindrical explosion in a rarefied plasma with $M_A > 1$ was solved in [11] on the basis of the equations of [12]. The use of these equations to describe the plasma motion in the experiments of [1-5] presumes a number of restrictions and reservations. First, this is because a shock wave forms in the external plasma when $M_A > 1$ [11], and there are presently no generally accepted relations like the classical Rankine–Hugoniot conditions for a collisionless shock wave front in a plasma. Second, the parameters of these experiments are such that the Larmor radius of cloud ions often proves to be greater than R_*. Therefore, the Chew–Goldberger–Low equations cannot be used to describe the flow at scales of $\sim R_*$, and just these scales are the most interesting in a study of the mechanism of interaction of a dispersing cloud with a surrounding plasma.

The problem of a cylindrical explosion in a rarefied plasma is solved numerically in the present report on the basis of a one-dimensional hybrid model [13-15] in which the ion motion is described by the Vlasov equation, while the electron component is described as a massless fluid. It is assumed that the magnetic field strength vector H_0 is directed along the axis of symmetry; the particle velocities are perpendicular to H_0. Of course, in such a statement (a cylindrical explosion) one cannot obtain a quantitative description of the experimental results of [1-5], in which the initial energy release is concentrated at a point; however, the qualitative pattern of the developing flow and the mechanism of interaction of the dispersing cloud with the magnetized external plasma can be investigated in sufficient detail.

1. Description of the Model and Statement of the Problem

Within the framework of the hybrid model, the motion of the plasma ions is described by the Vlasov equation

$$\frac{\partial f_\alpha}{\partial t} + \mathbf{v}\frac{\partial f_\alpha}{\partial \mathbf{r}} + \frac{e_\alpha}{m_\alpha}\left(\mathbf{E} + \frac{1}{c}[\mathbf{v} \cdot \mathbf{H}]\right)\frac{\partial f_\alpha}{\partial \mathbf{v}} = 0, \tag{1.1}$$

where f_α is the distribution function for ions of the type α. The index α will distinguish ions of the dispersing cloud ($\alpha = 1$) and ions of the plasma of the medium ($\alpha = 2$). The average characteristics of the ion motion are determined by the equations

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$$n = \sum_{\alpha} d\mathbf{v} \cdot f_{\alpha}, \quad \mathbf{u} = \frac{1}{n} \sum_{\alpha} \int d\mathbf{v} \, (\mathbf{v} \cdot f_{\alpha}). \tag{1.2}$$

Henceforth we will assume that the plasma is singly ionized, so that the condition of local electrical neutrality of the plasma has the following form:

$$n_e = n = \sum_{\alpha} \int d\mathbf{v} \cdot f_{\alpha}. \tag{1.3}$$

The electron component of the plasma is described as a massless fluid, so that the equation of motion for it is

$$-e\left(\mathbf{E} + \frac{1}{c}[\mathbf{v}_e \cdot \mathbf{H}]\right) + m_e \nu (\mathbf{u} - \mathbf{v}_e) - \frac{1}{n_e} \nabla p_e = 0, \tag{1.4}$$

where ν is the frequency of collisions of electrons with ions (classical or anomalous [14]); p_e is the electron pressure.

Equations (1.1)-(1.4), together with the equation for the electron pressure p_e

$$\frac{\partial p_e}{\partial t} + (\mathbf{v}_e \cdot \nabla) p_e = -\gamma_e p_e \operatorname{div} \mathbf{v}_e + (\gamma_e - 1) \sigma^{-1} j^2 + (\gamma_e - 1) \left(\operatorname{grad} k \ \operatorname{grad} \frac{p_e}{n}\right) + k \Delta \frac{p_e}{n} = 0$$

(σ and k are the conductivity and electron thermal conductivity, respectively [14], and $\gamma_e = 5/3$) and the Maxwell equations

$$\operatorname{rot} \mathbf{H} = \frac{4\pi}{c} e \sum_{\alpha} \int d\mathbf{v} \, (\mathbf{v} \cdot f_{\alpha}) - \frac{4\pi}{c} e n_e \mathbf{v}_e \equiv \frac{4\pi}{c} \mathbf{j}, \tag{1.5}$$

$$\operatorname{rot} \mathbf{E} = -\frac{1}{c} \frac{\partial \mathbf{H}}{\partial t},$$

in which the displacement current is omitted, form a complete system of equations to describe the motion of a rarefied plasma.

If we are confined to the investigation of plasma flows with $\beta_e = 8\pi n_e T_e / H^2 \ll 1$ then the system (1.1)-(1.5) can be reduced to the form

$$\frac{\partial f_{\alpha}}{\partial t} + \mathbf{v} \frac{\partial f_{\alpha}}{\partial \mathbf{r}} + \frac{\mathbf{F}}{m_{\alpha}} \frac{\partial f_{\alpha}}{\partial \mathbf{v}} = 0; \tag{1.6}$$

$$\mathbf{F} = \frac{e[(\mathbf{v} - \mathbf{u}) \cdot \mathbf{H}]}{c} + \frac{[\operatorname{rot} \mathbf{H} \cdot \mathbf{H}]}{4\pi n} + \frac{e\eta_m}{c} \operatorname{rot} \mathbf{H}; \tag{1.7}$$

$$\frac{\partial \mathbf{H}}{\partial t} = \operatorname{rot}\left\{[\mathbf{u} \cdot \mathbf{H}] - \eta_m \cdot \operatorname{rot} \mathbf{H} - \frac{c}{4\pi e n}[\operatorname{rot} \mathbf{H} \cdot \mathbf{H}]\right\}, \tag{1.8}$$

where $\eta_m = c^2/4\pi\sigma$ is the magnetic viscosity ($\sigma = e^2 n / m_e \nu$).

The following can be taken as the initial ion distribution functions in the problem of a cylindrical explosion:

$$f_1(\rho, v_\rho, v_\varphi, t = 0) = \frac{N\delta(\rho)}{\pi\rho u_0^2} v_\rho \delta(v_\varphi) \tag{1.9}$$

$$(0 \leqslant v_\rho \leqslant u_0),$$

where N is the total number of cloud ions per unit length along the axis of symmetry ($[N] = \mathrm{cm}^{-1}$), and

$$f_2 = n_* \delta(v_\rho) \delta(v_\varphi). \tag{1.10}$$

An initial distribution function chosen in the form of (1.9) for the cloud ions assures that the solution of the kinetic equation as $t \to 0$ coincides with the asymptotic solution of the problem of the dispersion of a gas cylinder into a vacuum [16], while the gasdynamic characteristics of the medium obtained from (1.10) coincide with the initial data of the problem of an explosion in a medium without a counterpressure [11]. We note that in a cylindrical explosion the characteristic radius of retardation R_* introduced above is determined by the equation $R_* = \sqrt{N/\pi n_*}$.

The solution of the stated problem depends only on the distance ρ to the axis of the cylinder, so that the system of equations (1.6)-(1.8) is considerably simplified. By introducing the component A_φ of the vector potential of the electromagnetic field we obtain, in place of (1.8),

$$\frac{\partial A_\varphi}{\partial t} + \frac{u_\rho}{\rho} \frac{\partial}{\partial \rho} (\rho A_\varphi) = \eta_m \left\{\frac{\partial^2 A_\varphi}{\partial \rho^2} + \frac{1}{\rho} \frac{\partial A_\varphi}{\partial \rho} - \frac{A_\varphi}{\rho^2}\right\}, \tag{1.11}$$

where

$$A_\varphi(\rho,\ t=0)=H_0\rho/2. \qquad (1.12)$$

The sole vector component of the magnetic field strength is $H_z=\frac{1}{\rho}\frac{\partial}{\partial\rho}(\rho A_\varphi)$. For the kinetic equation (1.6) the system of characteristics has the form

$$\frac{d\rho}{dt}=v_\rho, \qquad (1.13)$$

$$\frac{dv_\rho}{dt}=\frac{v_\varphi^2}{\rho}+\frac{eH_z}{m_\alpha c}(v_\varphi-u_\varphi)-\frac{1}{nm_\alpha}\frac{\partial}{\partial\rho}\frac{H_z^2}{8\pi},$$

$$\rho\left(v_\varphi+\frac{eA_\varphi}{m_\alpha c}\right)=P_\alpha(0)=\text{const.}$$

The value of the generalized momentum $P_\alpha(0)$ is determined by the initial conditions (1.9), (1.10), and (1.12).

The dissipative processes due to the finite conductivity of the plasma do not play an important role if the magnetic Reynolds numbers are large ($Re_m \gg 1$). The characteristic spatial scale in this problem is the Larmor radius R_H of the ions, while the velocity scale is u_0. Using the Buneman equation [17] to estimate the conductivity, for $t \gg R_*/u_0$ we obtain the following estimate for $Re_m=\frac{u_0R_H}{\eta_m}$: $Re_m \approx M_A u_0 m_i/cm_e$. For $M_A \gg 1$, therefore, one can set $\eta_m=0$ in (1.11), and this was done in the numerical solution of the system (1.11)-(1.13).

2. Mechanism of Energy Transfer from the Cloud to the Surrounding Plasma

In the early stage of dispersion ($t \ll \Omega_\alpha^{-1}$) one can assume that the cloud ions disperse freely, i.e., for $\rho < u_0t$

$$f_1(\rho,\ v_\rho,\ v_\varphi,\ t)=n_*\left(\frac{R_*}{u_0t}\right)^2\delta\left(v_\rho-\frac{\rho}{t}\right)\delta(v_\varphi). \qquad (2.1)$$

The ions of the surrounding plasma are immobile, i.e., their distribution function coincides with (1.10).

Two-velocity ion flow develops in the region of $\rho < u_0t$. The electrons, owing to the small mass, undergo drift motion. The radial electron velocity is determined from the equation

$$v_{\rho e}=u_\rho=\frac{\rho}{t}\left\{1+\left(\frac{u_0t}{R_*}\right)^2\right\}^{-1},\qquad \rho<u_0t. \qquad (2.2)$$

For larger radii ρ, $v_{\rho e}=u_\rho=0$.

Equation (1.11) for the component A_φ of the vector potential of the electromagnetic field, with u_ρ defined by Eq. (2.2) and $\eta_m=0$, has the following solution:

$$A_\varphi(\rho,\ t)=\begin{cases}0, & 0<\rho<R_1(t),\\ H_0\frac{\rho}{2}\left[1+\left(\frac{R_*}{u_0t}\right)^2-\left(\frac{R_*}{\rho}\right)^2\right], & R_1(t)<\rho<R(t),\\ H_0\frac{\rho}{2}, & \rho>R(t),\end{cases} \qquad (2.3)$$

where $R_1(t)=u_0t\{1+(u_0t/R_*)^2\}^{-1}$; $R(t)=u_0t$.

Thus, at the boundary $R(t)=u_0t$ between one- and two-velocity flows there is a jump ΔH_z in the magnetic field strength vector equal to

$$\Delta H_z=H_z-H_0=H_0(R_*/u_0t)^2.$$

The magnetic field is fully displaced from the region of $\rho < R_1(t)$. In the region of $R_1 < \rho < R$ a uniform magnetic field is formed with a strength which varies with time by the following law:

$$H_z=H_0\{1+(R_*/u_0t)^2\}. \qquad (2.4)$$

The result obtained has a simple interpretation. The electrons of the cloud plasma and the external plasma do not mix because of the small Larmor radius. With free dispersion of the cloud ions the boundary $R_1(t)$ separating the electrons can be found from the equation

$$N = 2\pi \int_0^{R_1} d\rho\rho n_e = \pi R_1^2 n_* \left\{1 + \left(\frac{R_*}{u_0 t}\right)^2\right\},$$

which follows from the condition of local electrical neutrality of the plasma. In the dispersion of the cloud all the electrons of the external plasma which are set into motion are in a cylindrical layer bounded by the radii $R_1(t)$ and $R(t)$ at a time t. The electron concentration in this region is $n_e = n_* \{1 + (R_*/u_0 t)^2\}$. Since the magnetic field is "frozen" into the electron component of the plasma (1.4), $H_z = H_0 n_e / n_*$ $(R_1 < \rho < R)$, and this gives Eq. (2.4). The magnetic flux connected with plasma electrons of the cloud is equal to zero. Since by the time t these electrons expand to a radius $R_1(t)$, $H_z = 0$ at $\rho < R_1$. In the region of $\rho > R$ the magnetic field evidently is not disturbed.

It should be noted that as $t \to 0$ there is a singularity $(H_z|_{t\to 0} \sim H_0(R_*/u_0 t)^2)$ in the solution of (2.4) connected with the approximation $\eta_m = 0$ which was used. It is obvious that dissipative effects cannot be neglected if

$$\Delta R(t) = R(t) - R_1(t) = \left(\int_0^t \eta_m(t')\,dt'\right)^{1/2} \tag{2.5}$$

i.e., if the size ΔR of the region of a "compressed" field is less than the distance to which the magnetic field penetrates, owing to the finite conductivity of the plasma. One can estimate the time t_0 starting with which the solution (2.3) is applicable. Using the Buneman equation for the conductivity, from (2.5) we obtain $t_0 \lessgtr (R_*/u_0)(R_*\omega_{pi}/c)^{-1/2}$. For the parameters of the experiments of [1-5] $u_0 t_0 / R_* \ll 1$.

In the dispersion stage under consideration (in a first approximation with respect to M_A^{-2}) only a vortical electric field E_φ, different from zero in the region of $R_1 < \rho < R$, develops:

$$E_\varphi = H_0\{(R_*/u_0 t)^2(\rho/ct)\}.$$

In a time $\Omega_2 t \ll 1$ the ions move little from their initial position ρ_0, and therefore, using the conservation of the component $P_2 = \rho(v_\varphi + eA_\varphi/m_2 c)$ of the generalized momentum (1.13) and the expression (2.3) for A_φ we obtain

$$v_\varphi(t) = \begin{cases} \Omega_2 \dfrac{\rho_0}{2}, & 0 < \rho_0 < R_1, \\ \Omega_2 \dfrac{\rho_0}{2}\left[\left(\dfrac{R_*}{\rho_0}\right)^2 - \left(\dfrac{R_*}{u_0 t}\right)^2\right], & R_1 < \rho_0 < R, \\ 0, & \rho_0 > R. \end{cases} \tag{2.6}$$

One can also show that in this stage the radial velocity component v_ρ is considerably less than v_φ. On the basis of (2.6) we estimate the fraction of energy transferred from the dispersing cloud to the ions of the medium through the mechanism described above,

$$\frac{W_2}{W_0} = \frac{2\pi}{W_0}\int_0^\infty d\rho_0\rho_0 \frac{n_* m_2 v_\varphi^2}{2} = \frac{1}{2}\frac{m_2}{m_1}\left(\frac{\Omega_2 R_*}{u_0}\right)^2 \left\{\ln\left[1 + \left(\frac{R}{R_*}\right)^2\right] - \frac{R^2}{R_*^2 + R^2}\right\},$$

where $W_0 = m_1 N u_0^2/4$ is the initial energy of cloud ions.

Thus, the fraction of transferred energy is determined by the value of the parameter $\delta = (R_* \Omega_2/u_0)^2$. A condition for intense interaction of the dispersing cloud with the surrounding plasma is the fulfillment of the inequality $\delta \gtrsim 1$.

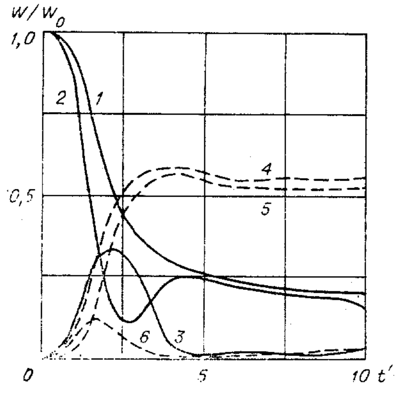

Fig. 1

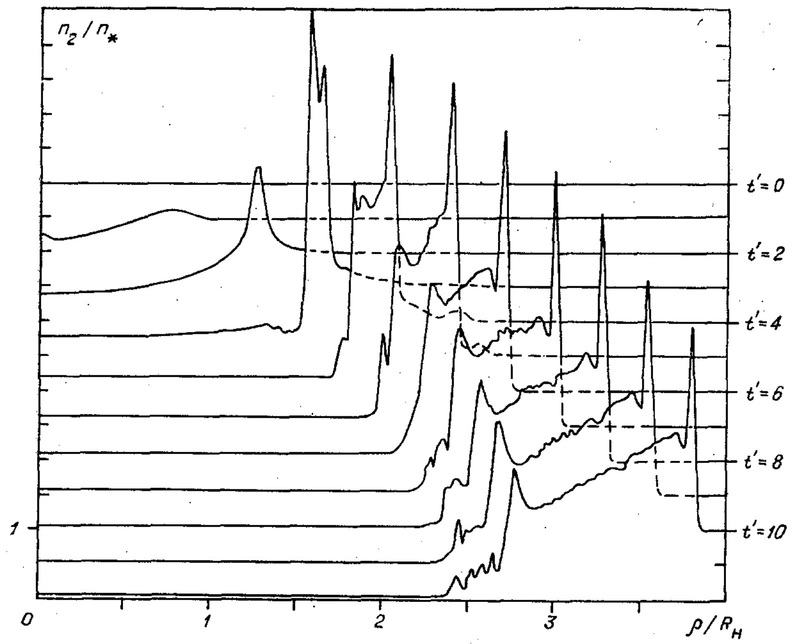

Fig. 2

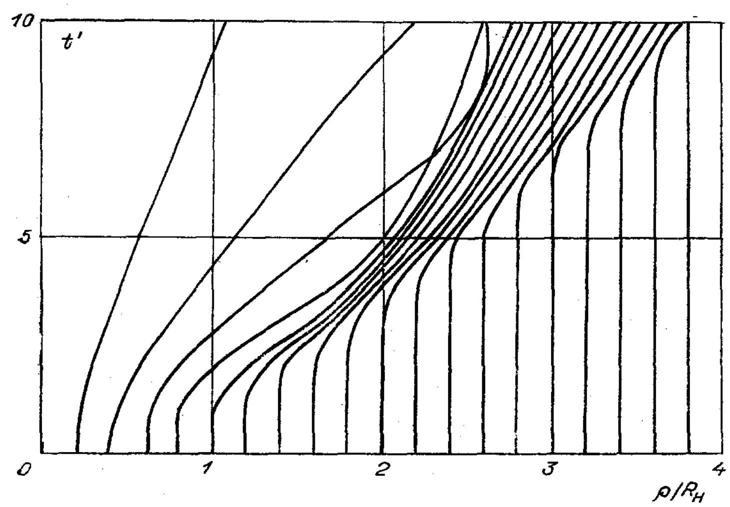

Fig. 3

3. Discussion of Results of the Numerical Calculation

Changing to the dimensionless variables and functions $t = t'/\Omega_2$, $\rho = \rho' u_0/\Omega_2$, $\nu = u_0\nu'$, $f_\alpha = (n_*/u_0^2)f'_\alpha$, $n_\alpha = n_* n'_\alpha$, and $E = (u_0 H_0 / c)E'$ in the initial equations and initial conditions (1.9)-(1.13), one can ascertain that the solution of the problem will depend on the following dimensionless parameters:

$$\delta = (R_* \Omega_2/u_0)^2, \quad \gamma = m_1/m_2, \quad M_A = u_0/v_A.$$

The solution of the kinetic equation for the ions was modeled using the method of particles in a cell [18]. The number of modeling particles was varied in the range of 10^3-10^4. The law of conservation of the total energy of the system,

$$\frac{d}{dt}\left(\sum_\alpha W_\alpha + 2\pi \int_0^\infty d\rho\rho \frac{H_z^2}{8\pi}\right) = 0,$$

which was satisfied with an accuracy of 1%, was used as a control on the accuracy of the solution.

The results of the calculations for the values of the parameters $\delta = \gamma = 1$ and $M_A = 10$ are presented in Figs. 1-6. The time dependences of the integral energy characteristics of the plasma are presented in Fig. 1 (in units of the total energy of the explosion):

$$W_\alpha^\rho = 2\pi \int_0^\infty d\rho\rho \int_{-\infty}^{+\infty} dv_\rho \int_{-\infty}^{+\infty} dv_\varphi \frac{m_\alpha v_\rho^2}{2} f_\alpha$$

(curves 2 and 5 for $\alpha = 1$ and 2, respectively),

$$W_\alpha^\varphi = 2\pi \int_0^\infty d\rho\rho \int_{-\infty}^{+\infty} dv_\rho \int_{-\infty}^{+\infty} dv_\varphi \frac{m_\alpha v_\varphi^2}{2} f_\alpha$$

(curves 3 and 6 for $\alpha = 1$ and 2),

$$W_\alpha = W_\alpha^\rho + W_\alpha^\varphi$$

(curves 1 and 4 for $\alpha = 1$ and 2).

It is seen that by the time $t' = 5$ the dispersing cloud loses ~75% of its total energy, with ~20% of the energy going into energy of the magnetic field. Subsequently the exchange of energy practically ceases and almost all the kinetic energy of the ions is concentrated in the radial motion.

The space–time pattern of the distribution of density n_2' of ions of the medium is illustrated by Fig. 2. At an early stage ($t' \leqslant 1$) in the vicinity of the point $\rho' = 1$, where the velocity v_φ is maximal (2.6), a region of increased plasma concentration is formed. The stable form of the concentration profile forms with the course

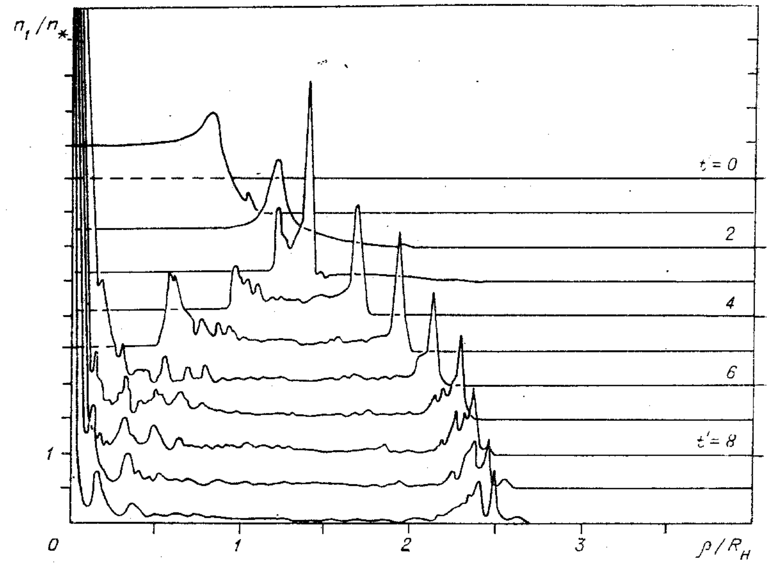

Fig. 4

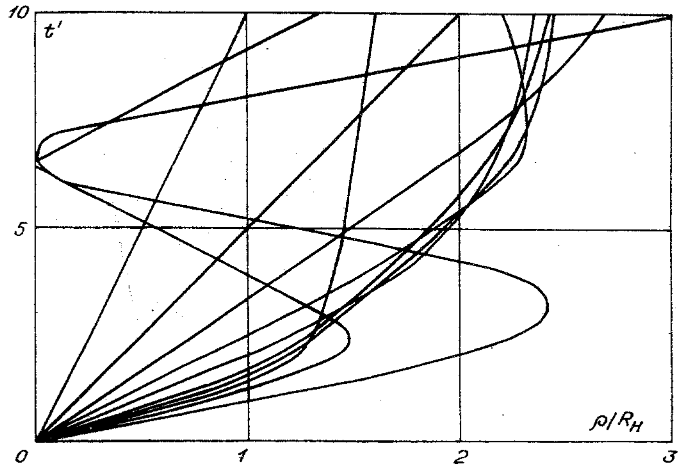

Fig. 5

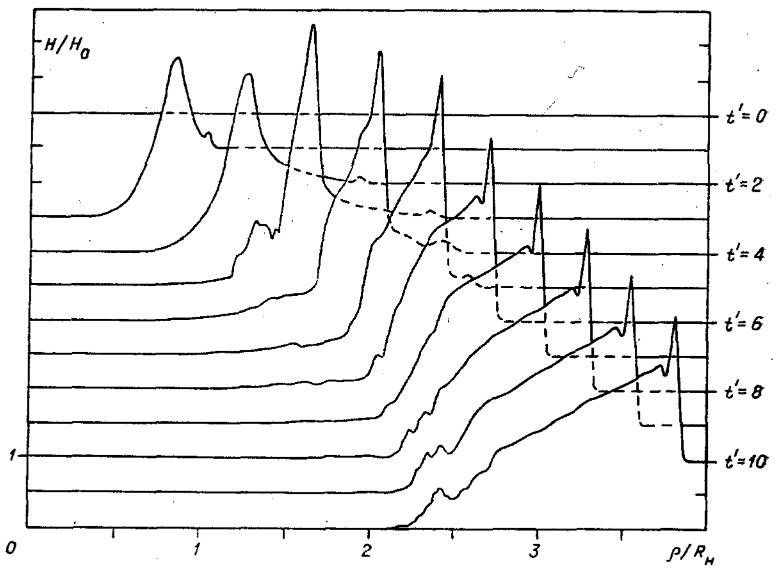

Fig. 6

of time ($t' > 5$). The sharp leading front moves in accordance with the law $\sim \sqrt{t'}$. The region of increased concentration expands monotonically. The developing plasma flow is additionally illustrated by Fig. 3, in which the trajectories of the model particles are presented.

The concentration distribution of cloud ions is shown in Fig. 4. If $\delta \leqslant 1$, then the maximum velocity of the ions of the medium for $t' \leqslant 1$ is less than the velocity of cloud ions at the leading front. Therefore, most of the cloud ions enter a region where the magnetic field is slightly disturbed. Moving in it along Larmor circles, by a time $t' \approx \pi$ they turn around and start to move toward the center, forming a region of high concentration there at times $t' > 2\pi$. Those of the ions which have velocities on the order of the velocity of the region of a compressed field form a condensation at the leading front of the cloud. The pattern of motion described above is clearly illustrated in Fig. 5, where the trajectories of particles modeling the plasma of the cloud are plotted.

The space–time pattern of variation of the magnetic field is presented in Fig. 6. The presence of a sharp magnetic field gradient at the leading front leads to a jump $\Delta\Phi'$ in the potential of the polarization electric field $E'_\rho = -\frac{1}{M_A^2 n'}\frac{\partial}{\partial\rho'}\frac{(H'_z)^2}{2}$. Since $H'_z \sim n'$ at the leading front, $\Delta\Phi'$ is estimated from the equation

$\Delta\Phi' \simeq M_A^{-2}(n'_2 - 1)$, where n'_2 is the concentration of the external plasma beyond the front. At $t' \geq 5$, $\Delta\Phi' \leqslant 0.02$, while the velocity of the front is $u'_f \simeq 0.3$. Consequently, $\frac{1}{2}m'_2(u'_f)^2 > \Delta\Phi'$ and there is no reflection of ions from the jump in potential at the front; the flow developing in the surrounding medium has an essentially one-velocity character. It should be noted that reflected ions should appear at larger values of M_A, according to [19]. This leads to a more complex flow pattern in the external plasma. In subsequent reports it is proposed to investigate the solution of the problem under consideration in this region of values of the parameter M_A.

LITERATURE CITED

1. H. W. Friedman and R. M. Patrick, "Momentum transfer in plasma flows at high Alfvén Mach numbers," Phys. Fluids, 14, No. 9, 1889 (1971).
2. P. T. Pumsby, J. W. M. Paul, and J. Beaulieu, "Shock wave studies using laser produced driving plasma," Bull. Am. Phys. Soc., 16, No. 10, 1191 (1971).
3. A. Y. Cheung, R. R. Goforth, and D. W. Koopman, "Magnetically induced collisionless coupling between counterstreaming laser-produced plasmas," Phys. Rev. Lett., 31, No. 7, 429 (1973).
4. Operation Argus [in Russian], Atomizdat, Moscow (1960).
5. Operation Starfish [in Russian], Atomizdat, Moscow (1964).
6. Yu. P. Raizer, "Retardation and energy conversions of a plasma expanding into an empty space in which there is a magnetic field," Zh. Prikl. Mekh. Tekh. Fiz., No. 6 (1963).
7. V. A. Pilipenko, "Boundary shape of a freely dispersing plasma in a magnetic field," Zh. Prikl. Mekh. Tekh. Fiz., No. 2 (1967).

8. J. W. Poukey, "Expansion of a plasma shell into a vacuum magnetic field," Phys. Fluids, 12, No. 7, 1452 (1969).
9. S. M. Bakhrakh, E. V. Gubkov, V. A. Zhmailo, and V. A. Terekhin, "Dispersion of a plasma cloud in a uniform magnetic field," Zh. Prikl. Mekh. Tekh. Fiz., No. 4 (1974).
10. S. A. Colgate, "The phenomenology of the mass motion of a high altitude nuclear explosion," J. Geophys. Res., 70, No. 13, 3161 (1965).
11. V. P. Korobeinikov, "Problems of the theory of a point explosion in a gas," Tr. Mat. Inst. Akad. Nauk SSSR, No. 119 (1973).
12. G. F. Chew, M. L. Goldberger, and F. E. Low, "The Boltzmann equation and the one-fluid hydromagnetic equations in the absence of particle collisions," Proc. R. Soc. London, A, 236, 112 (1956).
13. D. W. Forslund and J. P. Freidberg, "Theory of laminar collisionless shocks," Phys. Rev. Lett., 27, No. 18, 1189 (1971).
14. R. A. Chodura, "Hybrid-fluid-particle model of ion heating in high-Mach-number shock waves," Nucl. Fusion, 15, 55 (1975).
15. A. G. Sgro and C. W. Nielson, "Hybrid model studies of ion dynamics and magnetic field diffusion during pinch implosion," Phys. Fluids, 19, No. 1, 126 (1976).
16. Ya. B. Zel'dovich and Yu. P. Raizer, Physics of Shock Waves and High-Temperature Hydrodynamic Phenomena [in Russian], Nauka, Moscow (1966).
17. A. A. Galeev and R. Z. Sagdeev, "Nonlinear plasma theory," in: Problems of Plasma Theory [in Russian], Atomizdat, Moscow (1973).
18. R. L. Morse, "Modeling of a multidimensional plasma using the method of a particle in a cell," in: Computational Methods in Plasma Physics [Russian translation], Mir, Moscow (1974).
19. D. Biskamp and H. Welter, "Numerical studies of magnetosonic collisionless shock waves," Nucl. Fusion, 12, No. 6, 663 (1972).

CALCULATION OF THE EFFECTIVE THERMAL DIFFUSION COEFFICIENTS OF A NONLINEAR MULTIELEMENT PLASMA

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Knowledge of the transport properties of an equilibrium plasma consisting of a mixture of different chemical elements with a temperature from several thousand to several tens of thousands of degrees and with a concentration of charged particles of 10^{18}-10^{21} cm^{-3} is necessary in connection with the design of gas-phase nuclear reactors, powerful MHD generators, and thermal protection of spacecraft. In the indicated range of parameters the interaction energy of the charged particles of the plasma is of the order of the thermal energy – the plasma is nonideal. Experimental data of the transport properties of a nonideal plasma are very limited; calculation of rigorous theoretical expressions for the transport coefficients of an equilibrium plasma is possible on the condition of weak interparticle interaction; therefore, modeling approaches to the determination of the kinetic coefficients of a nonideal plasma acquire an important role.

The gasdynamics problems associated with the investigation of the flows of a nonideal plasma in the devices listed are usually solved in the approximation of local chemical equilibrium, and the plasma is assumed to be quasineutral. It is advisable in this approximation to switch from the diffusion equations of the components to the diffusion equations of the chemical elements if the number of components in the plasma is greater than the number of chemical elements. The mass flows of chemical elements and the "molecular" thermal flux are determined with the help of the introduction of effective transport coefficients in terms of the temperature gradient, the fractions of chemical elements, and the electric field in the plasma (the pressure is assumed to be constant). We emphasize that the calculation of the transport coefficients is correctly determined from the solution of a specific gasdynamics problem, which appreciably simplifies its formulation; the necessary transport properties of the plasma are compactly tabulated as a function of the pressure, temperature, and fractions of the chemical elements forming the plasma. The complete system of gasdynamics

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