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Hydrodynamic instability and spontaneous magnetic fields in a spherical laser plasma

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We consider the magnetic fields generated in a spherical laser target by the appearance of thermoelectric currents in the region between the thermal-wave front and the surface on which the hydrodynamic velocity is equal to the local velocity of sound (the Jouguet point). Noncollinear temperature and density gradients are produced by small perturbations and by the development of a Rayleigh-Taylor instability. For targets in the form of glass shells, which are being studied in contemporary experiments, these fields amount to $\sim 10^6$ G at laser-irradiation energies $\sim 10^2$ J [N. G. Basov *et al*., JETP Lett. 23, 428 (1976)] and are capable of magnetizing the plasma in the indicated region ($\omega_e \tau_e \sim 1$). Outside the target, the fields decrease rapidly to less than one gauss at a distance ~ 1 mm.

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1. Magnetic fields of appreciable magnitude (up to 10^6 G) have been observed in experiments on the interaction of laser radiation with the matter in the produced plasma.[1-5] Although many experiments have not been uniquely interpreted, the possibility of formation of strong magnetic fields in a laser produced plasma is subject to no doubt. The presence of a field of $\sim 10^6$ G can significantly alter the transport coefficients of the plasma, influence the transport of "fast electrons" from the "corona" of the target, deform the profile of the electron density in the corona, and change the character of the evolution of the hydrodynamic perturbations. The strongest effect can be exerted on the target compression by the plasma magnetization in the region between the zone where the laser radiation is absorbed and the front of the thermal wave, since this region is responsible for the formation of the pressure pulse.

The present paper deals with the generation of the fields in this region. The field generation can be the result of a number of mechanisms: a) resonant absorption of plane-polarized light[5] and anisotropy of the light pressure[7]; b) magneto-thermal instability[8,9]; thermoelectric power in inhomogeneous plasma.[3] We shall show that at radiation fluxes $\sim 10^{14}$ W/cm^2 and at the plasma parameters in the indicated zone ($N_e \sim 10^{21-22}$ cm^{-3}, $T_e \sim 1$, eV) the main contribution to the generation of the field is made by a mechanism connected with the thermoelectric power in the inhomogeneous plasma. Let us estimate the fields for the aforementioned mechansims.

a) In the general case, the magnetic-field generation is connected with the anisotropy of the energy-momentum tensor T^{ik} of the electromagnetic field[5,7]:

$$\frac{\partial \mathbf{B}}{\partial t} = \frac{c}{e} \operatorname{rot} \frac{\mathbf{f}}{N_e}, \qquad f^i = \frac{\partial T^{ik}}{\partial x^k}.$$

In particular, magnetic fields can result from an increase in the longitudinal component of the electric field in the incident wave as the critical density is approached.

Let us estimate the magnetic field in the case when the presence of the longitudinal component is due to oblique incidence of the wave on the critical surface,[10] while the saturation mechanism is connected with the fact that the electrons carry the field out of the generation region[5]:

$$B \leqslant \frac{c}{JL_0} \left(\frac{\omega}{\nu}\right)^2 \frac{E_0^2}{8\pi} \sin\theta \cos\theta \exp\left[-\frac{4}{3} k_0 L_0 \sin^3\theta - 2k_0 L_0 \left(\frac{\nu}{\omega}\right)^2\right],$$

where $J = eN_e v_e$, $L \sim 3\lambda_0$, $\nu \sim 0.1\omega_p$ ($\omega_p = (4\pi N_e e^2/m_e)^{1/2}$ is the plasma frequency), and $\theta = 10°$. For the specified parameters we obtain

$$B \leqslant 3 \cdot 10^4 \text{ G}.$$

b) The magneto-thermal instability mechanism causes the field to increase in the case when the direc-

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tions of the temperature and density gradients are the same.[8] This situation is realized on the periphery of the corona, in the plasma that moves away from the target at supersonic velocity. This field-generation mechanism does not operate in the region considered here.

c) In an inhomogeneous plasma, owing to the appearance of the thermoelectric currents, we have

$$\frac{\partial \mathbf{B}}{\partial t} \sim \frac{c}{eN_e}[\nabla T_e \times \nabla N_e].$$

At the characteristic parameters $T_e \sim 1$ keV, $N_e \sim 10^{21}$ cm^{-3}, $L \sim 10^{-2}$ cm, and $t \sim 10^{-9}$ sec we obtain $B \sim 10^6$ G.

When a laser interacts with a planar target, the onset of crossed temperature and density gradients (and hence, of thermoelectric power) is determined by the experimental geometry itself. In an ideal spherical target, the gradients are collinear and there are no fields. Actually, deviations from spherical symmetry of either the target shape and density or of the laser radiation flux are always present. The hydrodynamic instability can lead to an appreciable growth of these initially small perturbation. The possibility of magnetic-field generation via Rayleigh-Taylor instabilities was discussed in[11] using as an example the idealized problem of adiabatic expansion of a plane layer in the field of constant acceleration with oppositely directed temperature and density gradients, specified on the boundaries of the layer at the initial instant of time.

In this paper we consider the appearance of fields in a narrow zone near the front of a thermal wave in a spherical target, where the Rayleigh-Taylor instability growth rates are maximal. To formulate this problem it is necessary to know in detail the spatial distribution of the hydrodynamic parameters of the plasma between the critical density and the front of the thermal wave. The investigations were carried out both by analytic methods[12] and with numerical experiments.[13-16] The theoretical data agree well with the experimental results. The instability zone is determined by the presence of oppositely directed density and pressure gradients.[17] We calculate below the magnetic fields produced in this zone in a spherical target.

2. We derive, following,[6] equations that describe the generation of the magnetic fields. It is known[18] that the equation of motion of electrons in a magnetized plasma is of the form

$$m_e N_e \frac{d\mathbf{V}_e}{dt} = -\nabla p_e - N_e e\left(\mathbf{E} + \frac{1}{c}[\mathbf{V}_e \times \mathbf{B}]\right) + eN_e\frac{\mathbf{J}}{\sigma} + \mathbf{R}_T. \quad (1)$$

Here N_e, $\mathbf{V}_e$, and p_0 are respectively the density, velocity, and pressure of the electrons; E and B are the electric and magnetic fields, J is the conduction current, σ is the electric conductivity, and R_T is the thermoelectric power. We have left out of (1) a term with the viscous-stress tensor, inasmuch as the viscosity plays a minor role in the hot ($T \sim 1$ keV) and dense ($N_e \sim 10^{21-22}$ cm^{-3}) plasma which we are considering.

Assuming the flow to be quasistationary, we can neglect in (1) the inertia term, and then we readily obtain from Maxwell's equations and (1)

$$\frac{\partial \mathbf{B}}{\partial t} = \mathrm{rot}[\mathbf{V}_e \times \mathbf{B}] + \frac{c^2}{4\pi\sigma}\nabla^2\mathbf{B} + \frac{c}{e}\mathrm{rot}\left(\frac{\nabla \mathbf{p}_e}{N_e} - \frac{\mathbf{R}_T}{N_e}\right). \quad (2)$$

We have neglected the displacement current in the derivation, a legitimate procedure since $V_e \ll c$. The first term describes convection and the second diffusion of the field, while the third contains the field source due to the thermoelectric currents. The last term, which contains the thermoelectric power, includes, in particular, the removal of the field by the hot electrons.[8]

In the linear approximation, the electron density and the electron temperature in the corona can be represented in the form

$$N_e(r, \theta, \varphi, t) = N_0(r, t) + n(r, \theta, \varphi, t), \quad n \ll N_0, \quad (3)$$

$$T_e(r, \theta, \varphi, t) = T_0(r, t) + \tau(r, \theta, \varphi, t), \quad \tau \ll T_0. \quad (4)$$

The evolution of the perturbations of the density n and τ can be determined in principle from the solution of the linearized problem in analogy with[19]:

$$n(r, \theta, \varphi, t) = \sum_{lm} n_{lm}^{(0)}(r)\exp(\gamma_{lm}t)Y_{lm}(\theta, \varphi), \quad (5)$$

$n_{lm}^{(0)}$ is the initial perturbation, $Y_{lm}(\theta, \varphi)$ is a spherical harmonic, and γ_{lm} is the perturbation growth rate.

It is known[20] that in the absence of dissipative processes the perturbation growth rate is

$$\gamma = (gl)^{1/2}, \quad (6)$$

where g is the acceleration and l is the number of the harmonic, and this rate grows as $l \to \infty$. Stabilization of the short-wave modes is reached either through the presence of viscosity, or if the density and pressure gradients are finite.[21] The latter case corresponds to our situation. Thus, the greatest growth rate is possessed in this case by perturbations with wavelength $\lambda \sim L_{\text{char}}$, where L_{char} is the characteristic dimension of the gradient.[22] To estimate the field we can use the expression for the growth rate[19,22]

$$\gamma_{l\gg 1} = \left(\frac{\nabla p\, \nabla N_e}{M_i N_e^2}\right)^{1/2} \quad (7)$$

(M_i is the ion mass). A schematic radial distribution of the hydrodynamic quantities near the front of the thermal wave is shown in Fig. 1a. The growth rate (7) is positive in the region where the pressure and density gradients are oppositely directed (the instability zone). The perturbations are carried out of this zone by the plasma stream either into the interior of the target or into the corona, and at $V > c_s$ (c_s is the local sound velocity) the perturbations become damped. To estimate the field it is therefore natural to assume that the perturbations are concentrated only in the instability region and there are none outside this region.

3. Let us calculate the growth rate of the magnetic field during the linear stage of the perturbation development, without allowance for the saturation mechanisms.

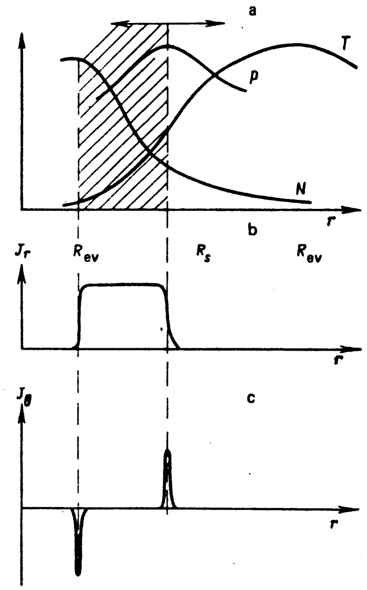

FIG. 1. Radial distributions of the temperature, pressure, and density (a) and of the radial (b) and angular (c) components of the currents near the front of the thermal wave. The arrows indicate the plasma motion direction. The shaded area is the instability zone, R_{ev} is the evaporation boundary, R_s is the Jouguent point, and R_{cr} is the critical radius.

Then, using (2) and (3) and neglecting the terms quadratic in the perturbations, we get

$$\frac{\partial \mathbf{B}}{\partial t}=\frac{c}{e}[\nabla T_e\times\nabla\ln N_e]=\frac{c}{e}\left([\nabla\tau\times\nabla\ln N_0]+\left[\nabla T_0\times\frac{\nabla n}{N_0}\right]\right). \tag{8}$$

Since the perturbation growth rate differs from zero in the instability region, and the gradients are maximal inside or near this region, the field turns out to be concentrated in a relatively narrow region, comparable with the size of the instability zone, near the front of the thermal wave. In this approximation the field has only φ and θ components (the rth component is of higher order of smallness the perturbation). For perturbations of the form (3) and (4) we get from (8)

$$\frac{\partial B_\theta}{\partial t}=-\sum_{lm}\dot{B}_{lm}(r,t)\frac{1}{\sin\theta}\frac{\partial Y_{lm}}{\partial\varphi},$$
$$\frac{\partial B_\varphi}{\partial t}=\sum_{lm}\dot{B}_{lm}(r,t)\frac{\partial Y_{lm}}{\partial\theta}, \tag{9}$$
$$\dot{B}_{lm}=\frac{c}{er}\left\{n_{lm}^{(0)}(r)\exp(\gamma_{lm}^{(n)}t)\frac{\nabla_r T_0}{N_0}-\tau_{lm}^{(0)}(r)\exp(\gamma_{lm}^{(\tau)}t)\nabla_r\ln N_0\right\},$$

the superscripts (n) and (τ) distinguish between the density and temperature growth rates, respectively.

Using (8) and Maxwell's equation with the displacement current neglected, we determine the eddy currents that generate the fields

$$\mathbf{J}=\frac{c}{4\pi}\operatorname{rot}\mathbf{B}. \tag{10}$$

This approximation, which is natural for all magnetohydrodynamics problems in general[18] is valid if $(V_e/c)^2\ll 1$ or $r_D/L\ll 1$ (r_D is the Debye radius and $L=T/|\nabla T|$). Both conditions are satisfied in our case with a margin of several orders of magnitude, meaning that the role of charge separation is negligibly small in the generation of the fields. It is easily seen that in the general case all three current components are different from zero:

$$J_r=\frac{c}{4\pi}\frac{1}{r\sin\theta}\sum_{lm}B_{lm}\left[\frac{\partial}{\partial\theta}\sin\theta\frac{\partial Y_{lm}}{\partial\theta}+\frac{1}{\sin\theta}\frac{\partial^2 Y_{lm}}{\partial\varphi^2}\right],$$
$$J_\theta=-\frac{c}{4\pi}\frac{1}{r\sin\theta}\sum_{lm}\frac{\partial rB_{lm}}{\partial r}\frac{\partial Y_{lm}}{\partial\theta}, \tag{11}$$
$$J_\varphi=-\frac{c}{4\pi}\sum_{lm}\frac{1}{\sin\theta}\frac{\partial Y_{lm}}{\partial\varphi}\frac{1}{r}\frac{\partial rB_{lm}}{\partial r}$$

where

$$B_{lm}=\int\dot{B}_{lm}\,dt.$$

4. Let us calculate the fields for the case when the initial perturbations of the density have axial symmetry and the temperature is uniform. We consider for simplicity the case when the distribution of the hydrodynamic quantities in the corona is quasistationary ($T_0=T_0(r); N_0=N_0(r)$); this is justified, since the characteristic time of removal of the perturbations (~10^{-10} sec) is smaller by one order of magnitude than the characteristic time of target compression (~10^{-9} sec). Then the time dependence is determined by the exponential terms

$$n_l(r,\theta,t)=\sum_l n_l^{(0)}(r)\exp(\gamma_l t)\cos l\theta. \tag{12}$$

In this case only the φ components of the field differs from the zero

$$B_\varphi=-\sum_l\frac{c}{er}\frac{n_l^{(0)}(r)}{N_0}\frac{\partial T_0}{\partial r}\frac{\exp(\gamma_l t)-1}{\gamma_l}l\sin l\theta=\sum_l B_l\sin l\theta. \tag{13}$$

The field amplitude is larger the larger the relative amplitude of the initial density perturbation $n_l^{(0)}/N_0$ and the radial temperature gradient, and the smaller the perturbation wavelength $\lambda_l=2\pi r/l$.

The current has r and θ components:

$$J_r=\sum_l\frac{J_l}{r^2}\frac{1}{\sin\theta}\frac{\partial}{\partial\theta}\sin\theta\sin l\theta,$$
$$J_\theta=\sum_l\frac{1}{r}\frac{\partial J_l}{\partial r}\sin l\theta, \tag{14}$$

where

$$J_l(r,t)=\frac{c^2}{4\pi e}\frac{l(\exp(\gamma_l t)-1)}{\gamma_l}\frac{n_l^{(0)}}{N_0}\nabla_r T.$$

According to,[17] the perturbations are concentrated in a narrow zone whose dimension is of the order of several microns, and the field growth time is much shorter than the laser pulse, and consequently the temperature gradient $\nabla_r T$ and the relative density-perturbation amplitude can be regarded as constant along the radius in this zone. The radial distribution of the currents for the axially symmetrical case is shown in Figs. 1b and 1c.

Using (14), we can construct the pattern of the eddy currents in the instability zone (see Fig. 2). These are annular currents and their direction varies with a period $2\pi/l$, while the rth component of the current has zeros at points whose values of θ are spaced π/l apart.

5. Let us examine the field in a spherical target constituting a thin-wall glass shell ($R \approx 50$ μm, $\delta R \sim 1-2$ μm). The absorbed energy in experiments with such targets is 3–10 J at laser pulse durations 0.5–2.5 nsec.[13] The amplitude of the initial perturbation, according to measurements of the nonsphericity of the glass microballoons used in the experiments, is several percent of the glass-shell thickness. Using the results of the numerical calculations[13] to determine the acceleration and the gradients in the corona, we can find that the maximum growth rates determined from formulas (6) and (7) are

$$\gamma \sim 10^{10} \text{ sec}^{-1}.$$

Thus, at $\gamma \sim 10^{10}$ sec^{-1}, at a relative density perturbations $n^{I[0]}/N_0 = 10^{-2}$, $\lambda \simeq 2 \cdot 10^{-4}$cm ($l = 2\pi R/\lambda = 150$), $\nabla_r T_0 = 10^{-6}$ erg/cm, and assuming $t \sim 1/\gamma_l$, we obtain from formula (13) the value $B_\varphi = 5 \times 10^6$ G. Thus, even one-percent perturbations of the density can lead to generation of substantial fields. The saturating mechanisms are connected with convection ($t_{\text{conv}} \sim x/V$), diffusion of the field, and rotation of the gradients[24] ($t_{\text{dif}} \sim 4\pi\sigma x^2/c^2$, $\sigma = 3.5 \times 10^{17}$ sec^{-1}). Estimates show that at the indicated plasma parameters the saturating mechanisms are capable of decreasing the calculated value to $B_\varphi = 10^8$ G.

A field of this size can lead to a severalfold decrease of the radial thermal conductivity (perpendicular to the field). In addition, the Larmor radius of the fast electrons produced in the laser-radiation absorption zone amounts to several microns in this field (at electron energies of several dozen keV). This means that the fast electrons will also be decelerated near the instability zone. These two effects raise the temperature in the indicated zone.[25] Convection of the field in the direction of decreasing density causes the magnetization of the plasma to be appreciable over distances much larger than the generation zone. Since B/ρ is constant (the field is frozen-in) and $\omega_e \tau_e \sim T_e^{3/2} B/\rho$, it follows that the magnetization can increase when the field drifts to the critical surface. We emphasize, however, that the processes that take place in this region are essentially non-one-dimensional.

The magnetic pressure in the instability zone turns out to be of the order of one-tenth the kinetic pressure and can influence the formation of the density profile in this region. The ponderomotive forces can exert a great influence on the development of the hydrodynamic instability. A correct calculation of the instability growth rates must take into account the action of the magnetic field.

6. Knowing the expressions for the currents in the target [(10) and (14)] and using the usual procedure of calculating the vector potential[26] we can obtain expressions for the magnetic field outside the target at distances greatly exceeding its radius ($r \gg R$). This procedure can be carried out even for the case of the arbitrary perturbations (3) and (4). It turns out that the dipole term ($B_\varphi \sim 1/r^3$) receives contributions only from the axially symmetrical perturbations ($m = 0$ and 1). Contributions to terms of higher order are made by perturbations of also high orders harmonics in φ.

By way of illustration we present a calculation of the field outside the target in the case when the perturbations are axially symmetrical (formulas (12)–(14)). In this case the field has only a φ component (we present below only the dipole term)

$$B_\varphi = \begin{cases} B_\varphi^t \dfrac{2\Delta R R^2}{r^3} \dfrac{1}{l(l^2-4)} \sin\theta, & l \text{ odd} \quad (l \geqslant 3) \\ 0, & l \text{ even} \end{cases} \tag{15}$$

(see Fig. 3). Here B_φ^t is the field amplitude in the target (13), r and θ are the distance and direction from the target to the measurement point; ΔR and R are the width of the zone in which the perturbations differ from zero and the target radius ($r \gg R$). In the calculation of terms of higher order in R/r, contributions are also obtained, naturally, from the even harmonics. Consequently, even at relatively short distances from the target (~10^{-1} cm) the fields are small ($\lesssim 1$ G) and may be difficult to measure in experiment.

The fact that $B_\varphi \neq 0$ only for odd harmonics can be understood by turning to the picture of the currents in the target for this case (Fig. 4). In fact, for even harmonics there exist equal numbers of oppositely directed currents, so that the fields produced by these currents cancel each other. To the contrary, for odd harmonics there remain uncancelled axial currents, and it is these that produce the resultant φ component of the field (Fig. 4). Comparing (13) and (15), we can note that the field in the target is stronger the larger the number of the harmonic, whereas outside the target the field for the higher harmonics of the perturbation is greatly decreased.

Formula (15) was obtained without allowance for the

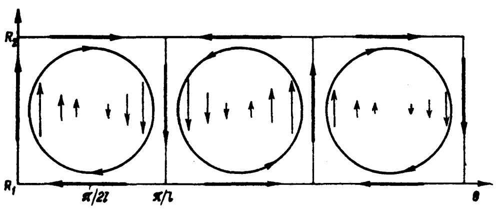

FIG. 2. Angular distribution of the currents in the instability zone for axially symmetrical perturbations.

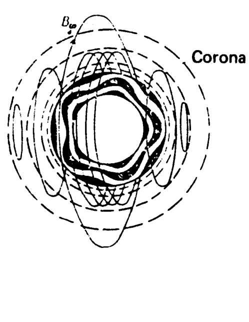

FIG. 3. Magnetic field at large distances from the target.

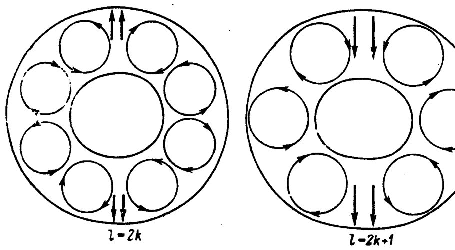

FIG. 4. Distribution of currents in instability zone in the cases when the number of harmonics is even and odd.

effect of the removal of the field by the plasma. Estimates show that the condition that the field be frozen-in is satisfied up to distances on the order of several hundred microns. Allowance for convection and diffusion of the field is in the general case a complicated matter. It follows, however, even from simple estimates that at the indicated distances the main effect is the dragging of the field by the plasma, and the diffusion is small. Under these conditions the field decreases in proportion to the density, $B\sim\rho\sim 1/r^{2-3}$, since $\rho \sim 1/vr^2$.

Thus, even at moderate laser-radiation fluxes there are present in the target appreciable magnetic fields that decrease sufficiently rapidly with increasing distance from the target.

We did not consider in this paper the possibility of field generation on the internal side of the shell, where Rayleigh-Taylor instability likewise arises at the instant of stopping. The produced fields can be strongly amplified by the compression effect. This question will be dealt with separately.

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