Hydrodynamic model of plasma "corona" produced when a charged-particle beam acts on a target

Yu. V. Afanas'ev, V. A. Isakov, and O. N. Krokhin

P. N. Lebedev Physics Institute, USSR Academy of Sciences
(Submitted 21 April 1981)
Zh. Eksp.Teor. Fiz. **81**, 1714–1722 (November 1981)

The interaction between charged-particle beams and plane and spherical targets is investigated theoretically. The parameters of the observed plasma "corona" are determined analytically within the framework of the hydrodynamic model, and the coefficient of conversion of beam energy into target kinetic energy is calculated as a function of the beam and target parameters. The results are compared with analogous data for the action of laser radiation on a target as applied to the problem of inertial thermonuclear fusion.

PACS numbers: 52.50.Dq

It has recently been suggested[1-4] that intense ion beams be used for inertial thermonuclear fusion. In the case of laser-mediated thermonuclear fusion (LTF), three stages of target compression by the ion beam can be distinguished in principle: corona formation accompanied by conversion of the beam energy into kinetic energy of the unevaporated part of the target material, compression, and thermonuclear burning. When ion beams are used, however, the physics of the corona formation exhibits a number of qualitative differences from the corresponding stage in both LTF and electron thermonuclear fusion; this is connected, naturally, with the radical difference in the character of the absorption of the supplied energy.[5-7] It is possible nevertheless, by starting from very simple assumptions concerning the mechanism whereby the ion-beam energy is transferred to the plasma, to construct a hydrodynamic model of the corona, in analogy with the LTF case.[8]

We consider in this paper several models that describe the formation of plasma when a beam of charged particles acts on a plane or spherical target. We determine in analytic form the physical parameters that characterize the process of evaporation and energy transfer to the unevaporated part of the target; these make it possible to compare the beam- and laser-mediated thermonuclear fusion. The dependence of the hydrodynamic coefficient of energy transfer on the beam and target parameters is calculated in analogy with Ref. 9.

I. PLANAR PROBLEM

1. We consider first the following problem. Let a planar particle beam with initial energy E_0 (per particle) and with flux density q_0 be incident on the half-space $x<0$ filled with matter of density ρ_0. The law that governs the stopping of the particles, which corresponds to Coulomb scattering by the electrons of the produced plasma, is represented in the form

$$dE/dx=a\rho/2E, \tag{1}$$

where ρ is the plasma density, $E=E(x)$ is the energy of the incident particle, $a=4\pi\Lambda e_j^2e^2z_im_j/m_eM_i$; m_j and e_j are the mass and charge of the beam particle; z_i and M_i are the ionization multiplicity and the mass of the plasma ions, and Λ is the Coulomb logarithm. We neglect next the scattering of the incident beam, i.e., $n_jv_j=\text{const}$ ($n_j(x)$ is the density of the beam particles), a procedure valid approximately only for heavy particles. It is clear also that (1) is valid so long as $v_j\gg v_i$.

From (1) it follows immediately that

$$dq/dx=q_0\rho/2m_0(1-m/m_0)^{1/2},\quad q=q_0(1-m/m_0)^{1/2},$$
$$m=\int_x^\infty\rho dx,\quad m_0=\int_0^\infty\rho dx=E_0^2/a. \tag{2}$$

If the density of the plasma $\rho\ll\rho_0$ and the beam power is high enough to be able to regard the plasma as fully ionized, two characteristic dimensional parameters q_0 and m_0 remain in the problem, which can then be regarded as self-similar.[10] As follows from (2), in this model (without allowance for other energy-transfer mechanisms) the mass of material that can be heated by the incident beam is limited and is equal to $m_0=E_0^2/a$, i.e., the problem reduces to unilateral expansion of a given mass m_0 that is bounded on the $x\leq 0$ side by a "rigid" wall impermeable to heat. Such problem with an arbitrary energy-release law $q=q(m/m_0)$ was solved by numerical methods in Ref. 10. Here we obtain an approximate analytic solution of this problem.

The system of hydrodynamic equations with energy release in the form (2) is

$$\frac{\partial\rho}{\partial t}+\frac{\partial}{\partial x}(\rho v)=0,\quad \frac{\partial}{\partial t}(\rho v)+\frac{\partial}{\partial x}(p+\rho v^2)=0,$$
$$\frac{\partial}{\partial t}\left[\rho\left(\varepsilon+\frac{v^2}{2}\right)\right]+\frac{\partial}{\partial x}\left[\rho v\left(\varepsilon+\frac{p}{\rho}+\frac{v^2}{2}\right)-q\right]=0, \tag{3}$$

$\varepsilon=p/(\varkappa-1)\rho$, $\varkappa$ is the adiabatic exponent.

Introducing the self-similar variable $\lambda=(m_0/q_0)^{1/2}xt^{-3/2}$ and the self-similar functions

$$\rho=q_0^{-1/2}m_0^{3/2}t^{-3/2}R(\lambda),\quad v=q_0^{1/2}m_0^{-1/2}t^{1/2}V(\lambda),$$
$$p=q_0^{1/2}m_0^{1/2}t^{-1/2}P(\lambda), \tag{4}$$

$$\frac{m}{m_0}=\int_\lambda^\infty R(\lambda)d\lambda,\quad \int_0^\infty R(\lambda)d\lambda=1,\quad \lambda\in[0,\infty),$$

we obtain

$$\frac{d}{d\lambda}\left[R\left(V-\frac{3}{2}\lambda\right)\right]=0,\quad \frac{1}{R}\frac{dP}{d\lambda}+\left(V-\frac{3}{2}\lambda\right)\frac{dV}{d\lambda}+\frac{V}{2}=0,$$
$$\frac{1}{\varkappa-1}\left(V-\frac{3}{2}\lambda\right)\frac{d}{d\lambda}\left(\frac{P}{R}\right)+\frac{P}{R}\frac{dV}{d\lambda}+\frac{1}{\varkappa-1}\frac{P}{R}-\frac{1}{2}\left(\int_0 Rd\lambda\right)^{-1/2}=0. \tag{3'}$$

The boundary conditions are of the form

0038-5646/81/110910-05$02.40

$$R(\infty)=P(\infty)=0, \quad V(0)=0 \tag{5}$$

(the rigid-wall condition).

From (3′) and from the second condition of (5) we have ($\varkappa=\frac{5}{3}$)

$$V=\frac{3}{2}\lambda, \quad \frac{dP}{d\lambda}=-\frac{3}{4}\lambda R, \quad P=\frac{R}{6}\left(\int_0^\lambda R\,d\lambda\right)^{-1/2}. \tag{6}$$

Since $P(0)$ is finite, we obtain $R(0)=0$.

We seek the solution with the aid of integral relations that express the mass, momentum, and energy conservation laws that take in the present case the form

$$\int_0^\infty R\,d\lambda=1, \quad \frac{1}{2}\int_0^\infty RV\,d\lambda=P(0), \quad \frac{1}{2}\int_0^\infty (3P+RV^2)\,d\lambda=1. \tag{7}$$

The problem (3′), with boundary conditions (5), has an approximate solution that satisfies exactly the integral conservation laws and can be represented in the form

$$V=\frac{3}{2}\lambda. \quad R\approx 4.0\lambda e^{-2,0\lambda^2}. \quad P\approx 0.47e^{-2,0\lambda^2}. \tag{8}$$

Assuming next that a rigid wall has a finite mass $M \gg m_0$, we obtain the value of the hydrodynamic transfer coefficient in the considered approximation. Obviously, the equation of motion of the incompressible part M is

$$M\frac{du}{dt}=p(0), \quad Mu=2q_0^{1/2}m_0^{1/2}t^{1/2}P_0,$$

from which we easily obtain an expression for the hydrodynamic efficiency:

$$\eta=\frac{Mu^2}{2}\Big/ q_0t=2P_0^2\frac{m_0}{M}\approx 0.44\frac{m_0}{M}, \quad \frac{m_0}{M}\ll 1. \tag{9}$$

2. The limiting value of the hydrodynamic efficiency in the case of LTF (planar problem) was calculated analytically in Ref. 11. In the approach used there, a hydrodynamic discontinuity boundary was introduced (an unevaporated rigid part–corona), on which the conditions of continuity of the mass, momentum, and energy fluxes are satisfied. The effect connected with the inertial forces that arise when the unevaporated part of a flat layer is accelerated were automatically take into account with the aid of an incident-radiation-flux "screening" coefficient q^*/q_0, where q_0 is the incident flux and q^* is the flux into the discontinuity. In our case, at a finite mass M of the rigid wall, this approach cannot be used, since the fluxes of all the quantities, including q^*, vanish on the discontinuity boundary, by virtue of the condition $v(0)=0$ (in the coordinate frame connected with the accelerating unevaporated part). To calculate the maximum hydrodynamic efficiency in the model where the beam interacts with a flat layer of mass $M_0=m_0+M$ it is necessary to take into account the acceleration of the unevaporated layer in explicit form. Since, however, this does not introduce any additional dimensional parameters in this situation, it is possible to generalize the problem considered above (and retain the self-similarity) to include also the case of a finite mass M at any value of the dimensionless parameter $\delta = m_0/M$.

In the coordinate frame connected with the unevaporated part of the layer M, moving with velocity $u=u(t)$ (in the laboratory coordinate frame), Eqs. (3) take the form

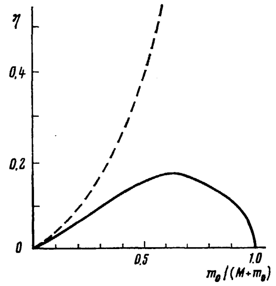

FIG. 1. Dependence of the hydrodynamic efficiency on the parameter $m_0/(M+m_0)$. The dashed curve corresponds to an approximation in which no account is taken of the acceleration of the unevaporated layer [see relation (19)].

$$\frac{\partial\rho}{\partial t}+\frac{\partial}{\partial x}(\rho v)=0, \quad \frac{\partial}{\partial t}(\rho v)+\frac{\partial}{\partial x}(p+\rho v^2)-\rho\frac{du}{dt}=0, \quad M\frac{du}{dt}=p(0),$$
$$\frac{\partial}{\partial t}\rho\left(\varepsilon+\frac{v^2}{2}-x\frac{du}{dt}\right)+\frac{\partial}{\partial x}\left[\rho v\left(\varepsilon+\frac{p}{\rho}+\frac{v^2}{2}-x\frac{du}{dt}\right)-q\right]=0. \tag{10}$$

The other relations for the functions of the self-similar variables take the form

$$V=\frac{3}{2}\lambda, \quad \frac{dP}{d\lambda}=-\frac{3}{4}\lambda R+\delta P(0)R, \quad P=\frac{R}{6}\left(\int_0^\lambda R\,d\lambda\right)^{-1/2}+\frac{1}{3}\delta P(0)\lambda R; \tag{11}$$

$$\int_0^\infty R\,d\lambda=1, \quad \frac{1}{2}\int_0^\infty RV\,d\lambda=(1+\delta)P(0),$$
$$\frac{1}{2}\int_0^\infty (3P+RV^2)\,d\lambda=1+\frac{4}{3}(1+\delta)\delta P^2(0). \tag{12}$$

The depencence of $\eta=2P_0^2\delta$ on the parameter m_0/M_0, obtained from (11) and (12), is shown in Fig. 1. As follows from the plots of the thermodynamic efficiency against m_0/M_0 in Fig. 1, allowance for the acceleration is essential and leads to perfectly natural results, namely the presence of a maximum η_m ($\eta_m\approx 16\%$) at $m_0/M_0\approx\frac{2}{3}$) and to $\eta=0$ at $m_0/M_0=1$ ($M=0$).

II. STATIONARY SPHERICAL PROBLEM

1. The planar problems considered above explain the main qualitative regularities of the interaction of an ion beam with the surface of a target under conditions when the energy dissipation is determined principally by the stopping of the beam. The latter is valid at sufficiently large values of

$$m_0=\int\rho\,dx\sim E_0^2,$$

exceeding the depth of heating via electronic thermal conductivity (a quantitative criterion will be given below). In the general case, however, the problem should be considered with the energy transport by electronic thermal conductivity also taken into account.[1)] The simplest physical model that permits an analytic treatment and includes the electronic thermal conductivity is connected with the problem of the stationary spherical corona, analogous to that considered in Ref. 8 as applied to LTF. In contrast to Ref. 8, however, where the energy release due to absorption of the incident radiation has a δ-like character at the point with critical density, the energy release in the case of a beam has in principle a distributed character. Nonetheless, the formal statement of the problem and the method of approximate analytic solution are quite similar in our case and in the case considered in Ref. 8.

In the spherical case we have in place of (2)

$$Q(R)=Q_0\left(1-g_0^{-1}\int_R^\infty \rho dr\right),\quad R>R_b,\quad g_0=\int_{Rb}^\infty \rho dr=E_0^2/a, \tag{13}$$

where $Q(R)$ is the beam energy flux per unit solid angle, $Q_0=Q(\infty)$, R_b is the radius corresponding to the vanishing of the flux $Q(R)$, i.e., $Q(R_b)=0$.

The condition

$$g_0=E_0^2/a=\int_{Rb}^\infty \rho dr \tag{14}$$

incorporates the principal physical and formal difference between our problem and the case of laser action.[8]

The equations for stationary flow of a plasma, with allowance for (13), and the boundary conditions, are given by

$$\begin{gathered}\rho v R^2=\rho_* v_* R_*^2,\quad \frac{d}{dR}(p+\rho v^2)=-\frac{2\rho v^2}{R},\\ \rho v R^2\left(\varepsilon+\frac{v^2}{2}+\frac{p}{\rho}\right)-R^2\varkappa_0 T^{5/2}\frac{dT}{dR}\\ =\begin{cases} Q_0\left(1-g_0^{-1}\int_R^\infty \rho dr\right), & R>R_b,\\ 0, & R<R_b,\end{cases}\end{gathered} \tag{15}$$

$$v(R_0)=T(R_0)=0,\quad v(\infty)=v_\infty=\text{const},\quad T(\infty)=0,\quad z_i T_*=M_i v_*^2,$$

where ρ, v, p, ε, and T are the density, velocity, pressure, specific internal energy, and electron temperature of the plasma; R_0 is the target radius ($\rho(R_0)=\infty$); ρ_*, v_*, and T_* are the values of the quantities at the Jouguet point[8] ($R=R_*$), and $\varkappa_0 T^{5/2}$ is the coefficient of electronic thermal conductivity.

Thus, the problem (15) jointly with the condition (14) contains five given parameters: $\varkappa_0$, Q_0, R_0, g_0, M_i/z_i (in the case of[8] we have $\varkappa_0$, Q_0, R_0, ρ_{cr}, and M_i/z_i). It is required to find the coordinate $R=R_b$ of the beam-absorption boundary, the velocity v_∞ of the material at infinity, and the scales of the quantities ρ_*, v_* and T_* together with the coordinate R_*.

2. We consider now an approximate similarity law that enables us to determine, accurate to numerical coefficients ~1, the scales of the quantities of interest to us.

Heating by electronic thermal conductivity can be characterized in the stationary case by the quantity g_T, which is comparable in dimensionality with the analogous quantity g_0 for a beam, i.e., with the "thickness" of the heating zone $g_T\approx\int\rho dr$:

$$g_T\sim\frac{\varkappa_0^{3/7}Q_0^{4/7}}{R_0^{4/7}}\left(\frac{M_i}{z_i}\right)^{3/2}. \tag{16}$$

Consequently, the dimensionless parameter of the problem, which determines the structure of the corona, takes in our case the form

$$\gamma_0\approx\frac{g_T}{g_0}\sim\frac{\varkappa_0^{3/7}Q_0^{4/7}}{g_0 R_0^{4/7}}\left(\frac{M_i}{z_i}\right)^{3/2}. \tag{17}$$

Indeed, it is physically clear that at $\gamma_0\gg1$ the characteristic parameters of the corona should not depend on g_0 and are determined completely by the thermal conductivity. Assuming in this case

$$\varkappa_0 T_*^{7/2}/R_*\sim\rho_* v_*^3\sim Q_0/R_*^2,\quad g_T\sim\rho_* R_*,\quad R_*\sim R_0,$$

we obtain

$$\begin{gathered}v_*^2\sim\varkappa_0^{-2/7}Q_0^{2/7}R_0^{-2/7}(z_i/M_i),\\ T_*\sim\varkappa_0^{-2/7}Q_0^{2/7}R_0^{-2/7},\\ \rho_*\sim\frac{\varkappa_0^{3/7}Q_0^{4/7}}{R_0^{11/7}}\left(\frac{M_i}{z_i}\right)^{3/2},\end{gathered} \tag{18}$$

which coincides naturally with the result of Ref. 8 under conditions when thermal conductivity predominates.

A relation for R_b can be obtained from the continuity equation

$$\begin{gathered}(\rho_b R_b)v_b R_b=(\rho_* R_*)v_* R_*,\\ g_0 R_b\sim g_T R_0 v_*/v_b,\quad R_b\sim R_0\gamma_0 v_*/v_b\sim R_0\gamma_0.\end{gathered}$$

In the case $\gamma_0\lesssim1$ we obtain

$$\begin{gathered}\rho_* R_0\sim g_0,\quad Q_0\sim\rho_* v_*^3 R_0^2,\\ v_*^2\sim\left(\frac{Q_0}{g_0R_0}\right)^{2/3},\quad T_*\sim\frac{M_i}{z_i}\left(\frac{Q_0}{g_0R_0}\right)^{2/3},\quad \rho_*\sim\frac{g_0}{R_0}.\end{gathered} \tag{19}$$

The formal procedure for solving the problem (14), (15) is similar to that used in Ref. 8. In particular, at $\gamma_0\ll1$, when the beam-absorption boundary reaches the target surface, (i.e., $R_b\approx R_0$), we obtain

$$\begin{gathered}R_*=1.15\,R_0,\quad \rho_*=1.1\,g_0/R_0,\\ v_*^2=0.24(Q_0/g_0R_0)^{2/3},\quad v_\infty^2=11.5\,v_*^2.\end{gathered} \tag{20}$$

III. ACCELERATION OF A THIN SPHERICAL SHELL BY AN ION BEAM

By way of a practical application of the results of Sec. II, it is of interest to consider the problem of acceleration, by an ion beam, of a target constituting a thin spherical shell. The solution of this problem will make it possible to compare the capabilities of ion and laser beams from the point of view of obtaining optimal hydrodynamic efficiency.

As applied to LTF, an approximate solution of this problem is given in Ref. 9. As shown above, the fundamental difference between a corona produced by laser radiation and that produced by particle beams appears under conditions when the role of thermal conductivity is small.

The equations of motion of the thin spherical shell take in this case a form that is perfectly analogous to that given in Ref. 9:

$$M\frac{du}{dt}=2\rho_* v_*^2 R_*^2,\quad \frac{dM}{dt}=-\rho_* v_* R_*^2, \tag{21}$$

where M is the mass of the unevaporated part of the target, u is the velocity of this part, and R_*, ρ_*, and v_* are the parameters of the corona and are determined by Eqs. (20), in which R_0 is now a function of the time. Its approximation is valid because dR_0/dt is small compared with the characteristic expansion rate $\sim v_*$ of the corona. The substitutions $\mu=M/M_0$, $w=u/v_0^*$, $x=R/R_0$, where M_0, R_0, v_0 are the initial values of the mass, of the radius of the unevaporated part of the target, and of the velocity v_*, and $R(t)$ is the running radius $[R_*=R_*(R),\ \rho_*=\rho_*(R),\ v_*=v_*(R)]$, reduce Eqs. (21) to the system

$$\mu w\frac{dw}{dx}=-2\alpha x^{1/3},\quad w\frac{d\mu}{dx}=\alpha x^{2/3} \tag{22}$$

with initial conditions

$$w(1)=0,\quad \mu(1)=1. \tag{23}$$

The system (22) contains a single parameter, equal to

$$\alpha\approx 1.3\,g_0/\Delta_0\rho_{sh}, \tag{24}$$

where Δ_0 and ρ_{sh} are the initial thickness and density of the shell. We have $\alpha<1$, since the model of the continuous evaporation is obviously valid at $g_0\ll\Delta_0\rho_{sh}$.

The approximate solution of Eqs. (22) with the initial conditions (23) is of the form

$$w=(3\alpha)^{1/2}(1-x^{4/3})^{1/2},\quad \mu=\exp\left[-\alpha^{1/2}F(x)\right],$$
$$\eta=0.33\,\alpha^{1/2}(1-x^{4/3})^{1/2}\exp\left[-\alpha^{1/2}F(x)\right], \tag{25}$$
$$F(x)=0.6\int_{\infty}^{1}\frac{x^{2/3}dx}{(1-x^{4/3})^{1/2}}.$$

Relations (25) can be extended to include the case $g_0\sim\Delta_0\rho_{sh}$, by replacing α with

$$\alpha^*=1.3\,g_0/(\Delta_0\rho_{sh}-g_0). \tag{26}$$

The introduction of α^* corresponds physically to dividing the shell-compression process into two stages. During the first stage the corona is formed by interaction of the ion beam with a part of the target shell with practically fixed mass

$$m_0\approx R_0^2\int\rho dr=R_0^2g_0.$$

In this sense the problem is planar and was solved in Sec. I of the present paper.

During the second stage the spherical shell with initial mass $\sim(M_0-m_0)$ undergoes an acceleration described by Eqs. (21)–(25) with α replaced by α^*. The large inertia of the dense shell compared with a rarefied corona ($V_{sh}\ll V_{cor}$) makes it possible to retain the condition $w(1)=0$.

The dependence of the maximum hydrodynamic efficiency in the collapse of the shell on α^* is shown in Fig. 2. The same figure shows the dependence of the target mass not yet evaporated by that instant on α^*. As follows from Fig. 2, the transfer coefficient does not exceed 0.085.

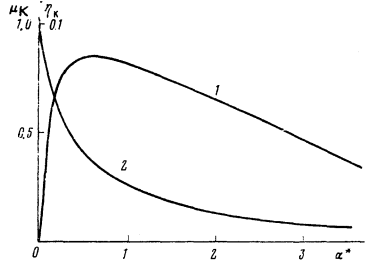

FIG. 2. Dependences of the hydrodynamic efficiency (1) and of the mass of the unevaporated part of the shell (2) at the instant of collapse (x = 0) on the parameter α^*.

With account taken of the two-stage character of the shell compression, we can write for the hydrodynamic efficiency the following approximate expression:

$$\eta=\eta_1\tau+\eta_2(1-\tau), \tag{27}$$

where η_1 and η_2 are the efficiencies reached during the stages of the "planar" and "spherical" motion (Figs. 1 and 2, respectively), and τ is the relative time interval during which the shell is in the stage of planar motion in the course of the compression. At $g_0\sim\Delta_0\rho_{sh}$ ($\alpha^*\gtrsim 1$) the contribution of the planar shell acceleration should predominate in the estimate (27) of the hydrodynamic efficiency. From the solution of the planar problem with a fixed value of the unevaporated mass it follows that $\eta_1\lesssim 16\%$. For the case of laser action[9]

$$\alpha_L\approx 1.5\frac{\rho_{cr}}{\rho_{sh}^L}\frac{R_0}{\Delta_0}\approx 0.1$$

[neodymium laser and thin ($R_0/\Delta_0\sim 10^2$) shell targets], corresponding to $\eta\approx 5\%$. It must be kept in mind, however, that the regime corresponding to $g_0\sim\Delta_0\rho_{sh}$ may turn out to be physically closer to the "exploding piston" regime.[13] Therefore a correct comparison of the processes of laser and beam acceleration of shell targets is possible only in the regime of continuous evaporation of the shell ($g_0\ll\Delta_0\rho_{sh}$). In this case $\alpha<1$.

To compare the actions of an ion beam and of laser radiation, we represent α in the form

$$\alpha=1.3\frac{g_0/R_0}{\rho_{sh}}\frac{R_0}{\Delta_0}$$

from which it follows that the analog of the critical density in the case of beams is the quantity g_0/R_0 ($g_0\ll\Delta_0\rho_{sh}$).

We present now an expression for $g_0/R_0\rho_{cr}$, which enables us to establish approximately the equivalent parameters of the beam and of the laser

$$\frac{g_0/R_0}{\rho_{cr}}=\frac{3}{\Lambda A_jz_j^2}\frac{1}{R_0[\text{cm}]}\left(\frac{E_0[\text{MeV}]}{\hbar\omega_0[\text{eV}]}\right)^2, \tag{28}$$

where Λ is the Coulomb logarithm, A_j and z_j are the atomic weight and the charge of the beam ions, and $\hbar\omega_0$ is the energy of the laser-radiation quantum.

It follows from (28), e.g., that in the case of a proton beam ($A_j=z_j=1$)

$$\frac{g_0/R_0}{\rho_{cr}}\approx\frac{0.3}{R_0[\text{cm}]}\left(\frac{E_0[\text{MeV}]}{\hbar\omega_0[\text{eV}]}\right)^2,$$

and therefore, at equal shell densities and R_0/Δ_0, and also at $g_0\ll\Delta_0\rho_{sh}$, a beam of protons with energy of several MeV is approximately equivalent to the emission of a neodymium laser ($\hbar\omega_0\approx 1$ eV) for targets with $R_0\lesssim 1$ cm.

In experiments with ion beams the targets are usually made of denser material than in laser experiments. The condition for the equivalence of the beam and the laser is in the relation $g_0/R_0\rho_{cr}\sim\rho_{sh}/\rho_{sh}^L$, which changes somewhat the estimate given above in favor of the laser radiations. An estimate of the satisfaction of the condition $g_0/\Delta_0\rho_{sh}<1$ in the case of a proton beam yields $\Delta_0 \rho_{sh} > 10^{-3} E_0^2$[MeV]. For targets with $R_0/\Delta_0 \sim 10^2$ this condition can be easily satisfied for protons of energy ~1 MeV at a target radius $R_0 \gtrsim 10^{-1}$ cm.

The authors thank V. G. Mikhaĭlov for help with the numerical calculations, and E. G. Gamaliĭ and V. B. Rozanov for a discussion and for helpful remarks.

[1)] Generally speaking, energy can be fed to the surface of the target also as a result of other kinetic effects that occur when an ion beam interacts with the corona plasma (see, e.g., Ref. 12). We confine ourselves here to allowance for the electronic thermal conductivity, since it will be shown below that the situation in which a substantial role in the release and transfer of energy is assumed by effects that are not connected with the stopping of the ion beam is not optimal from the point of view of reaching high efficiencies. Therefore allowance for these effects, from the viewpoint of comparing laser and beam LTF, introduces in principle nothing new.

[1]F. Winterberg, Plasma Phys. **17**, 69 (1975).
[2]J. M. Shearer, Nuclear Fusion **15**, 952 (1975).
[3]J. W. Poukey, J. R. Freeman, M. J. Clauser, and G. Yonas, Phys. Rev. Lett. **35**, 1806 (1975).
[4]L. I. Rudakov, Fiz. Plazmy **4**, 72 (1978) [Sov. J. Plasma Phys. **4**, 40 (1978)].
[5]S. L. Bogolyubskiĭ, B. P. Gerasimov, V. I. Liksonov, Yu. P. Popov, L. I. Rudakov, A. A. Samarskiĭ, V. P. Smirnov, and L. I. Urutskoev, Pis'ma Zh. Eksp. Teor. Fiz. **24**, 202 (1976) [JETP Lett. **24**, 178 (1976)].
[6]M. J. Clauser, Targets for Heavy Ion Fusion, Preprint SNAD-78-0472C, 1978.
[7]A. V. Dobkin, I. B. Kosarev, and I. V. Nemchinov, Zh. Tekh. Fiz. **49**, 1405 (1979) [Sov. Phys. Tech. Phys. **24**, 782 (1979)].
[8]Yu. V. Afanas'ev, E. G. Gamaliĭ, O. N. Krokhin, and V. B. Rozanov, Zh. Eksp. Teor. Fiz. **71**, 594 (1976) [Sov. Phys. JETP **44**, 311 (1976)].
[9]Yu. V. Afanas'ev, E. G. Gamaliĭ, S. Yu. Gus'kov, and V. B. Rozanov, FIAN Preprint No. 30, 1980.
[10]I. V. Nemchinov, Prikl. Mat. Mekh. **31**, 300 (1967).
[11]Yu. V. Afanas'ev, E. G. Gamaliĭ, O. N. Krokhin, and V. B. Rozanov, *ibid.* **71**, 594 (1975).
[12]A. V. Gurevich, Pis'ma Zh. Eksp. Teor. Fiz. **25**, 469 (1977) [JETP Lett. **25**, 441 (1977)].
[13]E. K. Storm, H. G. Anestrom, M. J. Boyle, *et al.*, Laser Fusion Experiments at 2 TW, Preprint UCRL-78581, 1976.

Translated by J. G. Adashko