DIFFUSION OF A STRONG MAGNETIC FIELD IN A DENSE PLASMA

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The diffusion of a strong magnetic field in a collisional plasma is important in many problems, such as plasma confinement by a magnetic field, magnetic dispersal of films, transforming during an electrical explosion into the plasma state, discharges arising on the surface of an insulator (magnetic pinch discharges) when a magnetic flux passes through it, etc.

The expansion into a vacuum of a thermally and electrically conducting gas formed from an electrical explosion of a plane conductor was studied in [1]. However, only the initial stage of the explosion was studied, since the electrical conductivity was assumed to decrease with increasing temperature.

In the present paper we consider the plane diffusion of a transverse magnetic field from a vacuum into a dense plasma in the collisional regime. A magnetic pinch discharge on the surface of an insulator can be considered as a special case of diffusion in a plasma of infinite density. We consider strong magnetic fields such that the plasma can be assumed to be completely ionized.

In the diffusion of a magnetic field in a dense plasma, three basic stages can be identified: 1) Radiation losses are small compared to Joule heating and the principal effects in limiting the skin-effect magnetic field are the electron thermal conductivity and thermoelectric effects (Nernst effect); 2) radiation losses begin to balance the Joule heating and a steady-state regime results; 3) radiation heats the inner layers of the plasma, thermal diffusion is determined by the radiative thermal conductivity, and the coefficients of magnetic diffusion and radiative thermal diffusivity become of the same order of magnitude.

We consider all quantities to be dependent on the coordinate X and time t, where the magnetic field H and electric field E are perpendicular to each other and to the X axis, and the characteristic times in the problem are large compared to gasdynamical times, such that the total pressure in the system can be equalized:

$$p+H^2/8\pi=H_0^2/8\pi; \tag{0.1}$$

where p is the thermal pressure and H_0 is the magnetic field strength on the boundary with the vacuum.

The equations for the magnetic and electric fields and the thermal balance equation of the plasma have the following forms (in Lagrangian variables)

$$\frac{\partial E}{\partial X}=-\frac{1}{c}\left(\frac{dH}{dt}-\frac{H}{\rho}\frac{d\rho}{dt}\right),\quad \frac{\partial H}{\partial X}=-\frac{4\pi}{c}j,\quad E=\frac{j}{G}-\frac{\beta_\Lambda}{e}\frac{\partial T}{\partial X}. \tag{0.2}$$
$$\rho\frac{d\varepsilon}{dt}-\frac{p}{\rho}\frac{d\rho}{dt}=-\frac{\partial Q}{\partial X}+jE-J,\quad Q=-\chi\frac{\partial T}{\partial X}+\frac{\beta_\Lambda T}{e}j,$$

where ρ is the density of the plasma, ε is the internal energy, σ, χ, β_Λ are the transverse electrical conductivity, thermal conductivity, and thermoelectric coefficient, respectively, J is the volume power of the radiation loss, and Q is the heat flux density. We assume that at the initial time the magnetic field is zero inside the plasma and that the plasma is initially homogeneous.

1. Diffusion of a Magnetic Field in a Hydrogen Plasma at Small Values of the Time.

First we consider the diffusion of a magnetic field in a plasma for small times when radiation is unimportant and the electron transport coefficients play the principal role. In this case, the magnetic diffusion and thermal diffusivity coefficients become of the same order of magnitude when the degree of magnetization of the electrons is such that $\omega_e\tau_e \sim 1$. We choose as units of temperature and electron density N the following quantities

$$[T]=\left(\frac{H_0}{\sqrt{8\pi}}c\lambda_c ze^3\sqrt{m}\right)^{2/5}. \tag{1.1}$$

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$$[N] = \frac{H_0^2}{8\pi}\Big/[T], \tag{1.2}$$

where z is the ion charge and λ_c is the Coulomb logarithm. Using the variable

$$\xi = \frac{e^{1.1}m^{0.1}(\lambda_c z)^{0.2}}{(H_0^2/8\pi)^{0.65}} \frac{\int N dX}{\sqrt{t}} \tag{1.3}$$

and introducing the dimensionless functions

$$T = [T]\,\Theta(\xi),\ N = [N]\,n(\xi),\ H = H_0 h(\xi),$$
$$E = \frac{e^{0.1}m^{0.1}(\lambda_c z)^{0.2}(H_0^2/8\pi)^{0.35}}{c^{0.3}\sqrt{t}}\,\varepsilon(\xi), \tag{1.4}$$
$$Q = \frac{e^{0.1}m^{0.1}c^{0.7}(\lambda_c z)^{0.2}(H_0^2/8\pi)^{0.85}}{\sqrt{t}}\,q(\xi),\quad X = \frac{e^{0.1}m^{0.1}c^{0.7}(\lambda_c z)^{0.2}}{(H_0^2/8\pi)^{0.15}}\,x(\xi),$$

the system of equations (0.1) and (0.2) can be rewritten in the form

$$n\Theta(1 + 1/z) + h^2 = 1,\quad \frac{d\varepsilon}{d\xi} = \frac{\sqrt{2\pi}\xi}{n}\left(\frac{dh}{d\xi} - \frac{h}{n}\frac{dn}{d\xi}\right),$$
$$\frac{dh}{d\xi} = -\frac{3\Theta^{3/2}}{4\alpha n}\left(\sum + \beta n\frac{d\Theta}{d\xi}\right),\quad q = -\frac{\beta n\Theta}{\sqrt{2\pi}}\frac{dh}{d\xi} - \frac{3\gamma n}{4\sqrt{2\pi}}\Theta^{5/2}\frac{d\Theta}{d\xi}, \tag{1.5}$$
$$\xi\left[\frac{5}{4}(1 + 1/z)\,n\frac{d\Theta}{d\xi} + h\frac{dh}{d\xi}\right] = n\frac{dq}{d\xi} + \frac{n\varepsilon}{\sqrt{2\pi}}\frac{dh}{d\xi},\ \frac{dx}{d\xi} = 1/n,$$

where α, β, γ depend on the degree of magnetization

$$y \equiv \omega_e\tau_e = \frac{3h}{2n}\Theta^{3/2} \tag{1.6}$$

and are given by the approximate relations [2]

$$\alpha = 1 - \frac{\alpha_1^1 y^2 + \alpha_0^1}{\Delta},\quad \beta = \frac{y(\beta_1'' y^2 + \beta_0'')}{\Delta},\quad \Delta = y^4 + \delta_1 y^2 + \delta_0. \tag{1.7}$$

where the notation for the coefficients in (1.7) follows that in [2].

The boundary conditions for the system (1.5) are:

$$h(0) = 1,\ h(\infty) = 0,\ n(\infty) = n_\infty,\ q(0) = q(\infty) = 0. \tag{1.8}$$

Using (1.5) and the boundary conditions (1.8) we can obtain expansions for $n(\xi)$, $\Theta(\xi)$, $q(\xi)$ as $\xi \to 0$ by taking into account that when $\xi \to 0$, we have $n \to 0$, $y \to \infty$, and by using analytical expressions for the kinetic coefficients for strong magnetization [3]:

$$n \sim \xi^k,\quad \Theta \sim \xi^{4k-2},\ q \sim \xi^{5k-2},\quad k = \frac{5z^2 + 4(\sqrt{2}-1)z + 4}{9z^2 + (8\sqrt{2}-7)z + 10}. \tag{1.9}$$

This stage of diffusion plays a role only in a hydrogen plasma (z = 1). For z > 1 there are large radiative losses, and the transition to the steady-state regime occurs early, when condition (0.1) is still not satisfied and the inertia of the material cannot be neglected.

We consider the solution of (1.5) for z = 1. In this case the expansion (1.9) gives

$$n \sim \xi^{0.457},\ \Theta \sim \xi^{-0.172},\ q \sim \xi^{0.286},$$

and the temperature at the boundary with the vacuum goes to infinity. The numerical solution of (1.5) for $n_\infty = \infty$ is shown in Fig. 1. The magnitude of the electric field on the boundary with the vacuum is shown as a function of n_∞ in Fig. 2. The electric field ε_0 approaches the constant value $\varepsilon_0 \simeq 2.04$ at large n_∞, and at small n_∞ it becomes proportional to $n_\infty^{3/4}$, as would be expected.

However, when

$$n_\infty \ll \left(\frac{m}{M}\right)^{0.1}\Big/z^{0.4}$$

(M is the ion mass) the ion thermal conductivity is more important than the electron thermal conductivity and the plasma can be considered as isothermal with temperature $\Theta = 1/[n_\infty(1 + 1/z)]$. We transform to the new variables

$$\xi' = \Theta^{7/4}(1 + 1/z)\xi,\ \varepsilon' = \Theta^{3/4}\varepsilon,\ n' = \Theta(1 + 1/z)n,\ x' = \Theta^{3/4}x,$$

and the system of equations (1.5) takes the form

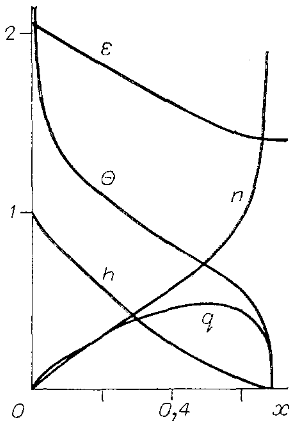


Fig. 1

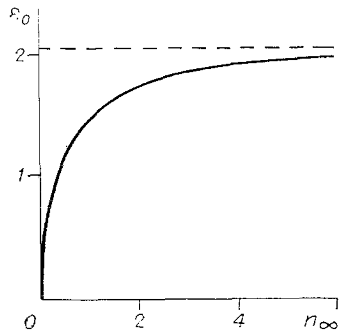


Fig. 2

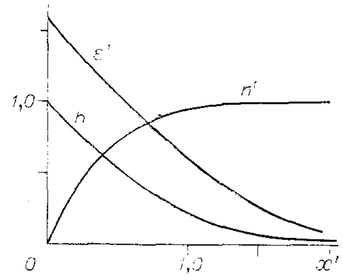


Fig. 3

$$\frac{dh}{d\xi'} = -\frac{3}{4}\frac{\varepsilon'}{(1-h^2)}, \quad \frac{d\varepsilon'}{d\xi'} = -\frac{3\sqrt{2\pi}}{4}\xi'\frac{(1+h^2)}{(1-h^2)^3}\varepsilon', \quad \frac{dx'}{d\xi'} = 1/(1-h^2). \tag{1.10}$$

The numerical solution of (1.10) with the boundary conditions $h(0) = 1$, $h(\infty) = \varepsilon(\infty) = 0$ is shown in Fig. 3. Note that $\varepsilon(0) = 1.58$ closely agrees with $\Theta_\infty^{3/4}\varepsilon(0)$ of Fig. 2 in the limit $n_\infty \to 0$, $\Theta_\infty \to \infty$: $(\Theta_\infty^{3/4}\varepsilon(0))_{n_\infty\to 0} \simeq 1.51$. This means that the isothermal approximation is valid in a low-density plasma even when only the electron transport coefficients are taken into account.

2. Steady-State Stage. As the thickness of the discharge region increases, the rate of Joule heating per unit volume decreases and the radiation losses, which are determined by the temperature (1.1) and density (1.2), remain the same. Therefore, as time increases the discharge goes into a steady-state phase when the Joule heating is balanced by radiation losses. It is easy to see that in this stage the thickness of the discharge region is small compared to the radiation path length and we can neglect surface effects in treating the radiation.

For strong fields in a hydrogen plasma, bremsstrahlung plays the principal role, and its power is given by

$$J_B = \frac{32}{3}\sqrt{\frac{2T}{\pi m}}\frac{zN^2e^6}{mc^3\hbar}.$$

Dimensionless quantities for this stage of the discharge are conveniently chosen according to (1.1)-(1.4), where in place of the time t we use the variable

$$\tau = \frac{H_0^2}{8\pi J_B([T],[N])}. \tag{2.1}$$

In the steady-state case, the second equation of (1.5) becomes

$$\varepsilon = \text{const}, \tag{2.2}$$

while the last equation can be written in the form

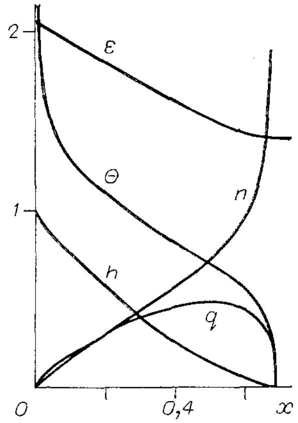
$$\frac{dq}{d\xi} = -n\sqrt{\Theta} - \frac{\varepsilon}{\sqrt{2\pi}}\frac{dh}{d\xi}, \tag{2.3}$$

and the other equations remain unchanged. The solution of (1.5) is shown in Fig. 4 for this case. The electric field has the value $\varepsilon = 1.16$.

For plasmas with atomic numbers $z_0 > 1$, recombination and line-emission are important effects. For the temperatures and densities (1.1), (1.2) and megagauss magnetic fields, the volume radiation of the plasma can be approximated by [4]

$$J_R(T, N) = R\frac{N^2}{\sqrt{Tm}}\frac{z_0^4}{z}\frac{e^{10}}{\hbar^3c^3}, \tag{2.4}$$

where R is a dimensionless constant. Then we use in (1.5) analogous with (2.1) for the dimensionless quantities (1.1)-(1.4) the following quantity in place of t, in analogy with (2.1):



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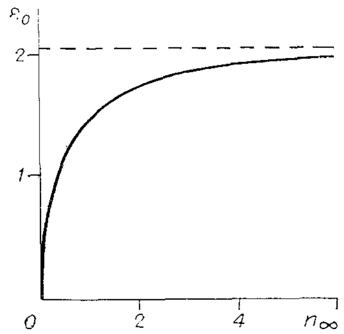

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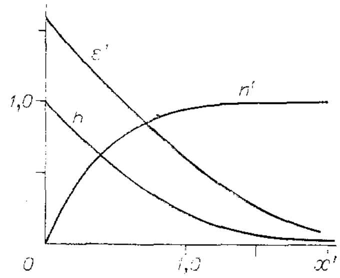

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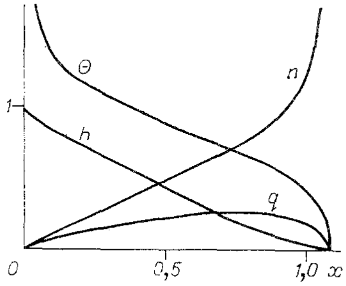

Fig. 4

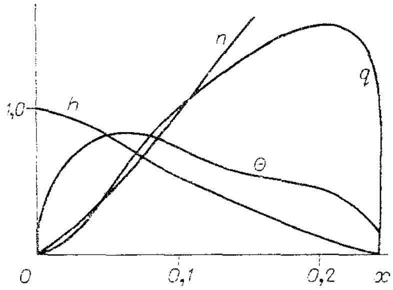

Fig. 5

$$\tau = \frac{H_0^2}{8\pi J_R([T], [N])} \tag{2.5}$$

and Eq. (1.5) remains unchanged except for the second which becomes (2.2) and the last, which will have the form

$$\frac{dq}{d\xi} = -n/\sqrt{\Theta} - \frac{\varepsilon}{\sqrt{2\pi}} \cdot \frac{dh}{d\xi}. \tag{2.6}$$

The solution of (1.5) for this case with the coefficients (1.7) is shown in Fig. 5 for $z = \infty$. The electric field in this case is $\varepsilon = 5.58$. A discharge of the type of Fig. 5 can also be interpreted as a magnetic pinch discharge at the surface of an insulator. We note that at large z the temperature on the plasma–vacuum boundary goes to zero (unlike the case z = 1). This behavior follows from expansions (1.9) and is due to the more significant (compared to the case z = 1) effect of thermoelectric heat fluxes on the thermal balance near the boundary $\xi = 0$.

We present typical numerical values of the quantities for discharges in hydrogen and organic glass ($H_8C_5O_2$) in magnetic fields in the megagauss range.

For hydrogen in megagauss fields the Coulomb logarithm is $\lambda_C = 7.5$, and

$$[T] = 74\ \text{eV} \cdot H_0^{0,4}(\text{MG}),\ [N] = 3{,}3 \cdot 10^{20}\ \text{cm}^{-3} \cdot H_0^{1,6}(\text{MG}).$$

In Fig. 1 the unit of distance is

$$[X] = 0.19\ \text{cm}\ \sqrt{t(\mu\text{sec})}/H_0^{0,3}(\text{MG}),$$

and the unit of electric field is

$$[E] = 1.9\frac{\text{kV}}{\text{cm}} \cdot H_0^{0,7}(\text{MG})/\sqrt{t(\mu\text{sec})}.$$

In Fig. 4 the unit of distance is $[X] = 0.095$ cm/H_0 (MG) and the unit of electric field is

$$[E] = 3.9\frac{\text{kV}}{\text{cm}} \cdot H_0^{1,4}(\text{MG}).$$

If we assume that the start of the transition from the regime of Fig. 1 to that of Fig. 4 is determined by the equality of the electric field E_∞ of Fig. 4 and E of Fig. 1, and the end of the transition is determined by the equality of the field E_0 of Fig. 1 and E of Fig. 4, then the characteristic initial and final times of the transition will be given by

$$t_i = 0.35\ \mu\text{sec}/H_0^{1,4}(\text{MG}),\quad t_f = 0.75\ \mu\text{sec}/H_0^{1,4}(\text{MG}).$$

The time corresponding to the start of the regime of Fig. 1, when the inertia of the material can be neglected and condition (0.1) begins to be satisfied, is

$$t \sim 6 \cdot 10^{-3}\ \mu\text{sec}/H_0\ (\text{MG}).$$

For organic glass in megagauss fields, the average ionic charge is given by $z = \Sigma n_i z_i^2/n$, where n_i is the density of ions with charge z_i (cf [5]). The Coulomb logarithm $\lambda_C \simeq 5.5$, and the units of temperature and density in Fig. 5 are

$$[T] = 120\ \text{eV} \cdot H_0^{0,4}(\text{MG}),\quad [N] = 2{,}1 \cdot 10^{20}\ \text{cm}^{-3} \cdot H_0^{1,6}(\text{MG}).$$

The quantity z_0^4/z, which is to be substituted in (2.4), is given by [4] $z_0^4/z = \Sigma n_i z_{0i}^4/n$, and the numerical constant R can be obtained using radiation tables for low-density plasmas [5]. For $H_8C_5O_2$ we have $z_0^4/z \simeq 350$, $R \simeq 50$ in the temperature region of interest to us. With these values, the unit of distance in Fig. 5 is

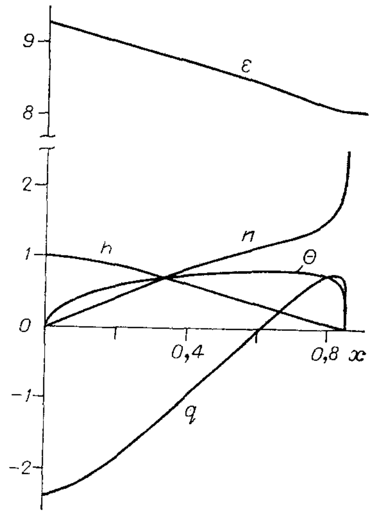

Fig. 6

$$[X] = 0.80 \cdot 10^{-2}\,\text{cm}\,/H_0^{0,8}\,(\text{MG}),$$

and the unit of electric field is

$$[E] = 74\,\frac{\text{kV}}{\text{cm}} \cdot H_0^{1,2}\,(\text{MG}).$$

The time corresponding to the start of the regime of Fig. 5, when the Joule heating begins to be balanced by radiation losses and the inertia of the material can be ignored, is given by:

$$t \sim 0{,}2 \cdot 10^{-3}\,\mu\text{sec}/H_0(\text{MG}).$$

3. Magnetic Diffusion Accompanied by Radiative Thermal Conduction. For dense plasmas or insulators, radiation from the steady-state regions of the discharge gradually heats up the inner layers, increasing their electrical conductivity, and the magnetic field begins to diffuse inward, heating the plasma, and heat transport to the neighboring layers occurs. This is how the next stage of magnetic field diffusion occurs; we consider the diffusion of a strong field in organic glass.

We consider a gradual dependence of the equation of state, radiation path length λ, and magnetic diffusion coefficient $\varkappa$ on temperature:

$$p/\rho = 0{,}17\,T^{1,19}/\rho^{0,06},\quad \lambda = 2 \cdot 10^{-9} T^{2,14}/\rho^{1,86},\quad \varkappa = 0.17/(T^{0,86}\rho^{0,14}),$$

where the adiabatic exponent is $p/\varepsilon\rho + 1 = 4/3$ and the system of units is g, cm, μsec, and eV for temperature. We choose units of temperature [T] and density [ρ] so that the thermal diffusivity and magnetic diffusion coefficients are of the same order

$$\sigma_{SB}\,[T]^4 \lambda\,([T],[\rho])/(H_0^2/8\pi) = \varkappa\,([T],[\rho])$$

(where $\sigma_{SB} = 1.03 \cdot 10^{-6}$ is the Stefan–Boltzmann constant) and the thermal pressure is the same order as the magnetic pressure

$$p\,([T],[\rho]) = H_0^2/8\pi.$$

Then

$$[T] = 17\,\text{eV} \cdot H_0^{0,62}\,(\text{MG}),\quad [\rho] = 5.7 \cdot 10^{-3}\,\frac{\text{g}}{\text{cm}^3} \cdot H_0^{1,35}\,(\text{MG}).$$

Using the self-modeling variable

$$\xi = \frac{1000 \int \rho dX \left(\frac{\text{g}}{\text{cm}^2}\right)}{\sqrt{t\,(\mu\text{sec})}\,H_0^{0,99}\,(\text{MG})}$$

and introducing the dimensionless functions

$$T = [T]\,\Theta(\xi),\quad \rho = [\rho]\,n(\xi),\quad H = H_0 h(\xi),\quad E = 1.8\,\frac{\text{kV}}{\text{cm}} \cdot \frac{H_0^{0,64}\,(\text{MG})}{\sqrt{t\,(\mu\text{sec})}}\,\varepsilon(\xi),$$

$$X = 0.18\,\text{cm} \cdot \frac{\sqrt{t\,(\mu\text{sec})}}{H_0^{0,36}\,(\text{MG})}\,x(\xi),\quad Q = 1.8 \cdot 10^4\,\frac{\text{W}}{\text{cm}^2} \cdot \frac{H_0^{1.64}\,(\text{MG})}{\sqrt{t\,(\mu\text{sec})}}\,q(\xi),$$

the system of equations (0.1) and (0.2) can be written in the form

$$\Theta^{1.19}/n^{0.94}+h^2=1,\ \frac{d\varepsilon}{d\xi}=\frac{\sqrt{2\pi}}{n}\xi\left(\frac{dh}{d\xi}-\frac{h}{n}\frac{dn}{d\xi}\right),$$
$$\frac{dh}{d\xi}=-\frac{1}{2\sqrt{2\pi}}\frac{\Theta^{0.86}\varepsilon}{n^{0.86}},\quad q=-\frac{16}{3}\frac{\Theta^{5.14}}{n^{0.86}}\frac{d\Theta}{d\xi}, \tag{3.1}$$
$$\frac{dq}{d\xi}=\frac{\varepsilon^2}{4\pi}\frac{\Theta^{0.86}}{n^{0.86}}+0.59\xi\left(3\frac{\Theta^{0.19}}{n^{0.06}}\frac{d\Theta}{d\xi}-\frac{\Theta^{1.19}}{n^{1.06}}\frac{dn}{d\xi}\right),\quad \frac{dx}{d\xi}=\frac{1}{n}.$$

From (3.1) and the boundary conditions $h(0) = 1$, $\varepsilon(0) = \text{const}$, $q(0) = \text{const}$, we have the following expansions for $\Theta(\xi)$ and $n(\xi)$:

$$\Theta(\xi)|_{\xi\to 0}\sim\xi^{0.23},\ n(\xi)|_{\xi\to 0}\sim\xi^{0.51}.$$

The solution of the system (3.1) with the boundary condition $n(\infty) = \infty$ is shown in Fig. 6. The characteristic times corresponding to the start and end (t_i and t_f) of the transition from the regime of Fig. 5 to that of Fig. 6 can be estimated by equating the electric field of Fig. 5 to the fields $E(\infty)$ and $E(0)$ of Fig. 6:

$$t_i=0.0012\ \mu\text{sec}/H_0^{1.12}(\text{MG}),\ t_f=0.0015\ \mu\text{sec}/H_0^{1.12}(\text{MG}).$$

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LOSS OF EQUILIBRIUM AND THE QUASISTATIONARY STATE IN AN EXPANDING RECOMBINING PLASMA

G. M. Zhinzhikov, G. A. Luk'yanov, and N. O. Pavlova

UDC 533.9

Many problems of modern gasdynamics and technical physics are concerned with thermodynamic nonequilibrium states of a medium and conditions with thermodynamic nonequilibrium states of a medium and conditions for obtaining the nonequilibrium state. A problem of this type, whose importance comes from its application to creation of effective plasma lasers [1], is the relaxation of a low-temperature plasma during an adiabatic expansion. The conditions for the loss of equilibrium are well-known for some typical situations in a plasma [2, 3]. As a rule these are situations when the cause of the loss of equilibrium is the steady-state effect of perturbing factors on the parameters of the problem. In the present paper we consider the loss of equilibrium in a nonsteady plasma. Criteria are obtained for the loss of ionization equilibrium, the equilibrium distribution of levels, and thermal equilibrium for the expansion of a plasma which is initially in equilibrium. We also study the closely related conditions for a quasistationary occupation of the excited states in a recombining plasma.

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