Saturation Mechanisms for the Generated Magnetic Field in Nonuniform Laser-Matter Irradiation

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Spontaneously generated magnetic fields B can be generated by nonuniform laser irradiation causing sources such as $\nabla T_e \times \nabla n_e / n_e e$. By considering convective losses and dissipation including the triggering of lower hybrid drift turbulence a universal diagram of saturated B versus $L_\perp$, the characteristic transverse scale length for the nonuniformity, is found. For $L_\perp$ less than the ion collisionless skin depth c/ω_{pi} the saturated B scales as $L_\perp^\alpha$ $\alpha \leq 1$. At higher $L_\perp$ values B scales as $L_\perp^{-1}$ because convective losses dominate. For resistive plasmas (e.g., high Z, low temperature) B will be limited to $(\mu_0 k T_e / e\eta)(L_\perp / L_\parallel)$ where $L_\parallel$ is $n_e/\nabla n_e$. [S0031-9007(96)01987-4]

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An important feature of any departure from one-dimensional symmetry in laser-plasma interactions and in inertial confinement in general is the possibility of the spontaneous generation of magnetic fields. Such fields can be in the several megagauss range [1–4] and are even predicted to be in the 10^8 gauss range with intense short pulse lasers [5–7]. Such magnetic fields can strongly affect energy transport leading to hot spots [4], fast electrons, and fast ions [8], as well as effectively freezing in initial laser imprints which can seed the Rayleigh-Taylor instability, and even producing a significant magnetic pressure.

The source is principally due to a localized supply of energetic electrons, thermal in origin for low intensity long pulse lasers, and very suprathermal, indeed relativistic, for short, intense laser pulses, to which must be added the electron quiver energy in the region of laser interaction. These energetic electrons will try to leave the heated region. In one-dimensional symmetry electrostatic forces arising from the breakdown of exact electrostatic neutrality will return the electrons, such fields also causing a slower acceleration of the ions, leading to a quasineutral expansion or ablation.

In the absence of one-dimensional symmetry the electrons will not return along the same path, and a net circulating current will flow. The resulting magnetic field B that grows will induce a back EMF to oppose this current flow. The fact that the electric field is no longer purely the electrostatic $-\nabla\Phi$ field alone but has in addition an induced $-\partial \mathbf{A}/\partial t$ term is useful in deriving the rate of generation of $\mathbf{B}$. This is done by combining Faraday's law

$$\frac{\partial \mathbf{B}}{\partial t} = -\nabla \times \mathbf{E} \tag{1}$$

with Ohm's law

$$-\mathbf{E} = \boldsymbol{\nu} \times \mathbf{B} - \frac{\mathbf{J} \times \mathbf{B}}{n_e e} + \frac{\mathbf{q}_e \times \mathbf{B}}{\frac{5}{2} p_e} + \frac{\nabla p_e}{n_e e} + \beta \frac{\nabla T}{e} - \eta \mathbf{J} - \frac{m_e}{n_e e^2} \frac{\partial J}{\partial t} + \frac{1}{n_e e} \nabla \cdot (\boldsymbol{T}_e + n_e m_e \boldsymbol{\nu}_e \boldsymbol{\nu}_e + \langle n_e m_e \dot{\boldsymbol{\xi}} \dot{\boldsymbol{\xi}} \rangle) + \langle \dot{\boldsymbol{\xi}} \times \tilde{\mathbf{B}} \rangle, \tag{2}$$

where inclusion of the term in the electron heat flux q_e allows the simplification [9] of the otherwise tensor resistivity and thermoelectric terms of Braginskii [10] used here in corrected form [11]. We include the electron traceless stress tensor $\boldsymbol{T}_e$ and the ponderomotive stress associated with the quiver velocity $\dot{\boldsymbol{\xi}}$ which arises from the high frequency Lorentz force associated with the laser fields $\tilde{\mathbf{E}}$ and $\tilde{\mathbf{B}}$,

$$m_e \ddot{\boldsymbol{\xi}} = -e(\tilde{\mathbf{E}} + \dot{\boldsymbol{\xi}} \times \tilde{\mathbf{B}}). \tag{3}$$

The time average $\langle\rangle$ of the last term leads to the conventional radiation pressure given in the last term in Eq. (2).

By inserting the curl of Eq. (2) into Eq. (1) we can identify all the sources of magnetic field, the convective loss terms, and the dissipative loss terms. This was undertaken in an earlier review paper [12]. The new features presented here are (i) identifying and classifying into various regimes the dominant terms in Eq. (2) that can grow and lead to the saturation of the magnetic field amplitude and (ii) including the possibility of the triggering of lower hybrid drift turbulence at high current density, which will be an additional saturation mechanism for the magnetic field. It will be found that microturbulence can be dominant for transverse scale lengths less than the ion collisionless skin depth.

The source term of greatest interest arises from the ∇p_e term in Eq. (2). On taking the curl of just this term and inserting into Eq. (1) we obtain the well known result

$$\frac{\partial \mathbf{B}}{\partial t} = \frac{\nabla T_e \times \nabla n_e}{n_e e}, \tag{4}$$

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which demonstrates that if the electron temperature and density gradients are not parallel, a spontaneous generation of magnetic field occurs. Usually experiments and modeling are conducted with azimuthal symmetry with a focused laser beam propagating in the $\hat{\mathbf{z}}$ direction with maximum intensity on the axis (at $r = 0$) and interacting with a planar solid surface normal to the axis [1–4]. The dominant component of ∇n_e is $\partial n_e/\partial z \hat{\mathbf{z}}$ the magnitude of which we denote by $n_e/L_\parallel$, and the magnitude of the relevant transverse component of ∇T_e, i.e., $\partial T_e/\partial r \hat{\mathbf{r}}$ we denote by $T_e/L_\perp$. If the ionic Z charge varies spatially then a genuine thermoelectric emf can arise with $\partial B/\partial t$ equal to $\nabla\beta \times \nabla T_e/e$. Reference [9] gives $\beta(Z, \Omega_e\tau_e)$ such that β varies from 0.703 ($Z = 1$) to 1.5 ($Z = \infty$) for $\Omega_e\tau_{ei} = 0$.

Other source terms particularly important at high laser intensity are the radiation pressure term and the ponderomotive force, if they are nonuniform in the transverse direction [6,13–15]. It is difficult in an experiment or a full numerical simulation to distinguish the magnetic field generated by these source terms and by the electron pressure or stress tensor terms which soon grow to be equal in magnitude [5,7].

Magnetic fields have also been predicted to occur in spatially uniform laser irradiation either through resonance absorption at oblique incidence [16] or through instabilities. The instabilities are caused by nonlinear heat flow and the associated anisotropy in the distribution function; this then essentially drives off-diagonal terms in the electron stress tensor either through the collisionless Weibel instability [17–19], the collisional Weibel instability [20,21], or the thermal instability [22], depending on the collisionality. For the last case this occurs in regions of heat flow where the mean free path of the electrons, λ_{mfp}, is less than the electron collisionless skin depth c/ω_{pe}. (A thermomagnetic instability theory [23] derived without proceeding to the stress tensor is incomplete and is subsumed in Refs. [20] and [21].)

In what follows we will take the $\nabla p_e/n_e e$ term in Eq. (2) as the archetypal source term, it being relatively easy to replace the electron pressure by the laser radiation pressure or ponderomotive stress as appropriate for a particular situation.

The resistive term $\eta\mathbf{J}$ is the main dissipative term, and we can identify regime 1 in which the saturation of the magnetic field occurs when this term balances $\nabla p_e/n_e e$. For classical resistivity and employing Ampère's law

$$\nabla \times \mathbf{B} = \mu_0 \mathbf{J} \tag{5}$$

we arrive at the result [24]

$$B = \frac{\mu_0 k T_e}{e\eta} \frac{L_\perp}{L_\parallel}. \tag{6}$$

The value of the Hall parameter $\Omega_e\tau_{ei}$ to which this corresponds is

$$\Omega_e\tau_{ei} \cong \frac{\lambda_{mfp}^2}{(c/\omega_{pe})^2} \frac{L_\perp}{L_\parallel}, \quad \propto \frac{T_e^4}{n_e z^2} \frac{L_\perp}{L_\parallel}, \tag{7}$$

where λ_{mfp} is the electron mean-free-path.

However, classical resistivity is valid only for low current densities. If through some triggered microturbulence the electron drift velocity is limited to, say, the ion sound speed $c_s = (ZkT_e/m_i)^{1/2}$, a very different result pertains; inserting $|\mathbf{J}| = n_e e c_s$ into Eq. (5) yields

$$B = \mu_0 n_e e \left(\frac{ZkT_e}{m_i}\right)^{1/2} L_\perp. \tag{8}$$

We label this as regime 2.

A more specific result for this regime can be obtained if we assume that the most probable turbulence arises from the lower hybrid drift instability since the current density and magnetic field are mutually orthogonal. At first, when there is no magnetic field, the ion-acoustic or even the two-stream instability might be triggered, but as the magnetic field grows, it is almost certain that the lower hybrid mode will dominate. The linear growth rate for this instability has been studied by various authors with a local approximation for electrostatic modes with unmagnetized ions [25,26], and later the inclusion of electromagnetic modes, finite β, and ∇B electron drifts [27,28]. More recent work includes magnetized ions and ∇B ion drift [29]. Comparison with a nonlocal simulation [30] shows that the fastest mode has a growth rate of $0.5(\Omega_e\Omega_i)^{1/2}$ at an optimum wave number of $(\Omega_e\Omega_i)^{1/2}/v_i$ where v_i is the ion thermal speed $(2kT_i/m_i)^{1/2}$ in the regime $\omega_{pe} > \Omega_e$. In the regime $\omega_{pe} < \Omega_e$ we expect the growth rate to be characterized by the ion plasma frequency ω_{pi} and the optimum wave number by the reciprocal of the Debye length. In applying the results of lower hybrid turbulence, care must be taken to ensure that the characteristic times and wavelengths of the instability are much less than the laser pulse length and inhomogeneity length scale.

Saturation of the instability has been followed in an electromagnetic implicit particle code [30]; at early times $(1000\omega_{pe}^{-1})$ the anomalous resistivity is three times that at late times $(3000\omega_{pe}^{-1})$. Taking the latter value in the regime $\omega_{pe} > \Omega_e$, the numerical simulation gives an effective collision frequency for lower hybrid turbulence, ν_{LH},

$$\nu_{\mathrm{LH}} = \frac{0.02}{\beta_i} (\Omega_e\Omega_i)^{1/2} \left(\frac{v_d}{v_i}\right)^2, \tag{9}$$

where $\beta_i = 2\mu_0 n_i m_i v_i^2/B^2$ and $v_d = \mathbf{J}/n_e e$. If we employ this in an anomalous resistivity η_{LH} and balance $\eta_{\mathrm{LH}} J$ against $\nabla p_e/n_e e$ in Eq. (2) and again use Eq. (5), we find that Eq. (8) has to be modified to $B \propto L_\perp^{1/3}(L_\perp/L_\parallel)^{1/6}$. [(If we are in the regime $\omega_{pe} < \Omega_e$ and if it is appropriate to replace $(\Omega_e\Omega_i)^{1/2}$ in Eq. (9) by ω_{pi} we find that $B \propto L_\perp^{2/5}(L_\perp/L_\parallel)^{1/5}$.)]

There is another model of lower hybrid drift wave saturation based on nonlinear mode coupling [31]. The result of this purely electrostatic model is to give an effective collision frequency ν_{LH} given by

$$\nu_{\mathrm{LH}} = 2.4(\Omega_e\Omega_i)^{1/2} \left(\frac{v_d}{v_i}\right)^2 \tag{10}$$

which, when balancing the anomalous $\eta_{\mathrm{LH}} J$ with $\nabla p_e / n_e e$, leads to the saturated $B \propto L_\perp^{1/2} (L_\perp / L_\parallel)^{1/4}$ in the $\omega_{pe} > \Omega_e$ regime. [In the $\omega_{pe} < \Omega_e$ regime, replacing $(\Omega_e \Omega_i)^{1/2}$ by ω_{pi} leads to $B \propto L_\perp^{2/3} (L_\perp / L_\parallel)^{1/3}$.]

There are three convective terms in Eq. (2) that can act as a loss to balance the source term. The first of these is the $\boldsymbol{\nu} \times \mathbf{B}$ term. To estimate the ablation velocity $\boldsymbol{\nu}$ we note from the fluid equation of motion

$$\rho \left[\frac{\partial \boldsymbol{\nu}}{\partial t} + (\boldsymbol{\nu} \cdot \nabla) \boldsymbol{\nu} \right] = -\nabla p \tag{11}$$

that the flow in the vicinity of the critical surface will be sonic with respect to c_s, i.e., $\nu = c_s$. On balancing $\nabla \times (\nu \times B)$ with the $\nabla T_e \times \nabla n_e / n_e e$ source we arrive at the result [8,32] for regime 3,

$$B = \frac{1}{L_\perp} \left(\frac{m_i k T_e}{Z e^2} \right)^{1/2}. \tag{12}$$

The $B \propto L_\perp^{-1}$ variation clearly indicates that this is the saturation mechanism for large $L_\perp$. If the ablation velocity is reduced by a factor of $\alpha < 1$ due, for example, to the gas fill in a hohlraum the magnetic field is increased by α^{-1}.

At this point we can combine the scaling laws found in Eqs. (8) and (12), where B is proportional to $L_\perp$ and to $1/L_\perp$, respectively. The intersection of these two scaling laws occurs at $L_\perp = c/\omega_{pi}$, the ion collisionless skin depth; at this point the magnetic field B is $(\mu_0 n_e k T_e)^{1/2}$, i.e., the magnetic and electron pressures are almost equal and the Hall parameter $\Omega_e \tau_{ei}$ at this point is

$$\Omega_e \tau_{ei} = \frac{\lambda_{mfp}}{c/\omega_{pe}}. \tag{13}$$

Figure 1 illustrates these two scaling laws. We can also include the saturated B for classical resistive diffusion given by Eq. (6); for $L_\perp = L_\parallel$, B is independent of $L_\perp$ and depends only on temperature and ionic charge number Z. Three lines are shown, corresponding to (a) high temperature, low Z or more precisely $\Omega_e \tau_{ei} > 1$, (b) the triple point where $\Omega_e \tau_{ei} = 1$, and (c) low temperature, low Z, i.e., $\Omega_e \tau_{ei} < 1$. In case (c) the saturated magnetic field is determined by the lower envelope of the three equations (8), (6), and (12) covering regimes 2, 1, and 3, respectively. In case (c) at $L_\perp = c/\omega_{pi}$ we have $\lambda_{mfp} < c/\omega_{pe}$, which is the necessary condition for the triggering of a thermal instability [22].

Strictly we should replace Eq. (8) for regime 2 by an equation employing the anomalous collision frequency ν_{LH} given by Eqs. (9) and (10). The intersection of the $B \propto L_\perp^{1/3}$ curve of Eq. (9) with the $L_\perp^{-1}$ curve of Eq. (12) occurs at

$$L_\perp = \frac{c}{\omega_{pi}} \left(\frac{L_\parallel}{L_\perp} \right)^{1/8} \left(\frac{Z m_e}{m_i} \right)^{1/16} \left(\frac{Z T_e}{T_i} \right)^{1/4} \left(\frac{0.01}{4} \right)^{1/8}, \tag{14}$$

which for $L_\parallel = L_\perp$, $Z m_e / m_i = 1/3672$, and (a) $T_e = 10 T_i$; and $Z = 10$ gives $L_\perp = 0.90 c/\omega_{pi}$ and for (b) $T_e = T_i$ and $Z = 1$ gives $L_\perp = 0.28 c/\omega_{pi}$. Similarly for the formula

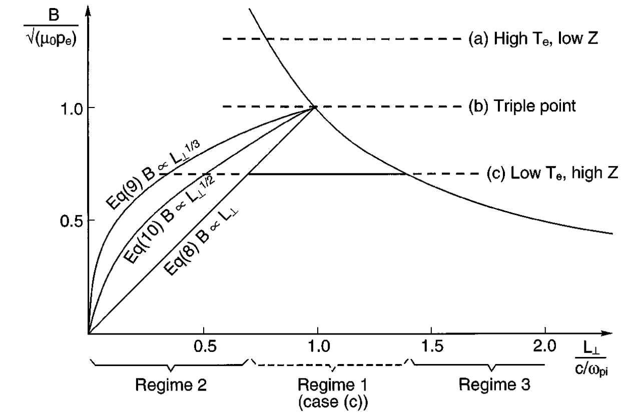

FIG. 1. Universal diagram of the saturated magnetic field B versus transverse length $L_\perp$ (in dimensionless units).

of Eq. (10) the $B \propto L_\perp^{1/2}$ curve intersects Eq. (12) at

$$L_\perp = \frac{c}{\omega_{pi}} \left(\frac{L_\parallel}{L_\perp} \right)^{1/6} \left(\frac{Z m_e}{m_i} \right)^{1/12} \left(\frac{Z T_e}{T_i} \right)^{1/6} 1.2^{1/6}. \tag{15}$$

Again for $L_\parallel = L_\perp$, $Z m_e / m_i = 1/3672$ and (a) $T_e = 10 T_i$ and $Z = 10$ gives $L_\perp = 1.09 c/\omega_{pi}$ and for (b) $T_e = T_i$ and $Z = 1$ gives $L_\perp = 0.50 c/\omega_{pi}$. We conclude that the intersection points are approximately coincident, especially for $T_e \gg T_i$ and high Z. The curves corresponding to Eqs. (9) and (10) for $\omega_{pe} > \Omega_e$ are included in Fig. 1 for case (a) above.

The second convection term is the Hall effect. Since the current density $\mathbf{J}$ obeys $\nabla \cdot \mathbf{J} = 0$ and the $\mathbf{J}$ streamlines are closed in the poloidal plane, the magnetic field is essentially rotated around by this term. Its inclusion in simulations [33] can modify the field distribution, but since $\mathbf{J}/n_e e$ should be limited by microturbulence to values of order c_s it is not likely to lead to a new scaling regime.

The third convection term, $\mathbf{q}_e \times \mathbf{B} / (\frac{5}{2} p_e)$, is more accurately written [9] as $\beta_\wedge \mathbf{q}_e \times \mathbf{b} / e \kappa_\perp$, where $\mathbf{b}$ is the unit vector in the magnetic field direction. If the heat flux $\mathbf{q}_e$ is dominated by the term $-\kappa_\perp \nabla T$, the ratio of this convective term to $\nabla p_e / n_e e$ is the dimensionless quantity $\beta_\wedge (L_\parallel / L_\perp)$. $\beta_\wedge$ is a function of $\Omega_e \tau_{ei}$ and has a maximum value of 0.285 at $\Omega_e \tau_{ei}$ of 0.89 for $Z = 1$, and of 0.491 at $\Omega_e \tau_{ei}$ of 0.23 for $Z = \infty$ [11]. The Nernst term is therefore significant but requires a value of $L_\parallel / L_\perp$ of 2 to 4 (for $Z = \infty$ to 1) in order to balance the source term. Since for inertial confinement the heat flux inwards towards the ablation surface must exceed the outward advected enthalpy flow $\frac{5}{2} p_e \nu \cong \frac{5}{2} p_e c_s$ it is clear that in general the heat flow is nonlinear and can convect the magnetic field both inwards towards the ablation surface as well as outwards in the corona outside the critical surface so long as there are sufficient collisions for the theory to hold. Indeed the inward convection can theoretically lead to considerable amplification of the magnetic field [34,35], a phenomenon confirmed in nonlinear heat flow using Fokker-Planck calculations [36,37]. However, because the expected heat flux while exceeding $\frac{5}{2} p_e c_s$ will scale in a similar way the scaling of the saturation value of B must be similar to Eq. (8), i.e., regime 2.

The terms involving electron inertia in Eq. (2) allow penetration of B to an electron collisionless skin depth c/ω_{pe}, this being manifested as a radially propagating surface layer in a PIC simulation [38], and in a reversed magnetic field in the overdense region in an analytic relativistic case [6]. The apparent discrepancy between Refs. [5] and [6] might be attributed to a kinetic broadening over a time t of the region carrying the "return" current; the anomalous skin effect formula [39,40] $(c^2 \nu_{Te} \pi t/\omega_{pe}^2)^{1/3}$ fits the simulation [5] well. Here $(\pi t)^{-1}$ has replaced the em frequency ω of Ref. [40]. These fast electrons can be energized by resonant absorption [41] or by "vacuum" heating proposed by Brunel [42] due to p-polarized light [7].

In summary, we have shown (see Fig. 1) that the spontaneously generated magnetic field can be limited by lower hybrid turbulence for transverse inhomogeneity scale-lengths less than c/ω_{pi}, scaling as $L_\perp^\alpha$ where α is $\frac{1}{3}$ to 1 (regime 2) depending on the model used; for $L_\perp > c/\omega_{pi}$ the magnetic field is determined by convection (regime 3), scaling as $L_\perp^{-1}$. At the intersection of these regimes the magnetic pressure can equal the plasma pressure, and at this point $\Omega_e \tau_e = \lambda_{mfp}/(c/\omega_{pe})$ provided that $\Omega_e \tau_e > 1$. For a colder or high Z plasma, dissipation can dominate, leading to regime 1 in which the magnetic field is limited to $(\mu_0 k T_e/e\eta)\ (L_\perp/L_\parallel)$. We have limited our consideration to the generalized Ohm's law and have not explored the coupling to the electron energy equation or indeed carried out full simulations over the range of physics considered in this Letter, which would be very interesting. Thus the estimate of $L_\perp$ for a given inhomogeneity in laser irradiation, and the subsequent importance or otherwise of magnetic fields in laser fusion targets requires further study. But this model gives, for example, in regime 3 fields of several magagauss for $L_\perp$ of $2c/\omega_{pi}$($\approx$ 20 laser wavelengths at critical density) for $T_e = 1$ keV which will perturb pressure and transport.

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