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ON THE POSSIBILITY OF ULTRAHIGH MAGNETIC FIELDS PRODUCTION IN IMPLODING PLASMA: MAGNETIC FLUX RETENTION

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A b s t r a c t

The problem of the axial magnetic flux retention in a low-β plasma imploded by the cylindrical liner has been studied. It is shown that the rate of magnetic flux losses in such a system becomes anomalously high due to the formation of a high-density plasma sheath at the liner wall. Analytical solution of the plasma transport equations based on the self-similar nature of the sheath profile evolution in time has been obtained. The enhanced magnetic flux losses result in rather hard requirements to the liner velocity needed for the production of ultrahigh magnetic fields in imploding plasma.

1. Introduction

Recently the new scheme for production of ultrahigh (~100 Mgs) magnetic fields has been proposed [1,2], based on the implosion of the plasma with a frozen magnetic field by the liner. In this method the field amplification is provided by currents induced in the plasma, and it results in some advantages in comparison with the traditional way of megagauss fields production by magnetic flux compression with a conducting liner.

The principal problem is now as follows: what implosion velocity it needs for "frozen in field" condition to be valid? It has been considered in [1,2] that it needs the magnetic Reynolds number S_m to be much greater than unity:

$$S_m = R_L(t)\, U_L(t) / D_m(t) \gg 1 \qquad (1)$$

Here R_L is the radius of the liner, $U_L = /\dot{R}_L/$ - implosion velocity, $D_m = c^2/4\pi\sigma$ - magnetic visco sity coefficient of the plasma, σ -plasma conductivity. But as it was noted in [3], the rate of the magnetic flux losses in such a system becomes anomalously high in comparison with the simple estimation that results in criteria (1). The reason is that the growing magnetic field in the main volume of system pushes plasma to the liner, so the high-density plasma sheath is formed near the wall. The corresponding plasma flow leads to the convective transport of the magnetic field to the sheath where the field is lost due to resistive magnetic diffusion. As a result the magnetic field remains frozen into the plasma if the following condition, rather more hard than criteria (1), is satisfide:

$$S_{eff.} = R_L U_L / D_{eff} \gg 1 \qquad (2)$$

where the effective magnetic visco sity of the plasma, $D_{eff.}$, greatly exceeds its classical value $D_m = c^2/4\pi\sigma$.

These features of the magnetic field diffusion in imploding plasma have been confirmed by the numerical simulation results presented in [4]. As to the analytical solutions of a plasma transport equations obtained in [4], they are irrelevant for the problem under discussion. The reason is that these self-similar solutions correspond to the homogeneous compression while the effect mentioned above results from the plasma redistribution during its implosion.

In this paper it is shown that the problem of the magnetic field evolution in a low-β plasma imploded by the liner posseses the analytical solution for an arbitary sub-alfvenic implosion low $R_L(t)$. It is based on the fact that the thickness of a high-density near-wall plasma sheath, where all the gradients are significant, remains small as compared to the plasma radius. So the evolution in time of the parameters of a homogeneous plasma may be obtained from the simple equations that describe the balance of energy, magnetic flux and number of particles between the sheath and the main plasma volume. At first the sheath structure is determined which is described by the self-similar solution of the plasma transport equations in the magnetic field. Then the equation for the magnetic field in imploding plasma is obtained and the possibility of the ultrahigh fields production in such a scheme is discussed.

2. Basic equations and near-wall sheath structure

Let us consider the cylindrically symmetric plasma implosion with the axial magnetic field. At the initial moment (before compression) inside the liner with radius R_o we have a homogeneous plasma with the density n_o, temperature T_o and seed magnetic field H_o. As to the sense of the problem it is supposed that the plasma pressure is small as compared with the magnetic one, so that $\beta_o \equiv 8\pi n_o T_o / H_o^2 \ll 1$. In the compressed plasma the magnetic field greatly exceeds its value at the liner (for the non-conducting liner the magnetic field a. the liner remains to be equal to H_o). At the same time for the sub-alfvenic implosion (only this one is of practical interest here) the process is almougt adiabatic and the total pressure remains to be approximately homogeneous:

$$\frac{\partial}{\partial r}\left(nT + H^2/8\pi\right) \approx 0 \qquad (3)$$

Under the condition $\beta_o \ll 1$ it means that although in the main volume plasma pressure is small ($\beta_i = 8\pi n_i T_i / H_i^2 \ll 1$), the high-pressure plasma sheath bears near the liner wall where $nT \sim H_i^2/8\pi$ and $n \gg n_i$ (here and later on the quantities H_i, T_i and n_i mark the parameters of the homogeneous plasma in the centre of the system). The qualitative sketch of the magnetic field and plasma pressure profiles is shown in Fig. 1.

The sheath thickness Δ is much less than the plasma radius ($\Delta \ll R$, see Appendix, 1), so the problem of the sheath structure description may be considered to be plane. Let us use the reference frame connected with the liner, its wall is the plane $x = 0$ and the plasma occupies the region $x > 0$. Then the plasma transport equations [5] take the form:

$$nT + H^2/8\pi = H_i^2/8\pi \qquad (4)$$

$$\frac{\partial n}{\partial t} + \frac{\partial}{\partial x}(nv) = 0 \qquad (5)$$

$$\frac{\partial H}{\partial t} = \frac{\partial}{\partial x}\left(\frac{c^2}{4\pi\sigma}\frac{\partial H}{\partial x} + \frac{c}{e}\beta_\wedge\frac{\partial T}{\partial x} - vH\right) = -\partial q_H/\partial x \qquad (6)$$

$$\frac{\partial}{\partial t}\left(\frac{3}{2}nT + H^2/8\pi\right) = \frac{\partial}{\partial x}\left\{\varkappa_\perp\frac{\partial T}{\partial x} + \frac{cT}{4\pi e}\beta_\wedge\frac{\partial H}{\partial x} - \frac{5}{2}nTv - \frac{H}{4\pi}q_H\right\} = -\frac{\partial q_W}{\partial x} \qquad (7)$$

(here the notations of [5] are used).

The boundary conditions for the set (4-7) require that at the liner wall ($x = 0$) the magnetic field is equal to the seed one: $H(0,t) = H_o$, and the plasma velocity $v(0,t) = 0$. We shall consider also the heat capacity of the liner to be large enough, so that its temperature remains constant: $T(0,t) = T_L$. Without the sheath (i.e. at $x \gg \Delta$) the plasma and the magnetic field are homogeneous, so we are interested in such the solution of set (4-7) that $H = H_i$, $n = n_i$ and $T = T_i$ when $x \to +\infty$.

The small thickness of the sheath mentioned above leads to the fact that the magnetic field and energy fluxes q_H and q_W draged by the plasma flow to the wall remain approximately constant inside the sheath (see Appendix, 2):

$$\Delta q_H = -\int_\Delta \frac{\partial H}{\partial t}dx \ll q_H, \quad \Delta q_W = -\int_\Delta \frac{\partial}{\partial t}\left(\frac{3}{2}nT + H^2/8\pi\right)dx \ll q_W$$

So we can use the following constant flux relations insted of eq's (6) and (7):

$$q_H = vH - \frac{c^2}{4\pi\sigma}\frac{\partial H}{\partial x} - \frac{c}{e}\beta_\wedge\frac{\partial T}{\partial x} = v_i H_i \qquad (6')$$

$$q_W = -\varkappa_\perp\frac{\partial T}{\partial x} - \frac{cT}{4\pi e}\beta_\wedge\frac{\partial H}{\partial x} + \frac{5}{2}nTv + q_H\frac{H}{4\pi} = v_i\frac{H_i^2}{4\pi} \qquad (7')$$

where $v_i = v(+\infty)$ is the velocity of the homogeneous plasma flow to the wall. It is convenient now to use the following dimensionless variables: $h = H/H_i$, $\rho = n/n_i$, $u = v/v_i$,

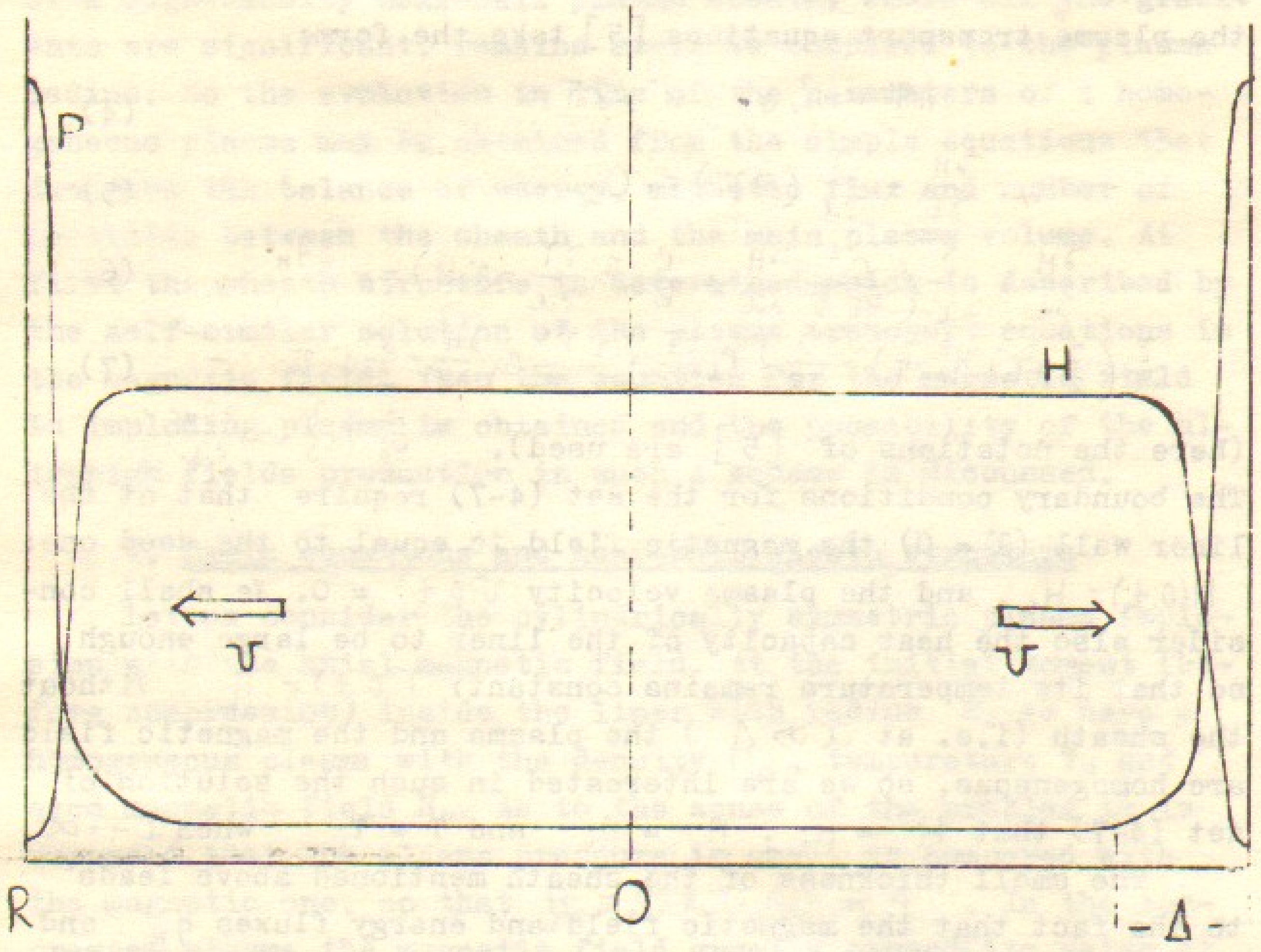

Fig. 1. The plasma pressure and the magnetic field profiles in imploding plasma

$\theta = T/T_H$ with $T_H = H_i^2/8\pi n_i$, and $\xi = x|v_i|/\mathcal{D}_H$ ($\mathcal{D}_H = c^2/4\pi\sigma(T_H)$ -magnetic viscosity of a plasma with the temperature $T = T_H$). Then eq's (4), (6') and (7') take the form:

$$\rho\theta + h^2 = 1 \tag{8}$$

$$uh + \theta^{-3/2}\frac{dh}{d\xi} + \alpha\frac{d\theta}{d\xi} = 1 \tag{9}$$

$$\varkappa\frac{d\theta}{d\xi} + \alpha\theta\frac{dh}{d\xi} + \frac{5}{4}\rho\theta u + h = 1 \tag{10}$$

It is important for the problem under consideration that although in the main part of the volume the plasma is strongly magnetized (the parameter $\omega_H\tau \gg 1$, where ω_H is the cyclotron frequency and τ - the scattering time for the particles), the magnetic field reduction and the plasma density growth in the sheath make the plasma unmagnetized at the liner wall, so that $\omega_H\tau \ll 1$ at $x = 0$. At the same time it is well known [5] that the plasma transport coefficients have a quite different dependence on plasma parameters and magnetic field in the magnetized and unmagnetized plasma. Therefore we use here rather simple expressions for $\beta_\wedge$ and $\varkappa_\perp$ that give correct results for these two limiting cases ($\omega_H\tau \gg 1$ and $\omega_H\tau \ll 1$). For that purpose let us consider separately three domains of a plasma parameters: I - the plasma with magnetized electrons and ions, where $(\omega_H\tau)_{e,i} > 1$; III - unmagnetized plasma with $(\omega_H\tau)_{e,i} < 1$; II - the intermediate domain with magnetized electrons but unmagnetized ions. Since $(\omega_H\tau)_i = \mu^{1/2}(\omega_H\tau)_e$, where $\mu \equiv m_e/m_i \ll 1$ is the mass ratio for electron and ion, the following inequalities are valid in II: $\mu^{1/2} < (\omega_H\tau)_i < 1$. After that the dimensionless coefficients α and $\varkappa$ in eq's (9) and (10) may be written in the following form [5]:

$$\varkappa = \begin{cases} \mu^{-1/2}\rho^2/h^2\theta^{1/2}, \ (\mathrm{I}), \ h\theta^{3/2}/\rho > \mu^{-1/2}\delta^{-1} \\ \delta\rho\theta/h, \ (\mathrm{II}), \ \delta^{-1} < \dfrac{h\theta^{3/2}}{\rho} < \mu^{-1/2}\delta^{-1} \\ \delta^2\theta^{5/2}, \ (\mathrm{III}), \ h\theta^{3/2}/\rho < \delta^{-1} \end{cases} \qquad \alpha = \begin{cases} \rho/h\theta^{3/2}, \ (\mathrm{I}, \mathrm{II}) \\ \delta^2 h\theta^{3/2}/\rho \ (\mathrm{III}) \end{cases} \tag{11}$$

Here $\delta \equiv (\omega_H\tau)_e(n_i, H_i, T_H) \gg 1$ - the large parameter, which (as it will be shown later) finally leads to the anomalous magnetic flux losses. Let us formulate now the boundary conditions in dimensionless variables notations. Since during the implosion the magnetic field in a compressed plasma dominates its value at the liner wall, it is possible to consider $h\ (\xi = 0) \simeq 0$. Further, the temperature of a homogeneous plasma $T_i \ll T_H = H_i^2/8\pi n_i$ (it follows from $\beta_i = 8\pi n_i T_i/H_i^2 \ll 1$), so $\theta\ (\xi \to \infty) \simeq 0$. At the same time the value of T_H rises during the compression ($T_H \propto H_i$) and finally significantly dominates the temperature of the liner wall T_L. This allows to put $\theta\ (\xi = 0) \simeq 0$. As a result the structure of a near-wall sheath may be described by the universal set of eq's (8)-(10) and boundary conditions

$$\rho(+\infty) = h(+\infty) = u(+\infty) = 1$$
$$u(0) = h(0) = \theta(0) = \theta(+\infty) = 0 \tag{12}$$

It means that for an any moment of the implosion process the profiles of the magnetic field H, density n and plasma temperature T in the sheath differ only by the scale, so the evolution of the sheath exhibits self-similar nature.

Since the plasma velocity u enters into eq's (9) and (10), they have to be solved, in general, together with the continuity equation (5). But in this case the situation greatly simplifies because we have $u = 1/\rho$ in that part of the sheath where the convective contribution to the magnetic and energy fluxes becomes significant (see Appendix, 3).

Let us start to solve eq's (8)-(10) from the unmagnetized plasma at the liner (domain III). Since the magnetic field here is rather small ($h = 0$ at $\xi = 0$), it follows from the pressure balance condition (8) that $\rho\theta \sim 1$. The main contribution to the magnetic (9) and energy (10) fluxes comes from the resistive magnetic diffusion and heat conductivity, so according to (11):

$$\varkappa \frac{d\theta}{d\xi} = \delta^2 \theta^{5/2} \frac{d\theta}{d\xi} \sim 1, \quad \theta^{-3/2} \frac{dh}{d\xi} \sim 1, \tag{13}$$

It follows from (12) and (13) that the distance from the liner $\xi \sim \delta^2 \theta^{7/2}$ and the magnetic field increase $\Delta h \sim \delta^2 \theta^4 d\theta$. Such a solution is valid up to $\theta \lesssim \theta_1 \sim \delta^{-2/5}$ where the electrons become magnetized (at $\theta \sim \theta_1$ the magnetic field $h \sim 1$). The thickness of this region $\Delta\xi = \xi_1 \sim \delta^2 \theta_1^{7/2} \sim \delta^{3/5}$. For $\xi > \xi_1$ the transition to domain (II) occurs, where the electrons are already magnetized ($\omega_{He}\tau_e > 1$) but the ions are not yet ($\omega_{Hi}\tau_i < 1$).

Here eq's (9) and (10) may be transformed, using (8) and (11), into the form:

$$\theta^{-3/2} \frac{dh}{d\xi} = -\frac{1}{2h\theta^{3/2}} \frac{d(\rho\theta)}{d\xi} = 1$$
$$\varkappa \frac{d\theta}{d\xi} = \delta \frac{\rho\theta}{h} \frac{d\theta}{d\xi} = (1-h) - \alpha\theta \frac{dh}{d\xi} = \rho\theta\left(\frac{1}{1+h} - \frac{1}{h}\right) \tag{14}$$

It is easy to see from the last relation that in region (II) the value of $d\theta/d\xi < 0$, so the temperature of a plasma reaches its maximum $\theta_{max} \sim \theta_1$ at $(\omega_H \tau)_e \sim 1$. Then, while the temperature drops on $\Delta\theta \sim \theta_1$, the plasma pressure decreases from $\rho\theta \sim 1$ to $\rho\theta \sim \mu^{1/2}$, where the ions become magnetized also (therefore in domain (II) the transition from the plasma with $\beta \sim 1$ to the low-β plasma occurs). It was assumed above that here the density of a plasma is high ($\rho \gg 1$), so convective contributions to the fluxes (9) and (10) are negligible. It is the case if $\mu^{1/2}\theta_1^{-1} \gg 1$, i.e. $\delta \gg \mu^{-5/4}$.[1]

As for the magnetized plasma (domain I), it follows from (8)-(11) that the plasma pressure decreases rapidly while the temperature remains approximately constant ($\theta \sim \theta_1$): $d\ln(\rho\theta)/d\ln\theta \sim \mu^{-1/2} \gg 1$. This is valid while the density is high: $\rho > 1$. For $\rho \simeq 1$ the temperature rapidly falls: $\theta \sim \mu^{-1}\xi^{-2}$. The typical thickness of the sheath $\Delta\xi$ is by the order of magnitude equal to $\delta^{3/5}$. The qualitative form of the whole solution is shown in Fig. 2.

For what follows it is important to know the amount of the plasma concentrated in the sheath, $\Pi_\Delta = \int(\rho - 1)d\xi$. It is easy to see from the solution obtained above that the main contribution to Π_Δ comes from the region of the sheath where the plasma pressure $\rho\theta \sim 1$ and the value of $\omega_{He}\tau_e \sim 1$:

$$\Pi_\Delta \sim \int_\Delta \rho\, d\xi \sim \theta_1^{-1}\xi_1 \sim \delta \gg 1 \tag{15}$$

Returning now to the usual variables and taking into account the cylindrical geometry of our system, the whole number of particles in the sheath may be written as follows:

$$N_\Delta = 2\pi R n_i \frac{\mathcal{D}_H}{v_i} \Pi_\Delta = 2\pi R n_i \frac{\mathcal{D}_H}{v_i} \alpha\delta = \alpha R c H_i / 2 e v_i \tag{16}$$

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[1] The final results are the same also for $\delta \lesssim \mu^{-5/4}$, and only the sheath structure slightly changes in this case.

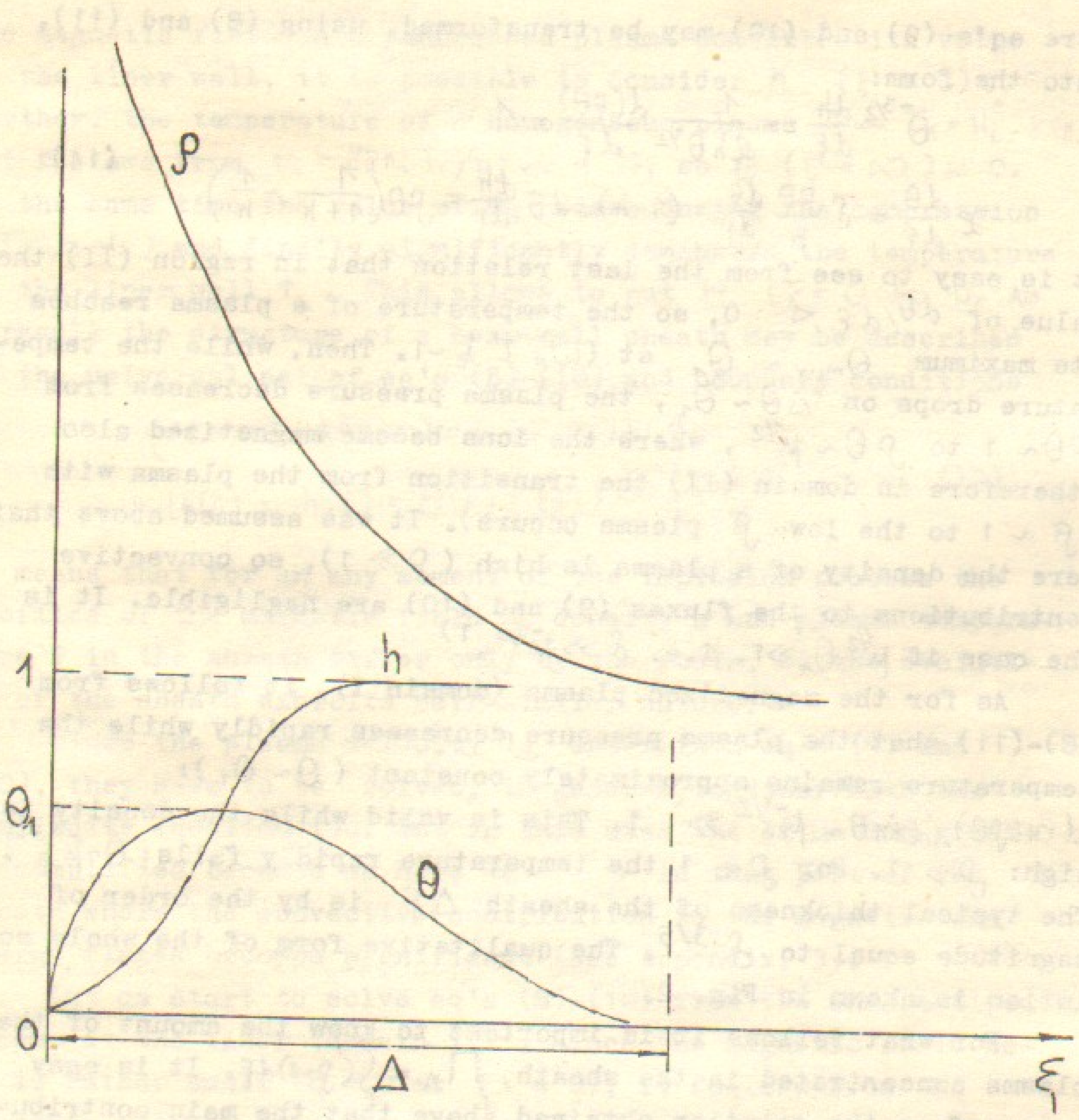

Fig. 2. The structure of the near-wall sheath

where α is a numerical factor which can be computed from the eq's (6') and (7') with the exact values of the plasma transport coefficients [5].

3. Magnetic field amplification in imploding plasma

Let us consider now the evolution in time of the homogeneus plasma during the implosion process. The retained magnetic flux $\Phi = \pi R^2 H_i$ decreases due to the convective flow into the sheath, so

$$\frac{d\Phi}{dt} = -2\pi R q_H = -2\pi R v_i H_i \qquad (17)$$

One more equation follows from the balance condition for the number of particles between the homogeneous plasma and the sheath:

$$\frac{dN_\Delta}{dt} = 2\pi R v_i n_i$$

Assuming now the implosion law $R_L(t)$ to be prescribed [2] and taking into account that the magnetic field is frozen into the homogeneous plasma: H_i/n_i = const = H_0/n_0 , it is possible to find the magnetic field $H_i(t)$ from the eq's (16)-(18). y introducing the new variables

$$\tau = t u_L/R_0, \quad r(\tau) = R/R_0 \quad x(\tau) = \Phi/\pi R_0^2 H_0, \quad y(\tau) = N_\Delta/\pi R_0^2 n_0$$

(here u_L is the typical velocity of a liner)
eq's (16)-(18) may be transformed as follows:

$$\frac{dx}{d\tau} = -\frac{1}{S_{eff}}\frac{x^2}{r^2 y}, \quad \frac{dy}{d\tau} = \frac{1}{S_{eff}}\frac{x^2}{r^2 y} \qquad (19)$$

where

$$S_{eff} = \pi n_0 e R_0 u_L / \alpha c H_0 \qquad (20)$$

is the effective magnetic Reynolds number. Since the sum x + y= const and at initial moment x = 1 and y = 0, the first of eq's (19) may be easely integrated:

$$\frac{1}{x} - 1 + \ln x = \frac{1}{S_{eff}} \int_0^\tau dt / r^2(t) \qquad (21)$$

It follows from (20) that the magnetic flux losses from the imp-

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2) The evolution in time of the radius of a plasma $R_L(t)$ has to be determined from the equation of motion for the liner and this is a separate problem.

loding plasma may be characterized by the effective magnetic diffusion coefficient

$$D_{eff} = \alpha c H_0/\pi e n_0 \qquad (22)$$

It is intresting to note that the value of D_{eff}. does not depend on the collision frequency in a plasma, though the flux losses are indebted to the finite plasma resistivity. The situation here is rather similar to that in a shock waves. The thickness of a shock front is formed that is sufficient to produce dissipation prescribed by Hugoniot relations. In our case the plasma resistivity and heat conductivity regulate the thickness of the sheath in such a way that the magnetic flux which enter in the sheath due to convection can be transfered to the wall by resistive diffusion.

For the illustration the solutions of eq.(21) for different effective magnetic Reynolds numbers S_{eff}. are shown in Fig.3. Here the magnetic field evolution in a plasma imploded by the liner with the constant velocity ($\zeta(\tau) = 1 - \tau$) is plotted. It is seen from Fig.3 that it needs $S_{eff} \gtrsim 30$ for the effecient ten-fold radii compression of the seed magnetic field.

The important difference between the results obtained above and the criteria (1) is that the value of S_{eff}. is proportional to the density of a plasma (while the usual magnetic viscosity of a plasma D_m doesn't depend on it). So, according to (20), it is profitable now to increase the initial density of a plasma n_0. Choosing approximately the same parameters of a system as discussed in [1,2], i.e. $R_0 \approx 10^{-1}$ cm and $H_0 \approx 0{,}2$ Mgs it follows from (20) that even for $n_0 \approx 10^{19}$ cm^{-3} it needs $u_L \approx 5\cdot 10^7$ cm/sec to achieve $S_{eff} \approx 30$, and this velocity greatly exceeds the values supposed in [1].3) It must be noted here that the requirements to the implosion rate may be even more hard with the liner deceleration at the final stage of the compression taken into account.

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3) The numerical factor α in (20) has been determined by the comparison of the graphs plotted in Fig.3 with the computer simulation results presented in [4]. It gives $\alpha \approx 0{,}4$.

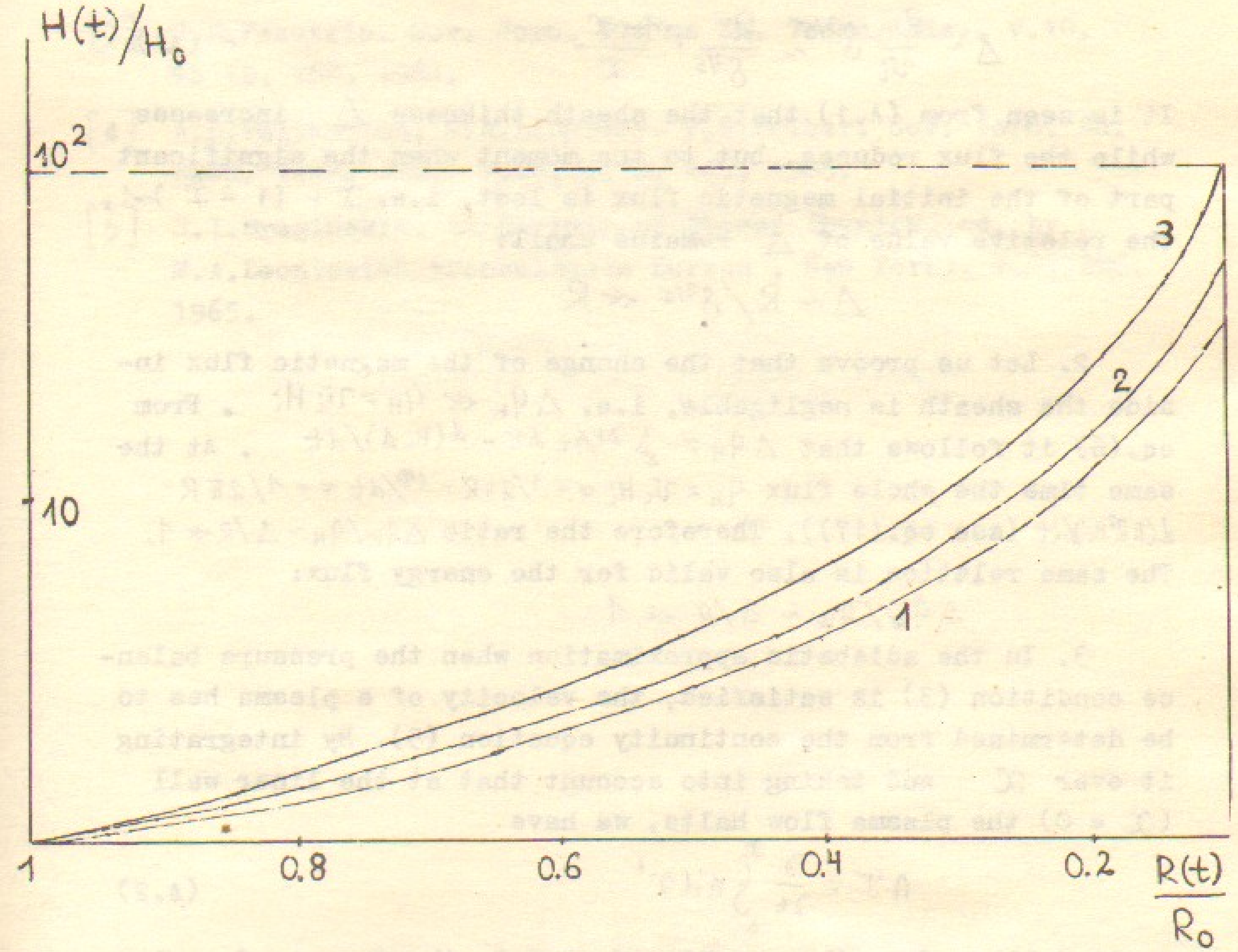

Fig. 3. The magnetic field compression for the different values of the effective magnetic Reynolds number S_{eff}. 1) $S_{eff.} = 10$; 2) $S_{eff.} = 30$; 3) $S_{eff} = \infty$.

Appendix

1. The solution described above was based on the fact the near-wall sheath thikness Δ is much less than the radius of a plasma R. Let us check this relation. According to estimations given in part 2, $\Delta \sim \delta^{3/5} \mathcal{D}_H / v_i$. Using now (16) and the dimensionless magnetic flux x , it is easy to show that

$$\Delta \sim \frac{\mathcal{D}_H}{v_i} \delta^{3/5} \sim \frac{R}{\delta^{2/5}} \cdot \frac{1-x}{x} \qquad (A.1)$$

It is seen from (A.1) that the sheath thikness Δ increases while the flux reduces, but to the moment when the significant part of the initial magnetic flux is lost, i.e. $x \sim (1 - x) \sim 1$, the relative value of Δ remains small:

$$\Delta \sim R / \delta^{2/5} \ll R$$

2. Let us proove that the change of the magnetic flux inside the sheath is negligable, i.e. $\Delta q_H \ll q_H = v_i H_i$. From eq.(6) it follows that $\Delta q_H = -\int_\Delta \partial H/\partial t \, dx \sim d(H_i \Delta)/dt$. At the same time the whole flux $q_H = v_i H_i = -1/2\pi R \cdot d\Phi/dt = -1/2\pi R \cdot d(\pi R^2 H_i)/dt$ (see eq.(17)). Therefore the ratio $\Delta q_H / q_H \sim \Delta / R \ll 1$. The same relation is also valid for the energy flux:

$$\Delta q_W / q_W \sim \Delta / R \ll 1$$

3. In the adiabatic approximation when the pressure balance condition (3) is satisfied, the velocity of a plasma has to be determined from the continuity equation (5). By integrating it over x and taking into account that at the liner wall (x = 0) the plasma flow halts, we have

$$n v = \frac{\partial}{\partial t} \int_0^x n \, dx' \qquad (A.2)$$

As it follows from the results of part 2, the amount of a plasma in the sheath is determined by the region with $x - x_1 \sim \xi_1 \mathcal{D}_H / v_i$ (see eq.(15)). So the integral in the r.h.s. of (A.2) ceases to depend on x when $x > x_1$. It means that for $x \gg x_1$ we have the constant plasma flux $n v = n_i v_i$, i.e., in dimensionless notations, $u = 1/\rho$ (and $u \leq 1/\rho$ for $x \leq x_1$). Therefore the convective contribution to the magnetic (9) and energy (10) fluxes is significant only in that part of a sheath where the density $\rho \sim 1$.

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ПОЛУЧЕНИЕ СВЕРХСИЛЬНЫХ МАГНИТНЫХ ПОЛЕЙ ПРИ СЖАТИИ ПЛАЗМЫ: УДЕРЖАНИЕ МАГНИТНОГО ПОТОКА

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