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Strong Magnetic Fields Produced by Composition Discontinuities in Laser-Produced Plasmas*

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Magnetic fields in excess of 10^7 G are expected to arise in laser-produced plasmas in which a discontinuity in atomic weight of the target occurs in the focal spot. The upper limit to B is controlled by fluid disassembly of the field layer due to magnetic pressure. This may provide a means for direct measurement of these fields in laser-produced plasmas.

The spontaneous generation of magnetic fields of order 10^6 G has been reported to occur[1-7] in high-density kilovolt plasmas produced by focusing a Nd laser pulse onto a solid target. Here we show that fields in excess of 10^7 G may be produced in such plasmas if a discontinuity in atomic weight occurs in the focal spot. Such high fields should be directly detectable via the Zeeman splittings of x-ray lines emitted from the plasma.

The magnetic field equation[5] is

$$\frac{\partial \vec{B}}{\partial t} = \nabla \times (\vec{V} \times \vec{B}) - \frac{c^2}{4\pi} \nabla \times [\overleftrightarrow{r} \cdot (\nabla \times \vec{B})] - \frac{c}{4\pi e} \nabla \times \left[\frac{1}{N_e} (\nabla \times B) \times \vec{B}\right] + \vec{S}, \quad (1)$$

where $\overleftrightarrow{r}$ is the resistivity and the source $\vec{S}$ has a thermal and a radiation-pressure part,

$$\vec{S} = -\frac{ck}{eN_e} \nabla N_e \times \nabla T_e - \frac{c}{eN_e^2} \nabla N_e \times (\nabla \cdot \overleftrightarrow{P}_R), \quad (2)$$

with the radiation pressure tensor given by[5,6]

$$\overleftrightarrow{P}_R = -\frac{1}{4\pi} \langle \epsilon \delta\vec{E} \delta\vec{E} + \delta\vec{B} \delta\vec{B} \rangle + \overleftrightarrow{I} \frac{1}{8\pi} \langle \delta\vec{E}^2 + \delta\vec{B}^2 \rangle, \quad (3)$$

where $\epsilon = 1 - \omega_e^2/\omega^2$, and $\delta\vec{E}$ and $\delta\vec{B}$ are the high-frequency electric and magnetic fields of the laser pulse. Now if a discontinuity in target material is arranged, such as in the sandwich target shown in Fig. 1, a very large electron density gradient, ∇N_e, is produced at the composition interface since the average ion charge Z ($N_e \cong ZN_i$) will in general differ in the two materials. This gradient is along the z axis, whereas ∇T_e and $\nabla \cdot \overleftrightarrow{P}_R$ are approximately along the beam direction at the center of the focal spot. This generates a large source of B field along the y direction.

Since the interface source layer is very thin (typically less then 1 μm), field diffusion out of the layer initially competes with the source S in Eq. (1). Field diffusion only slows down to speeds less than the fluid disassembly velocity V when the B field has spread through a layer of thickness of order 10 μm. We calculate the field values reached during the diffusion-dominated phase.

Thus, Eq. (1) can be solved if we (i) "freeze" the hydrodynamics, (ii) neglect the x dependence

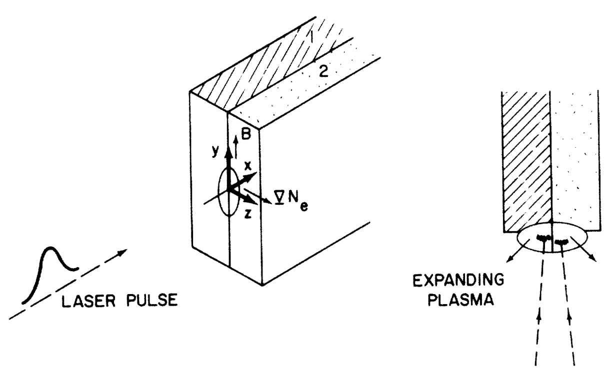

FIG. 1. Field produced by a focused laser pulse incident on a composition discontinuity. Although the critical depths for maximum energy deposition occur at different positions on the two sides, the displacement between them can be made less than the x temperature gradient scale for two materials not too widely spaced in atomic weight.

of B compared to its much faster z dependence, and (iii) neglect the t and z dependence of T and σ, and the nonlinear terms in B (these latter are of order $\beta^{-1} \leqslant 1$). For the geometry of Fig. 1, Eq. (1) then becomes

$$\frac{\partial B}{\partial t} \cong \frac{c^2}{4\pi\sigma} \frac{\partial^2 B}{\partial z^2} + S(z), \tag{4}$$

$$S = -\frac{c}{eN_e} \frac{\partial N_e}{\partial z}\left[\frac{\partial k T_e}{\partial x} + \frac{1}{N_e}(\nabla \cdot \overleftrightarrow{\mathrm{P}}_R)\cdot \hat{i}_x\right] \tag{5}$$

which represents the localized source together with diffusion away from the interface due to the finite conductivity σ ($\equiv \sigma_\perp$).

The electron density variation N_e is modeled by choosing

$$\frac{1}{N_e}\frac{\partial N_e}{\partial z} = \left|\frac{\Delta N_e}{N_{e0}}\right| \frac{1}{l_N} \exp\left(-\frac{z^2}{l_N^2}\right). \tag{6}$$

Similarly, a local temperature-gradient scale in the z direction is defined by

$$l_T = T_e/(\partial T_e/\partial x). \tag{7}$$

Note that l_T and l_N can be either positive or negative, and usually $|l_T| \gg |l_N|$. The solution to (4) with $B(t=0)=0$ then follows as

$$B = -2\tau_B\left(\frac{ckT_e}{el_N l_T}\right)\left|\frac{\Delta N_e}{N_{e0}}\right|\left[1 + Kl_T\frac{\mathcal{E}_{\mathrm{rad}}}{N_e k T_e}\right]$$
$$\times\left\{\left(1+\frac{t}{\tau_B}\right)^{1/2}\exp\left(-\frac{z^2}{l_N^2(1+t/\tau_B)}\right) - \exp\left(-\frac{z^2}{l_N^2}\right) + \frac{z\pi^{1/2}}{l_N}\left[\Phi\left(\frac{z}{l_N(1+t/\tau_B)^{1/2}}\right) - \Phi\left(\frac{z}{l_N}\right)\right]\right\}, \tag{8}$$

where

$$\tau_B = \pi\sigma l_N^2/c \tag{9}$$

is characteristic field diffusion time, the strong-field ($\Omega_e\tau_e \gg 1$) perpendicular conductivity σ is[8]

$$\sigma = 7.2\times 10^7 T_e^{3/2}/Z\ln\Lambda \tag{10}$$

with T_e in degrees Kelvin, and $\Phi(y) = 2\pi^{-1/2}\int_0^y dx\,\exp(-x^2)$. The radiation-pressure part of (8) was written[5]

$$\hat{i}_x\cdot(\nabla\cdot\overleftrightarrow{\mathrm{P}}_R) \cong K\mathcal{E}_{\mathrm{rad}}, \tag{11}$$

where $\mathcal{E}_{\mathrm{rad}}$ is the radiation energy density in the focal spot and K the absorption coefficient.[9] Radiation reflection[5] also contributes to (11) but has been neglected.

The solution (8) has features shown in Fig. 2. The field diffuses a distance

$$z_B \cong l_N(t/\tau_B)^{1/2} = c(t/\pi\sigma)^{1/2} \tag{12}$$

away from the interface in a time t after the laser pulse is initiated. The value of B at the origin ($z=0$), and for times $t \gg \tau_B$, becomes

$$B_0 \cong \frac{8.7\times 10^{-3} T_e^{7/4} t^{1/2}}{l_T(Z\ln\Lambda)^{1/2}}\left|\frac{\Delta N_e}{N_{e0}}\right|\left(1 + Kl_T\frac{\mathcal{E}_{\mathrm{rad}}}{N_e k T_e}\right). \tag{13}$$

After some time the field (13) can grow to a magnitude such that $\beta \cong 1$. If this occurs the field layer then acts as a piston and drives a shock into the adjacent plasma. Fluid expansion of the B layer then follows on a time scale $\sim z_B/V_A$ and competes with further growth in B. Only very short pulses (for which $\tau_L < z_B/V_A$) could give rise to fields in excess of this $\beta=1$ limit, i.e., for most cases the maximum B field is

$$B_M \cong 6\times 10^{-8}(N_e T_e + N_i T_i)^{1/2}. \tag{14}$$

Now typically we are interested in Nd pulses of energy > 50 J, focal spot radii < 50 μm, and

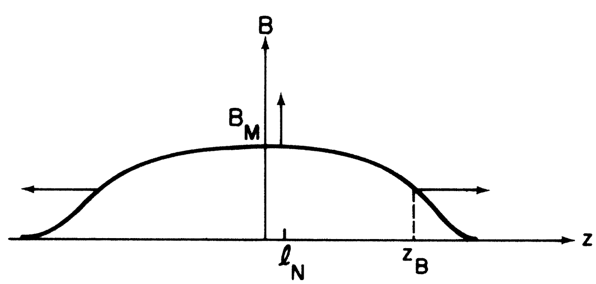

FIG. 2. Diffusion of magnetic flux away from its source near the origin.

pulse times ≤ 1 nsec incident on a high atomic weight target.[3] For $T_e \sim 5\times 10^7$ °K, $Z\sim 20$, $\ln\Lambda \sim 10$, the field diffuses a distance $z_B \sim 45\sqrt{t}$ cm, e.g., $z_B = 4.5$ μm for $t = 10^{-10}$ sec.

Very little diffusive broadening of the composition jump occurs. In a time t the interface broadens via ion diffusion to a thickness l_N given by $\lambda_i(t/\tau_i)^{1/2}$, where the scattering time τ_i for ions of charge Ze on side 1 by ions of charge $(Z+\delta Z)e$ on side 2 is[8] (for the simplest case $\delta Z \ll Z$)

$$\tau_i \cong \frac{17}{\ln\Lambda}\left(\frac{m_i}{m_p}\right)^{1/2}\frac{T_i^{3/2}}{N_i Z^4}. \quad (15)$$

For example, if $T_i = 10^7$ °K ($T_i < T_e$), $ZN_i = N_e = 10^{21}$ cm^{-3}, $Z\sim 20$, $\ln\Lambda \sim 10$, $m_i/m_p \sim 10^2$, then for $t = 0.3$ nsec, $l_N \cong 0.2$ μm.

For a jump with gradient scale $l_N \sim 1$ μm (note also $l_N \gg \lambda_D \sim 10^{-6}$ cm), the field diffusion time τ_B [Eq. (9)] is 4×10^{-12} sec, so that the large-t formula (13) applies for B_0 throughout most of the pulse duration. For example, if $T_e = 5\times 10^7$, $l_T = 4\times 10^{-3}$, $Z = 20$, $\ln\Lambda \sim 10$, we find

$$B_0 \cong 4.6\times 10^{12} t^{1/2}\left(1 + \frac{K l_T \mathcal{E}_{\rm rad}}{N_e k T_e}\right)\left|\frac{\Delta N_e}{N_{e0}}\right|, \quad (16)$$

$$B_M = 4.2\times 10^{-4} N_e^{1/2}.$$

The field B_0 reaches its maximum value B_M in a time

$$t_M \cong 10^{-32} N_e \left(1 + \frac{K l_T \mathcal{E}_R}{N_e k T_e}\right)^{-2}\left|\frac{\Delta N_e}{N_{e0}}\right|^{-2}. \quad (17)$$

Since for most cases T_e increases to its maximum on the same time scale as the laser pulse width, τ_L, an optimum situation involves matching τ_L with t_M. For a sample case $|\Delta N_e/N_{e0}|^2 = 0.1$ and $\mathcal{E}_R \cong 0$, these formulas show that B_0 would rise to a value $B_M \cong 1.3\times 10^7$ G in $t_M = 0.1$ nsec at the critical depth where $N_e = 10^{21}$ cm^{-3}. However, heating also penetrates into the overdense regions beyond the critical depth for laser energy deposition. If similar temperatures occur in the region $N_e = 10^{22}$, B_M becomes 4.2×10^7 G which is reached in $t_M = 1$ nsec. For such fields, Zeeman splitting of x-ray lines is about equal to their Doppler broadening [$\hbar\Omega_e = (1.2\times 10^{-8}$ eV 1 G$)B$], and should become detectable. Emission from the region of maximum $B(z)$, i.e., B_0, could be selectively observed by seeding the interface with a small amount of material, e.g., Al, for which hydrogenic and He lines have been identified. Smearing of the splitting would then arise only from the x and t dependence of B_0. However the effect may still appear as a splitting instead of a broadening in time-integrated profiles since the x-ray lines are principally emitted when T (and therefore B) is large.

Finally we note that the question of electrostatic instability and anomalous conductivity arises for the current layer associated with the B field. However, for highly charged ions [use of parameters below Eq. (15)] the ion collision frequency τ_i^{-1} is comparable to the ion plasma frequency at the critical depth [$\omega_i = \omega_e(Zm_e/m_i)^{1/2} \sim 2\times 10^{13}$]. Thus, ion sound instability is collisionally damped and the classical conductivity (10) is expected to apply.

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