Parametric generation of magnetic fields by action of strong radiation on a plasma

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(Submitted 4 December 1978)
Zh. Eksp. Teor. Fiz. **76**, 1586–1592 (May 1979)

We investigate the generation of a magnetic field in a plasma acted upon by powerful radiation. The generation mechanism is taken to be parametric excitation of the zeroth harmonic of a nonpotential field; this harmonic increases as a result of parametric resonance of the pump wave with the Langmuir oscillations. The estimate obtained for the nonlinear saturation level of the magnetic field points to a possibility of generating megagauss fields in a laser plasma.

PACS numbers: 52.35.Mw, 52.50.Jm

1. The generation of magnetic fields in a plasma acted upon by electromagnetic radiation are being intensively investigated of late. This phenomenon was first discussed in Refs. 1 and 2. The main cause of the urgency of the investigation of the generation of magnetic fields are the experiments performed in the program of laser-mediated controlled fusion, primarily for the reason that the produced strong magnetic fields can greatly influence the character of the penetration of the laser radiation into the plasma and the rate of the transport processes that take place in it. It must also be borne in mind that a strong magnetic field can play a significant role in the absorption, scattering, and depolarization of the electromagnetic radiation.

In the study of the question of the onset of magnetic fields in a laser plasma, the most widely discussed are the following two mechanisms. First, the magnetic field may be excited by thermoelectric power.[3,4] Second, the magnetic field can be generated by the p-component of the electric field in the laser radiation.[5,6] In Refs. 3–6 the onset of the magnetic field was attributed by the authors to a source that acts continuously in time. A more important role, however, is played by physical processes in which the magnetic field generation has a self-excitation character. An analysis of the process of excitation of a magnetic field from the point of view of the development of instability in a plasma was carried out in Refs. 7 and 8 for the case of a thermionic-current source of field generation and in Ref. 8 for a turbulent plasma.

A possibility of generating a field by development of a magnetic instability initiated by high-frequency radiation of a pump wave was observed in Ref. 10. The analysis was carried out in the approximation of strong collisional hydrodynamics under conditions when the electron mean free path was smaller than the wavelength of the perturbations, and the collision time was shorter than the characteristic time of field variation.

The present paper presents a further development of the concepts advanced in Ref. 10. The undertaken kinetic description of parametric generation of the zeroth harmonic of a nonpotential perturbation field has made it possible to analyze the linear stage of the instability in a wide range of wavelengths and frequencies. We estimate the nonlinear saturation level of the magnetic field.

2. We consider a homogeneous plasma subjected to the action of a high-frequency electric field

$$\mathbf{E}_0(t)=\mathbf{E}_0\sin\omega_0 t. \qquad (1)$$

The frequency ω_0 of the pump field will be assumed close to the plasma frequency ω_p [$\omega_p=(\omega_{Le}^2+\omega_{Li}^2)^{1/2}$, $\omega_{Le(i)}$ is the Langmuir frequency of the electrons (ions)].

Being interested in the possibility that the magnetic field will increase with time, we consider the onset of parametric instability accompanied by excitation of nonpotential perturbations in the pump-wave field (1).

Representing the nonequilibrium electric field $\delta\mathbf{E}$ in the form of an expansion

$$\delta\mathbf{E}(\mathbf{r},t)=e^{i\mathbf{kr}-i\omega t}\sum_{n=-\infty}^{\infty}\delta\mathbf{E}^{(n)}\exp\{-in\omega_0 t\}, \qquad (2)$$

we obtain for the components of the amplitudes of the expansion (2) from Maxwell's equations the following system of equations[11]

$$\begin{aligned}&\{(\omega+n\omega_0)^2\varkappa_l\varkappa_j[1+\delta\varepsilon_i(\omega+n\omega_0,\mathbf{k})]+[(\omega+n\omega_0)^2(1+\delta\varepsilon_i^{tr}(\omega+n\omega_0,\mathbf{k}))\\&-c^2k^2](\delta_{lj}-\varkappa_l\varkappa_j)\}\delta E_j^{(n)}+\sum_{m,r=-\infty}^{\infty}\frac{\omega+n\omega_0}{\omega+m\omega_0}\delta E_j^{(m)}J_{r-n}(\mathbf{kr}_E)J_{r-m}(\mathbf{kr}_E)\\&\times\{(\omega+r\omega_0)^2\delta\varepsilon_e^{tr}(\omega+r\omega_0,\mathbf{k})(\delta_{lj}-\varkappa_l\varkappa_j)+\delta\varepsilon_e(\omega+r\omega_0,\mathbf{k})\Big[\varkappa(\omega+n\omega_0)\\&+(n-r)\omega_0\frac{|[\mathbf{k}\times\mathbf{r}_E]|}{\mathbf{kr}_E}\chi\Big]_l\Big[\varkappa(\omega+m\omega_0)+(m-r)\omega_0\frac{|[\mathbf{k}\times\mathbf{r}_E]|}{\mathbf{kr}_E}\chi\Big]_j\Big\}=0.\end{aligned} \qquad (3)$$

Here $\delta\varepsilon_{e(i)}$ is the partial longitudinal dielectric constant of the electrons (ions), $\delta\varepsilon_{e(i)}^{tr}$ is the partial transverse dielectric constant of the electrons (ions), $\mathbf{r}_E=e\mathbf{E}_0/m_e\omega_0^2$ is the amplitude of the oscillations of the electrons in the pump-wave field (e and m_e are the charge and mass of the electrons), c is the speed of light, and $\varkappa$ and χ are mutually perpendicular unit vectors:

$$\varkappa=\frac{\mathbf{k}}{k},\quad \chi=\frac{[\mathbf{k}\times[\mathbf{E}_0\times\mathbf{k}]]}{k|[\mathbf{E}_0\times\mathbf{k}]|}, \qquad (4)$$

J_n is a Bessel function of the first kind with integer index n. Besides the system (3) we must also bear in mind the connection, which follows from Maxwell's equations, between the amplitudes of the electric and magnetic fields:

$$\delta\mathbf{B}^{(n)}=\frac{c}{\omega+n\omega_0}[\mathbf{k}\times\delta\mathbf{E}^{(n)}]. \qquad (5)$$

The system (3) will be used below to study the properties

0038-5646/79/050805-04$02.40

of low-frequency short-wave plasma perturbations, when $\omega_0 \gg |\omega|$ and $k \gg \omega_0/c$.

It follows from the system (3) that for oscillations whose electric vector $\delta\mathbf{E}$ is polarized perpendicular to the plane of the vectors $\mathbf{k}$ and $\mathbf{E}_0$ there is no stability.[11] To reveal the instability we consider therefore the case of a polarization wherein the electric vector of the perturbations lies in the plane of the vectors $\mathbf{k}$ and $\mathbf{E}_0$. For these oscillations, the amplitudes $\delta\mathbf{E}^{(n)}$ can be represented in the form of an expansion in the unit vectors (4)

$$\delta\mathbf{E}^{(n)}=\varkappa\delta E_l^{(n)}+\chi\delta E_{tr}^{(n)}, \qquad (6)$$

which corresponds to separation of the potential and nonpotential parts of the electric field of the perturbations.

We consider first the case of a perturbation propagation wherein $\theta > \theta_0$, where $\theta = \pi/2 - \alpha, \cos\alpha = \mathbf{k}\mathbf{r}_E/kr_E$, and

$$\theta_0=\frac{\omega_0}{ck}\left|\frac{1}{\varepsilon(\omega+\omega_0,k)}+\frac{1}{\varepsilon(\omega-\omega_0,k)}\right|^{-1/2}.$$

This angle range $\pi/2 > \theta > \theta_0$ covers practically all the propagation directions of the perturbations, with the exception of a narrow cone of directions for which the wave vector $\mathbf{k}$ is practically perpendicular to the external electric field.

Under the conditions $|\mathbf{k}\mathbf{r}_E| \ll 1$ we find from the system (3) that the following dispersion equation

$$\varepsilon(\omega,k)+\frac{1}{4}(\mathbf{k}\mathbf{r}_E)^2\delta\varepsilon_i(\omega,k)(1+\delta\varepsilon_e(\omega,k))\times\left(\frac{1}{\varepsilon(\omega+\omega_0,k)}+\frac{1}{\varepsilon(\omega-\omega_0,k)}\right)=0,$$

which coincides with that obtained in Ref. 12 if the nonpotential Langmuir oscillations are neglected, holds in the considered angle range. Here $\varepsilon = 1 + \delta\varepsilon_e + \delta\varepsilon_i$ is the longitudinal dielectric constant of the plasma. Thus, allowance for the nonpotential perturbations has little influence on the dispersion properties of the parametrically unstable plasma. The generation of a low-frequency magnetic field is the result of the parametric coupling of the harmonics $\delta E_l^{(n)}$ and $\delta E_{tr}^{(n)}$ via the pump field, a coupling that follows from the system (3). In particular, from these equations we get in accordance with (5) the approximate relation ($\mathbf{v}_E = \mathbf{r}_E\omega_0$)

$$\delta\mathbf{B}^{(0)}=\frac{[\mathbf{k}\times\mathbf{v}_E]}{2ck}(\delta E_l^{(1)}+\delta E_l^{(-1)}). \qquad (7)$$

It follows from (7) that the magnetic field is directed perpendicular to the plane made up of the vectors $\mathbf{k}$ and $\mathbf{E}_0$, and that it is not generated when the perturbations propagate strictly along the direction of the external electric field.

In accordance with formula (7), the generation of the magnetic field constitutes induced excitation by the field of the high-frequency plasma waves. The onset of the Langmuir field may be due to the development in the plasma of either ion-sound or aperiodic parametric instabilities. In the former case the growing magnetic field has a frequency on the order of the ion-sound frequency, and in the latter case the field increases aperiodically. The threshold field of the pump wave is determined by the known expressions[12] for the case of development of the potential periodic or aperiodic instabilities, respectively.

3. We have dealt above with perturbations propagating in a wide range of angles $\theta > \theta_0$. We proceed now to analyze the instability in the angle range $\theta < \theta_0$. Since this angle interval is small, it suffices to consider the case of strictly transverse propagation of the perturbations, when the wave vector k is perpendicular to the external electric field. We then have for the components of the electric field oriented along the external field, in accordance with relations (3),[11]

$$\{(\omega+n\omega_0)^2\varepsilon^{tr}(\omega+n\omega_0,k)-c^2k^2\}\varphi^{(n)}+{}^1/_4k^2v_E^2\Big\{(\varphi^{(n)}+\varphi^{(n+2)})\times\delta\varepsilon_e(\omega+(n+1)\omega_0,k)\left[1-\frac{\delta\varepsilon_e(\omega+(n+1)\omega_0,k)}{\varepsilon(\omega+(n+1)\omega_0,k)}\right]+(\varphi^{(n)}+\varphi^{(n-2)})\times\delta\varepsilon_e(\omega+(n-1)\omega_0,k)\left[1-\frac{\delta\varepsilon_e(\omega+(n-1)\omega_0,k)}{\varepsilon(\omega+(n-1)\omega_0,k)}\right]\Big\}=0. \qquad (8)$$

Here $\varphi^{(n)}=\mathbf{E}_0\cdot\delta\mathbf{E}^{(n)}/(\omega+n\omega_0)$ and $\varepsilon^{tr}=1+\delta\varepsilon_e^{tr}+\delta\varepsilon_i^{tr}$ is the transverse dielectric constant of the plasma.

Taking into account the smallness of the parameter $v_E/c \ll 1$, it suffices to restrict the analysis of the system (8) in the frequency region $\omega_0 \approx \omega_p$ to the harmonics $\varphi^{(0)}$ and $\varphi^{(\pm 2)}$. In this case we obtain for perturbations with wave number $k \gg \omega_0/c$

$$\varphi^{(0)}-\varphi^{(2)}-\varphi^{(-2)}=0,$$

$$\varphi^{(\pm2)}+\frac{1}{4}\left(\frac{v_E}{c}\right)^2(\varphi^{(0)}+\varphi^{(\pm2)})\frac{1}{\varepsilon(\omega\pm\omega_0,k)}=0. \qquad (9)$$

From the system (9) we get the dispersion relation

$$1+\left(\frac{v_E}{2c}\right)^2\left[\frac{1}{\varepsilon(\omega+\omega_0,k)+(v_E/2c)^2}+\frac{1}{\varepsilon(\omega-\omega_0,k)+(v_E/2c)^2}\right]=0. \qquad (10)$$

We analyze Eq. (10) under the assumption that the wavelength of the perturbations is small compared with the electron mean free path. We can then write the following expressions for the dielectric constants:

$$\varepsilon(\omega\pm\omega_0,k)=2\omega_0^{-1}[\Delta\omega_0(k)\pm(\omega+i\tilde{\gamma}(k))], \qquad (11)$$

where $\Delta\omega_0 = \omega_0 - \omega_p(1+\frac{3}{2}k^2r_{De}^2)$, r_{De} is the Debye radius of the electrons, and $\tilde{\gamma}$ is the high-frequency damping decrement of the plasma waves and is determined by the collisions of the electrons with the ions and by the Landau damping.

Taking (11) into account, we find that the solution of the dispersion equation (10) is $\omega = i\gamma$, where

$$\gamma=-\tilde{\gamma}+\left[-\left(\Delta\omega_0+\frac{v_E^2}{8c^2}\omega_0\right)^2-\frac{v_E^2}{4c^2}\omega_0\left(\Delta\omega_0+\frac{v_E^2}{8c^2}\omega_0\right)\right]^{1/2}. \qquad (12)$$

Thus, an aperiodic growth of the magnetic-field perturbations takes place in our case. For the instability to develop it is necessary to satisfy the condition $\Delta\omega_0 < 0$. We emphasize that the considered parametric instability is due entirely to allowance for the nonpotential character of the growing Langmuir oscillations. In the approximation of a nondissipative cold plasma ($\tilde{\gamma}=0, r_{De}=0$) an instability of this type was observed in Ref. 11.

Of greatest practical interest is the analysis of formula (12) in a wavelength range corresponding to the condition $k < k_{st}$, where k_{st} is the wave number at which the contribution of the Cerenkov damping and the decrement

$\tilde{\gamma}$ becomes comparable with the contribution of the electron–ion collisions. The reason is that at $k>k_{st}$ an exponential growth takes place with increasing Landau-damping wave number and leads, according (12) to a decrease of the increment and to a suppression of the instability. We note that under typical conditions of a laser plasma, the quantity $k_{st}r_{De}$ depends little on the temperature and turns out to be $k_{st}r_{De}\approx 0.3$.

At $k<k_{st}$ we have $\gamma=\nu_{ei}/2=\text{const}$, where ν_{ei} is the frequency of the electron–ion collisions. Taking this into account, we find that the maximum of the growth rate (12) is reached at $k=k_{\max}$:

$$k_{\max}=r_{De}^{-1}\left[\frac{2}{3}\frac{\omega_0-\omega_p}{\omega_p}+\frac{v_E^2}{6c^2}\right]^{1/2}, \tag{13}$$

and is determined according to (13) by the relation

$$\gamma_{\max}=\frac{v_E^2}{8c^2}\omega_0-\tilde{\gamma}. \tag{14}$$

The threshold of the nonpotential aperiodic instability exceeds the corresponding value for the potential aperiodic instability by a factor $c^2/r_{De}^2\omega_0^2$ and is determined by the formula

$$\frac{v_{E\,\text{thr}}^2}{c^2}=\frac{E_{0\,\text{thr}}^2}{4\pi n_e m_e c^2}=\frac{8\tilde{\gamma}}{\omega_0}=\frac{4\nu_{ei}}{\omega_0}. \tag{15}$$

At $E_0>E_{0\,\text{thr}}$ the instability develops in the wave-number range $k_-<k<k_+$, whose boundaries

$$k_\pm=\left[k_{\max}^2\pm\frac{2}{3}r_{De}^{-2}\frac{\tilde{\gamma}}{\omega_0}\left(\frac{E_0^2}{E_{0\,\text{thr}}^2}-1\right)^{1/2}\right]^{1/2}$$

move apart with increasing pump-field intensity.

As applied to the laser-plasma conditions, it is convenient to rewrite (15) in the form

$$q_{\text{thr}}=5\cdot 10^{15}\frac{Z}{T_e^{3/2}\lambda_0^3}\left[\frac{\text{W}}{\text{cm}^2}\right], \tag{16}$$

where q_{thr} is the threshold value of the energy flux density of the laser radiation. Here T_e is the electron temperature of the plasma in kiloelectron volts, λ_0 is the pump vacuum wavelength in microns, and Z is the ion charge multiplicity. The threshold (16) decreases sharply with increasing plasma temperature. At the same time, under typical experimental conditions realized at the present time the instability threshold (16) turns out to be high. For example, at $T_e\approx 1$ keV, $Z\approx 1$, and $\lambda_0\approx 1$ μm we obtain $q_{\text{thr}}\approx 5\times 10^{15}$ W/cm^2. At the same time, the magnetic-field generation connected with the development of the potential parametric instability[12] begins at much lower values of the pump-wave field intensity. For the conditions indicated above we have $q_{\text{thr}}\approx 10^{13}$ W/cm^2.

4. The analysis above has shown that the development of parametric instabilities in a plasma is accompanied by an exponential growth of the magnetic-field perturbations with time.

Since the buildup of the perturbations, which propagate almost perpendicular to the external field, is hindered by the high instability threshold for this direction, we must assume that the principal role in the generation of the magnetic fields is played by the parametric instability investigated in Ref. 12. Depending on the intensity of the pump electric field, the growth rate of this parametric instability varies in the range

$$\nu_{ei}\lesssim\gamma\lesssim\omega_{Li}^{2/3}\,\omega_{Le}^{1/3}. \tag{17}$$

It follows from (17) that the characteristic time of generation of the magnetic field does not exceed $(T_e^{3/2}\lambda_0^2/Z)$ psec (T_e is in kiloelectron volts and λ_0 is in microns), and under typical conditions of laser irradiation of targets it turns out to be much shorter than the duration of the laser pulse. Thus, in a laser plasma the magnetic field is determined by the nonlinear saturation of the parametric instability.

Under real conditions, the plasma is inhomogeneous. Therefore the magnetic field is generated in a region close to the critical densities. The spatial dimension of this region is determined by the condition that there be no strong Landau damping and turns out to equal $\Delta x\approx 3(k_{st}r_{De})^2L$, where L is the characteristic scale of the density inhomogeneity. For a laser plasma $\Delta x\approx 0.3L$ and constitutes an appreciable fraction of the corona.

In accordance with (7), the generated magnetic field is determined completely by the field of the Langmuir waves. The development of parametric instability can be accompanied by the onset of both random and regular Langmuir fields. There is therefore a possibility of exciting in the plasma not only a regular magnetic field, but also random microfields. Formula (7) can be used to estimate the quasistationary magnetic-field level that is established when parametric instability is saturated from the value of the quasistationary Langmuir field.

We now estimate the magnetic field established in the region of the buildup of oscillations with random phases when, for example, the ion-sound parametric instability is saturated. We take the nonlinear process that limits the exponential growth of the amplitude of the parametrically excited Langmuir oscillations to be their secondary parametric instability to the excitation of other Langmuir and ion-sound waves.[13] In this case the field intensity of the plasma waves in the buildup region is equal to the intensity of the electric pump field. We then obtain for the magnetic field from formula (7)

$$\delta B\sim c^{-1}v_E E_0. \tag{18}$$

Under laser-plasma conditions, the estimate (18) points to the possibility of generation of strong magnetic fields. Thus, for example, for a plasma obtained with the aid of a neodymium laser, at an energy flux density ~10^{16} W/cm^2, we have according to (18) $\delta B\sim 1$ MG. As seen from (18), the magnetic field intensity increases in proportion to the pump-wave energy flux density.

Besides the instabilities investigated above, secondary instabilities, connected with excitation of a low-frequency perturbation and of another Langmuir wave by the Langmuir wave, are also possible in the plasma. A magnetic field can also be generated on account of the development of these secondary instabilities. The role of the pump is assumed in this case by Langmuir wave itself. The connection between the generated magnetic field and the field of the Langmuir waves is determined by a formula similar to (7), and the intensity of the magnetic field turns out to have a quadratic dependence on the electric field intensity of the Langmuir oscillations.

Thus, in those regions where an intense Langmuir field is present, magnetic fields can be generated. In the corona of a laser plasma, this can occur not only at densities close to crictical, but also in the vicinity of a density close to one-quarter the critical value.

We note in conclusion that the onset of strong magnetic fields in a laser plasma can greatly influence the transport coefficients. If the magnetic field has a regular character and its spatial structure is characterized by a wavelength λ (which coincides according to (6) with the wavelength of the Langmuir oscillations), this influence will be appreciable if the Larmor radius of the electron is small compared with λ. This condition can be rewritten in the form

$$\delta B \geqslant (T_e^{1/2}/\lambda)\,[\mathrm{MG}], \tag{19}$$

where T_e is in kiloelectron volts and λ is in microns. For a plasma with an electron temperature ~1 keV and a perturbation wavelength amounting to some tenths of λ_0, the condition (19) is realized for a magnetic field of several megagauss. We note that the transport phenomena can be influenced not only by the regular but also by stochastic turbulent magnetic fields.[14]

The authors thank L. M. Gorbunov and V. P. Silin for interest in the work and for a useful discussion.

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Translated by J. G. Adashko

Theory of electron dragging by a spin wave in a layered medium

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(Submitted 21 September 1978)
Zh. Eksp. Teor. Fiz. 76, 1593–1601 (May 1979)

A surface spin wave produces in a semiconductor layer static extraneous currents of conduction electrons. If the layer is galvanically open-circuited, then two types of extrinsic currents flow: the dragging current and the thermoelectric current. These currents are calculated by time-averaging the phenomenological expression for the instantaneous current density. It is shown that the principal role is played by the dragging current. The dragging leads to formation of current eddies and to the onset of a magnetic moment of the sample. The problem of the distribution of the electric field over the layer, with the eddies taken into account, is solved. The longitudinal and transverse potential differences are calculated and the induced magnetic moment is estimated.

PACS numbers: 72.20.My, 75.30.Cr, 75.30.Ds

The dragging of electrons by a spin wave (SW) was first observed apparently in thin metallic nickel films.[1] This effect, which turned out to be quite sizable, was subsequently observed in the magnetic simiconductor $HgCr_2Se_4$ (Ref. 2) and ultimately in a structure consisting of flat layers of yttrium iron garnet (YIG) and n-InSb in contact with one another.[3] The dragging effect was later investigated experimentally in detail both in magnetic semiconductors[4-6] and in layered structures[7,8] (see also the review[9]). Despite the large number of experiments, the theory of the effect is seemingly only in its initial stage. All that were discussed mainly were the possible nonlinearity mechanisms that contribute to the time-averaged electron current in the SW field.[5,10,11]

The present paper is devoted to the theory of the dragging effect in layered semiconductor + ferrite structures. Starting from a phenomenological formula derived and discussed in Ref. 12 for the electron-current density, we obtain by time-averaging two types of extrinsic currents in a galvanically open-circuited sample. The first type is the dragging current $I^{(d)}$ that satisfies the Weinreich relation.[13]

$$I^{(d)} \sim \mu_d \alpha W / v_{\mathrm{ph}}, \tag{1}$$

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