$$T_{-}-T_{+}={}^{8}/_{5}T(t)\,\mu/\mu_{\perp}.$$

We see therefore that after a time

$$t=\frac{1}{5qm^{1/2}}\left(\frac{5}{8}T_{T}\frac{\mu_{\perp}}{\mu}\right)^{3/2}$$

(T_T is the thermal energy) the energy difference reaches the value T_T.

In the situation described above, the plasma becomes spin-polarized in phase space. Polarization in ordinary space can be the consequence of this process. In fact, as stated above, the particle moves mainly along the force line, i.e., almost along a straight line. Assume that the region in which the MHD waves propagate is finite, so that the particles ultimately leave this region and land in an unperturbed homogeneous magnetic field. Because of the different velocities of particles of different type, the leading front of the plasma, on entering the homogeneous field, is polarized in analogy with the situation in the Stern-Gerlach experiment (see the end of Sec. 2). A similar situation will obtain also when accelerated particles land in the region where the plasma can no longer be regarded as collisionless (for example, the spilling of particles from the earths magnetosphere into the ionsphere). Particles of different types, having different energies, have different mean free paths, as a result of which the plasma becomes polarized over a finite length.

The described acceleration and polarization of particles in phase space can take place in an interstellar plasma, and produce in turn circular polarization of the passing electromagnetic radiation.

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Advanced magnetothermal phenomena in a laser plasma

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The possibility of spontaneous generation of sizable magnetic fields in a laser plasma within periods much shorter than hydrodynamic times is demonstrated. The field growth has a threshold that depends on the geometrical factors of the corona and on the dimensions of the heat-release region. The magnetic fields cause a qualitative restructuring of the heat front and decrease the heat transfer substantially. In particular, heat can penetrate into the plasma in the form of a magnetothermal jet.

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1. The magnetic fields produced spontaneously in a laser plasma were investigated in a number of experimental and theoretical studies.[1-21] The hertofore considered magnetic-field generation mechanisms can be broken up into two groups.

The first group includes mechanisms connected in one way or another with the process of absorption of the laser radiation. Notice should be taken of the transfer of momentum from the light wave to the electrons in resonant absorption when the radiation is obliquely incident on the critical surface,[2,3] of the electromagnetic instability due to the anisotropy of the distribution of the electron velocities in the absorption zone,[4,5] of the parametric generation of magnetic fields by beats between the high-frequency motions of the electrons,[6,7] and of the generation of magnetic fields by the advanced Langmuir motions in the plasma.[8] The foregoing field-generation methods are of considerable interest from the point of view of their influence on the light absorption, and particularly on the diffraction of the absorbed energy and its distribution over the critical surface. The

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magnetic fields that can appear as a result of these mechanisms are small-scale and are localized in a relatively small region near the critical surface (with dimension of the order of the radiation wavelength). The indicated causes of the fields are connected in one way or another with the mechanisms of anomalous absorption of laser radiation near the critical surface. As a rule, the anomalous absorption processes is important at radiation intensities ($\lambda=1.06$ μm)$I \sim 10^{15}-10^{16}$ W/cm^2. At the intensities $\sim 10^{13}-10^{14}$ W/cm^2 considered by us, the contribution of the anomalous mechanism to the absorption is small. At any rate, the magnetic fields localized in the near-critical region do not influence significantly the heat transfer from the absorption zone to the ablation zone, in view of the small thermal resistance of this thin plasma layer.

The other group of mechanisms that generate magnetic fields is connected with thermoelectric phenomena in the plasma.[9,10] It has been pointed out[10,11] that these fields, which have characteristic scales of the order of the reciprocal gradients of the temperature and of the plasma density, can decrease substantially the thermal conductivity of the plasma, and also disturb the symmetry of its motion.

Further investigations in this direction are of great interest. On the one hand, reliable data on the intensities of the magnetic fields and on their influence on the plasma are still unavailable. This is due to the difficulty of observing even megagauss fields in a dense plasma with small dimensions and abrupt density gradients (10–100 μm) and with a short lifetime 0.1 – 1 nsec. Measurements[12,13] of magnetic fields with the aid of the Faraday rotation of the polarization plane of even the fourth harmonic of the fundamental radiation do not make it possible to determine reliably the field in a region in which the density greatly exceeds the critical value, owing to refraction in the inhomogeneous plasma. On the other hand, the reported numerical experiments (see, e.g.. Ref. 14) on laser irradiation of targets were performed under conditions (high radiation intensities, short pulse times) such that many other processes, including hydrodynamic, are substantial, and not only the generation of the magnetic fields. This does not provide information on the conditions under which the magnetic fields in the laser plasma can reach large intensities, on the observable effects that they can produce, on the physical causes that limit their growth, and on the extent to which their influence on the heat transfer is significant.

Theoretical investigations of the excitation of a magnetic field by an inhomogeneous thermoelectric power developed during that time in the following manner. It was shown in Ref. 10 that the thermoelectric mechanism that produces the fields can have the character of an instability. The factor that can exert a stabilizing influence on the development of instability and limit the growth of the fields in the generation region, as noted in Ref. 10 can be the outflow of the fields at a rate proportional to the temperature gradient.

On the basis of this remark, some investigators have reached the erroneous conclusion[15-17] that the magnetic fields cannot grow to appreciable values. Since this error gained wide acceptance in the literature (e.g., Refs. 2 and 8), we present in Sec. 2 simple qualitative arguments that reveal the physical cause of this error.

In the present paper, on the basis of a qualitative analysis and numerical calculations, we wish to investigate different variants of the evolution of magnetic fields in a laser plasma, and demonstrate in particular the critical character of the fate of the magnetic fields relative to the spatial scales of the corona and the dimensions of the heat-release region. One of the possible variants of heat transfer turns out to be formation of the magnetothermal jet.

2. It was established in Ref. 10 that when the inequality $l/a \gg (m/M)^{1/2}$ is satisfied (l is the mean free path, a is the inhomogeneity scale of the temperature and density of the plasma, and m and M are the masses of the electrons and ions respectively (the magnetic fields produced in an inhomogeneous plasma begin to influence the thermal conductivity (at $\omega_e\tau \sim 1, \omega_e=eB/mc, \tau$ is the time between the collisions of the electrons with the ions) long before the plasma expansion takes place. We can therefore separate the magnetothermal part of the complete problem from the hydrodynamic part, assuming the plasma to be immobile and its density to be a specified function of the coordinates.

The physical conditions corresponding to this inequality are that the electron heat fulx velocity be much larger than the plasma expansion velocity, i.e., the picture of the plasma expansion is essentially nonstationary. Under the same conditions we can neglect the energy exchange between the electrons and ions. In Ref. 10 was derived a closed system of simplified equations for the magnetic field and for the electron temperature, and this system described correctly the magnetothermal phenomena over periods shorter than the hydrodynamic times:

$$\frac{\partial \mathbf{B}}{\partial t}=\frac{c}{e}[\nabla T\times\nabla \ln n]-\frac{c}{e}\operatorname{rot}\frac{\mathbf{R}_T}{n}, \tag{1}$$

$$\frac{3}{2}n\frac{\partial T}{\partial t}=-\operatorname{div}\mathbf{q}+S. \tag{2}$$

Here n is the electron density, $\mathbf{R}_T=-\beta_\parallel^{uT}\nabla_\parallel T-\beta_\perp^{uT}\nabla_\perp T-\beta_\wedge^{uT}[\mathbf{h}\times\nabla T]$ is the thermopower, S is the density of the sources, $\mathbf{q}=\varkappa_\parallel^e\nabla_\parallel T-\varkappa_\perp^e\nabla_\perp T-\varkappa_\wedge^e[\mathbf{h}\times\nabla T]$ is the electron heat flux, $\mathbf{h}=\mathbf{B}/B$, and the expressions for β and $\varkappa$ are given in Ref. 22.

Equation (1) describes the growth of the magnetic field on account of the source (the inhomogeneous thermal emf $ce^{-1}[\nabla T\times\nabla \ln n]$) and its drift on account of the thermopower into the cold region together with the heat flow, as a result of the fact that the magnetic field is better frozen-in in hot electrons than in cold ones. At $\omega_e\tau \ll 1$ the drift velocity is $v_q \sim q/nT \sim \tau m^{-1}\nabla T$,[10] so that it might seem that during the time of the electronic heat conduction $\sim a^2/lv_{Te}$ the magnetic field reaches a stationary value with $\omega_e\tau \lesssim 1$, determined by the balance between the two terms of the right-hand side of (1). However, at $\omega_e\tau \gtrsim (\omega_e\tau)^*$, with $(\omega_e\tau)^* \sim 1$ only for hydrogen ($[\omega_e\tau]^*$ depends on the charge of the ions and, for example, in the case of an aluminum target $[\omega_e\tau]^*=0,1$), the magnetization of the electronic thermoconductivity is simultaneously accompanied by magnetization of the velocity of this thermal drift. At $\omega_e\tau \gg (\omega_e\tau)^*$ the drift velocity is small:

$$v_q \sim \frac{\tau}{m}\nabla T\frac{(\omega_e\tau)^{*2}}{(\omega_e\tau)^2}$$

and the generation cannot be offset by the drift. Therefore, if a field with $\omega_e\tau > (\omega_e\tau)^*$ is produced in the plasma, the field will increase up to hydrodynamic times.

The threshold character of the growth of the magnetic fields can be illustrated by using a model equation for $\omega_e\tau$, which is the qualitative analog of Eq. (1):

$$\frac{d}{dt}(\omega\tau)=A-C\omega_e\tau/(1+(\omega_e\tau)^2). \tag{3}$$

Here $A \sim C \sim lv_{Te}/a^2$. In the derivation of (3) it was assumed that the spatial structure of the distribution of the temperature and of the magnetic field does not change with time, so that to estimate the gradients all the spatial derivatives were placed by the ratios of the corresponding quantities to their spatial scales. If we make the time dimensionless relative to the electron thermal-conductivity time a^2/lv_{Te} which is characteristic of our problem, then all the terms in (3) become of the same order.

If $A < C/2$, the quantity $\omega_e\tau$, increasing from 0, saturates at a value

$$(\omega_e\tau)_0=\frac{C}{2A}-\left(\frac{C^2}{4A^2}-1\right)^{1/2},$$

with $(\omega_e\tau)_0 \sim 1$. On the other hand if $A > C/2$, there is no saturation and $\omega_e\tau$ increases without limit (linearly with time at periods larger than the heat-conduction time). This simplified analysis shows the extent to which even the qualitative picture of the growth of the magnetic field is sensitive to the numerical ratio of the parameters of the problem.

Using the model equation (3) we can estimate the saturation level of generated fields that have a smaller scale on account of the thermal drift. Let the characteristic scale of the inhomogeneity of the plasma density a be much larger than the characteristic scale $1/k$ of the temperature inhomogeneity and accordingly of the magnetic field. The generation term in (3) is then $A \sim lv_{Te}k/a$, and the drift term $C \sim lv_{Te}k^2$. Since $ka \gg 1$, it follows that $A \ll C$ and, in accordance with the foregoing analysis, saturation of the magnetic fields takes place at $(\omega_e\tau)_0 \sim A/C \sim 1/ka \ll 1$. Shorter-wave perturbations of the magnetic field saturate on a lower level of $\omega_e\tau$. Only at $ka \sim 1$ can values $\omega_e\tau \gtrsim 1$ be reached.

It is of interest to estimate the influence of the small-scale magnetic fields on the transport processes in an inhomogeneous plasma. The degree of the influence of the magnetic fields with scale $1/k$ on the transport is determined by the ratio of the Larmor radius of the electron r_k to the scale $1/k$. Since $\omega_e\tau \sim 1/ka$, it follows that $kr_k \sim k^2la \gg 1$, i.e., the small-scale magnetic field does not exert a substantial influence on the transport in the region of interest to us, and therefore classical values were used for the transport coefficients in (1) and (2). The question of reaching a stationary value of the magnetic field within heat-conduction times can be solved by numerical simulation, the simplest nontrivial case being the two-dimensional one.

3. Equations (1) and (2) were solved numerically in a rectangular region with dimensions 500 μm along the X axis and 50–200 μm along the Z axis. The distributions of the magnetic field, of the temperature, and of the plasma density along the Y axis were assumed to be uniform. A detailed description of the numerical experiments is given in Ref. 18. In the geometry in question, the magnetic field has one component along the Y axis. This geometry corresponds to a physical situation in which radiation propagating along the Z axis is focused by a cylindrical lens on a flat target only along one direction X.

The density was assumed to depend only on the coordinate Z and increased from the critical value 10^{21} cm^{-3} (for a neodymium laser) to 5×10^{22} cm^{-3}, corresponding to the electron density of metallic aluminum. The heat sources were distributed near the critical surface. The distribution of S over the critical surface was chosen to be Gaussian, $S \sim \exp(-x^2/b^2)$, b =17–200 μm. The maximum of the heat release corresponded to energy fluxes 3×10^{14} W/cm^2. The initial conditions were chosen to be those of a non-magnetized plasma with temperature 100 eV.

4. The calculations have shown that saturation of the generation of the magnetic field has a threshold governed by the geometrical factors of the experiment, particularly by the abruptness of the spatial distribution of the heat input. At a sufficiently homogeneous heat-source distribution (b =200 μm) the magnetic field increased to a value close to $(\omega_e\tau)^*$, and was saturated because of its removal by the heat flux. In the case of a more localized heat release (b =50 μm), i.e., the magnetic field increased practically linearly with time, at least up to hydrodynamic times, and reached after a time ~0.5 nsec a value ~5 MG with a maximum value $\omega_e\tau \gg (\omega_e\tau)^*$. Figure 1 shows plots of the maximal values of the magnetic field and of $\omega_e\tau$ against the time for these two cases.

In the case of a "breakthrough" of the magnetic field through the threshold ($\omega_e\tau \gg (\omega_e\tau)^*$) the maximum magnetic field depends weakly on the geometry of the heat

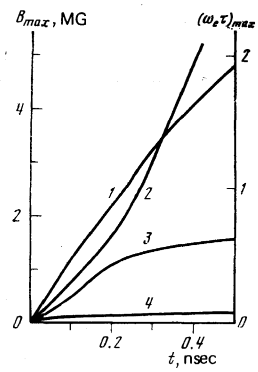

FIG. 1. Maximal values of the magnetic field B (curves 1, 3) and of the parameter $\omega_e\tau$ (curves 2, 4) for different dimensions of the heat-release region b: curves 1, 2—b = 50 μm, 3, 4—b = 200 μm.

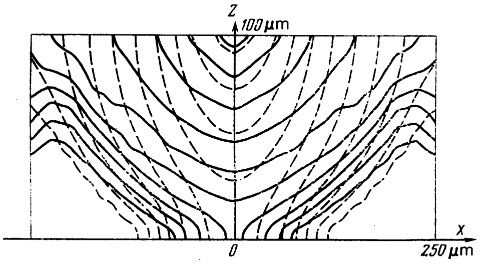

FIG. 2. Isotherms at $t=0.3$ nsec ($q=3\cdot10^{14}$ W/cm^2, $b=50$ µm) $(\nabla \ln n)^{-1}$ 25 µm) in a plasma with (solid lines) and without (dashed) allowance for the magnetic-field generation.

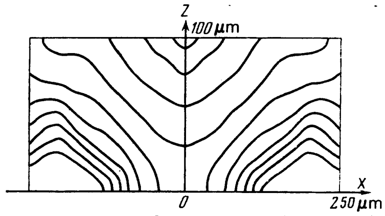

FIG. 4. Isotherms obtained when the heat-release source is turned on instantaneously ($t=0.5$ nsec, $q=3\times10^{14}$ W/cm^2, $b=50$ µm, $(\nabla \ln n)^{-1}=25$ µm).

release, on the initial temperature level, and on the magnetic field (up to several hundred kilogauss) and is determined principally by two parameters: the maximum density of the heat flux q and the spatial scale of the density distribution. These fields exert a substantial influence on the heat transport. We note that even at $\omega_e\tau\sim3$ the thermal conductivity of an aluminum plasma decreases by a factor 10^3 compared with its value in the absence of a field. Figure 2 shows the distributions of the plasma temperatures obtained with and without allowance for the generation of the magnetic field. It is seen that the magnetic field slows down the heat transfer into the dense region, producing so to speak magnetic "mirrors" with $\omega_e\tau\gtrsim(\omega_e\tau)^*$. Therefore the heat input spreads over layers close to the critical surface and overheats them. When the magnetic mirrors are heated they expand towards the denser plasma, since the most intense generation of the magnetic fields occurs near the front of the thermal wave.

In the case of heat release strongly localized on the X axis, the magnetic mirrors can hinder the propagation of the heat also along the X axis. Figure 3 shows the distribution of the temperatures in this case. A characteristic feature of localized heat release is the formation of a thermal jet. The growth of the magnetic fields leads to a unique focusing of the heat flux. The heat region takes the form of a needle that broaches the magnetic mirrors.

The character of the growth of the heat-releasing source with time turns out to have little effect on the maximum magnetic field, but is important for the determination of the spatial structure of the thermal wave. Instantaneous turning on of the heat source was compared with its slow (linear) growth to the maximum value in a time 0.1 nsec. Physically the abruptness of the turning on of the thermal source is determined by the ratio of the width of the leading front of the heating radiation to the heat-conduction time over the characteristic scale of the corona. Figures 4 and 5 show the characteristic isotherms for these two cases. In some variants, when the heat flux was turned on smoothly, a relative minimum of the temperature was observed at the center of the heat-release zone. Such a temperature distribution can be explained by considering the propagation of the heat inside a region with inhomogeneous thermal conductivity. In zones with low thermal conductivity (magnetic mirrors) the temperature gradients must be large to pass a given heat flux from a boundary. The appearance of two maxima at the boundary of the heat-release zone is closely connected with the character of magnetization of the thermal drift of the magnetic fields. Thus, in control variants with the thermal drift turned off, analogous temperature profiles appeared even when the heat release had a smooth distribution ($b=200$ µm).

In a number of variants, the diffusion of the magnetic field was taken into account. Its influence on the generation of the magnetic field and on the temperature distribution turned out to be insignificant everywhere with the possible exception of the densest plasma layers.

In the situation when the magnetic field is not stabilized by thermal drift, the intensities of the maximal fields, the distribution, and their influence on the plasma dynamics should be determined in a slow hydrodynamic time scale by simultaneously solving Eqs. (1) and (2) with the equations of motion of the plasma. The same time scales apply to generation of magnetic fields when oblique temperature and density gradients are produced by the inhomogeneous plasma motion itself. A most interesting example of such a situation is the generation of magnetic fields in the course of Rayleigh-Taylor instability, considered in Refs. 19–21. The question of the intensity of the magnetic fields in this case has not been finally answered to this day, inas-

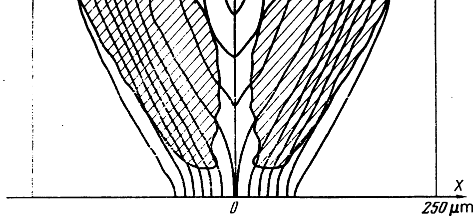

FIG. 3. Isotherms at $t=0.45$ nsec ($q=10^{14}$ W/cm^2, $b=17$ µm, $(\nabla \ln n)^{-1}=12.5$ µm). The region with $\omega_e\tau>(\omega_e\tau)^*$ is shaded.

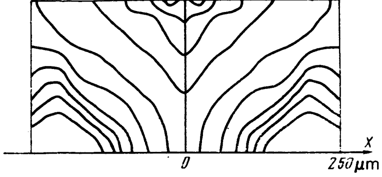

FIG. 5. Isotherms for slow (linear) growth of the heat-release source ($t=0.5$ nsec, $q=3\times10^{14}$ W/cm^2, $b=50$ µm, and $(\nabla \ln n)^{-1}=25$ µm).

much as in Refs. 19 and 20 they investigated only the linear stage of the instability, and the thermal drift was not taken into account in the numerical experiment of Ref. 21, so that doubts are case on the large magnetic fields obtained in the latter reference.

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Investigation of a magnetic phase transition in a layer of carbon monoxide molecules chemisorbed on platinum and gold surfaces in a magnetic field

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The thermomagnetic effect was used to study a magnetic phase transition in a monolayer of chemisorbed carbon monoxide molecules. A study was made of the dependence of the critical field H_c, in which the magnetic phase transition took place, on the surface concentration of the carbon monoxide molecules x_{CO} and the surface temperature T. A phase transition was also observed in a monolayer of chemisorbed hydrogen on the surface of platinum. An analysis was made of a change in the nonspherical interaction between the molecules and the surface which occurred at the phase transition. The observed dependences of H_c on x_{CO} and T could be interpreted as a consequence of a transition in a monolayer of chemisorbed particles from an antiferromagnetic to a ferromagnetic state in a field $H = H_c$.

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1. INTRODUCTION

Numerous investigations of the adsorption on the surfaces of solids have established that, in addition to the interaction between the adsorbed atoms or molecules and the surface atoms of the crystal lattice of the adsorbent, there are also forces acting between the adsorbed atoms or molecules. The existence of a strong interaction between chemisorbed atoms or molecules is supported by, for example, the concentration (coverage) dependences of the heat of adsorption and electron work function.[1] Low-energy electron diffraction investigations have shown that this type of interaction can produce a monolayer with a geometrically ordered structure.[1] The collective interaction of adsorbed atoms and molecules can also result in magnetic ordering in the adsorption monolayer. As demonstrated in our earlier investigation,[2] a system of this kind may exhibit magnetic phase transitions induced by a magnetic field. A phase transition in a monolayer of chemisorbed carbon monoxide molecules on gold and platinum surfaces was ob-

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