Lett. 35, 663 (1975).
[7]C. E. Max, W. M. Manheimer, and J. J. Thomson, Phys. Fluids 21, 128 (1978).
[8]W. Woo and J. S. De Groot, Phys. Fluids 21, 124 (1978).
[9]Yu. V. Afanas'ev, E. G. Gamaliĭ, I. G. Lebo, and V. B. Rozanov, Zh. Eksp. Teor. Fiz. 74, 516 (1978) [Sov. Phys. JETP 47, 271 (1978)].
[10]L. A. Bol'shov, Yu. A. Dreĭzin, and A. M. Dykhne, Pis'ma Zh. Eksp. Teor. Fiz. 19, 288 [JETP Lett. 19, 168 (1974)].
[11]B. A. Al'terkop, E. V. Mishin, A. A. Rukhadze, Pis'ma Zh. Eksp. Teor. Fiz. 19, 291 (1974) [JETP Lett. 19, 170 (1974)].
[12]D. A. Tidman and R. A. Shanny, Phys. Fluids 17, 1207 (1974).
[13]B. A. Al'terkop and E. V. Mishin, Phys. Lett. A 46, 319 (1974).
[14]A. F. Nastoyashchiĭ, At. Energ. 38, 27 (1975).
[15]S. I. Braginskii in: Voprosy teorii plazmy (Problems of Plasma Theory), ed. M. A. Leontovich, vol. 1, Atomizdat, 1963, p. 183.

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Fast waves in a laser plasma

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Heating of spherical targets by intense laser radiation in the Kal'mar installation produced, besides spherically symmetric emission of thermal plasma ions, as various methods have shown, also jet-like emission of groups of fast ions of energy ~0.5 MeV. It is shown by analysis of the possible acceleration mechanisms that the observed effect is due to ponderomotive acceleration of the ions in the region of the critical density by the resonantly amplified electric field of the laser radiation. A theory of resonant ion acceleration, in which account is taken of the nonlinear interaction of the plasma waves in the critical-density region, is proposed. Estimates of the energy and of the number of the fast ions agree with the experimental data.

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1. INTRODUCTION

An important problem in laser-mediated thermonuclear fusion is the determination of the redistribution of the absorbed energy in the plasma and the generation of the accelerated particles. The appearance of fast ions that carry away an appreciable fraction of the absorbed energy has been repeatedly reported in recent years.[1-3]

In experiments with the nine-channel "Kal'mar" laser installation in which solid and hollow shell targets were heated, a new effect was observed against the background of the spherically symmetrical expansion of the plasma corona, namely jetlike acceleration of particles to high energy (on the order of 0.5 MeV). The purpose of the present paper is to explain the physical nature of the observed effect and to calculate the main characteristics (number and energy) of the fast particles. It is shown on the basis of an analysis of the possible ion-acceleration mechanism in a nonstationary and inhomogeneous plasma that the formation of a jet of fast ions is due to electrostatic acceleration of the particles in the region of plasma resonance. The presently prevailing theory of ion acceleration in plasma resonance[4] can explain acceleration of the ions to an energy on the order of several dozen kiloelectron volt. To describe ion acceleration to higher energies it is necessary to take into account the nonlinear deformation of the structure of the electric field in the region of the plasma resonance, since the field pressure turns out in this case to exceed the gas kinetic pressure. Such a nonlinear theory of resonant acceleration of ions is developed in the present paper. The derived formulas explain the experimentally observed number and energy of the fast ions.

From the point of view of laser-mediated thermonuclear fusion, an important question is the influence of the ion jets on the symmetry of the compression of shell-like targets. In the case of resonant ion acceleration, this influence turns out to be small because the average momentum transferred to the shell during the time of the collapse is small compared with the gas-kinetic pressure of the plasma. This is due to the short duration of the ion-acceleration process.

2. INVESTIGATION OF THE DYNAMICS OF EXPANSION OF THE PLASMA CORONA

The experiments were performed on the nine-channel "Kal'mar" laser installation, the diagram and diagnostic assembly of which are described in detail in Refs. 5–7. The laser radiation was focused from nine directions on spherical targets of glass (SiO_2) of diameter from 80 to 200 μm, placed in the center of a vacuum chamber. With each light beam having in the region of the target a diameter ~150 μm and an optical

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energy $\lesssim 100$ J (the pulse duration at the base was $\tau \approx 2.5$ nsec), the flux density at the target was $q \sim 10^{14}$ W/cm^2.

To investigate the angular directivity and the rate of scattering of the plasma particle, various ultrahigh speed interferometric[8] and Schlieren[9] photography methods in the laser beam were used, as well as ion collectors mounted at various angles to the plasma.

The interferometric measurements performed at a pressure $P \sim 10^{-6}$–1 torr in the vacuum chamber have shown the plasma corona to have a high expansion symmetry (see Fig. 1). At high pressures of the residual gas (D_2) ($P \gtrsim 5$ Torr), however, in more than 50% of the experiments it was observed[10] that besides the spherical expansion of the plasma there appear several strongly peaked "jets," whose velocity exceeds noticeably the velocity of the spherical expansion of the plasma corona and reaches values $v \sim (2-3) \times 10^8$ cm/sec at the first instants after the end of the heating pulse (see Fig. 2). The number and direction of the jets registered by an interferometry method changed from experiment to experiment, but the angle between the directions of an individual jet and the nearest laser heating beam was in the range $\theta \sim 15$–$30°$.

Reduction of the interference pattern has made it possible to conclude that the electron density n_e inside the jet corresponds to complete ionization of the residual gas (deuterium) and increases with increasing pressure in the vacuum chamber. The lower limit of the sensitivity of the interferometric procedures was $n_e l \sim 5 \times 10^{16}$ cm^{-2}. (Here $l \sim 1$ mm is the characteristic jet diameter.) At these values, the jet could be registered only at $n_e \gtrsim 5 \times 10^{17}$ cm^{-3}, i.e., at a gas pressure $P \gtrsim 5$ Torr.

Another method of observing the dynamics of the region of ionization of the residual gas was multiframe Schlieren photography[9] for ~300 nsec after the end of the heating pulse. Figure 3 shows a characteristic 7-frame Schlieren photograph, the first frames of which demonstrate the increase of the transverse dimensions of the jet with time, explainable as due to formation of a cylindrical shock wave propagating from the gas ionization region. In the later stages of the observation, visualization of the jet becomes impossible. The reason is apparently that at a velocity $v' \lesssim 10^6$ cm/sec (starting with approximately 50 nsec) the cylindrical shock wave becomes non-ionizing and, as a result of the decrease of the gradients of the refractive index on its front, the deflection angle of the sounding radiation becomes smaller than the minimum value registered by the employed Schlieren method.

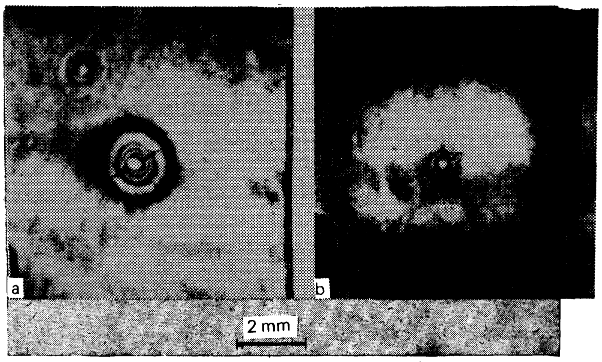

FIG. 1. Interference patterns of plasma corona spreading in vacuum ($P \sim 10^{-6}$ Torr): a) 5.4 nsec, b) 3.1 nsec.

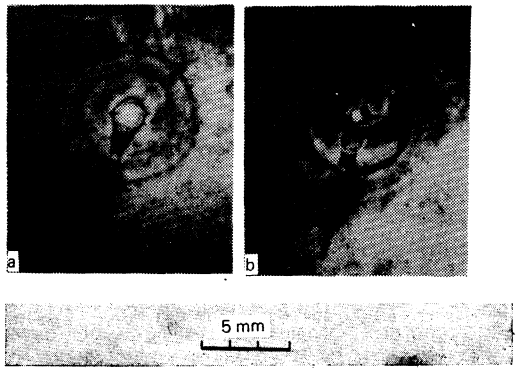

FIG. 2. Interference patterns of a plasma spreading in an atmosphere ($P \sim 10$ Torr) of deuterium: a) 3.1 nsec, b) 5.4 nsec.

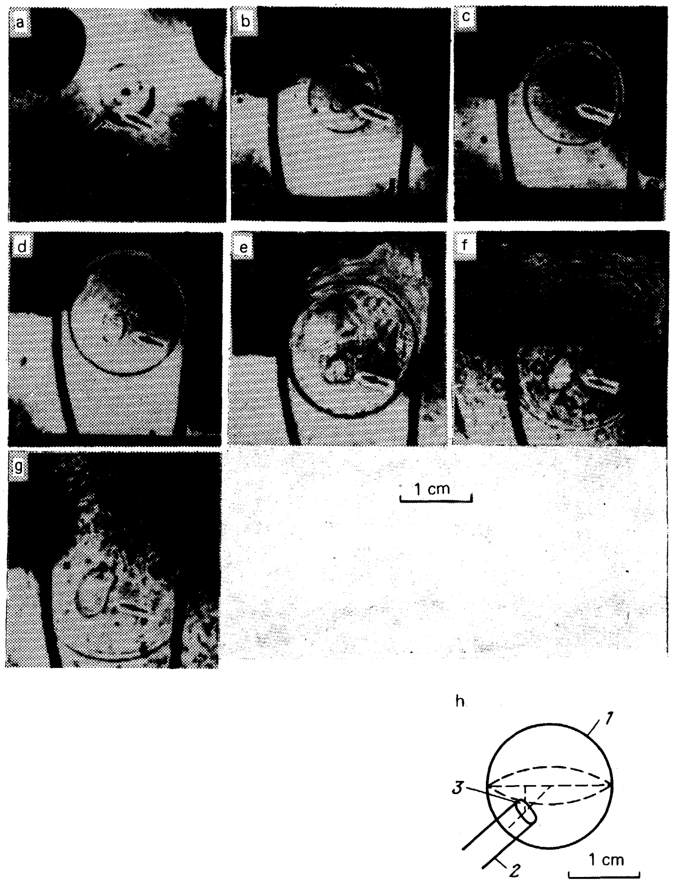

FIG. 3. Seven-frame Schlieren photograph of laser plasma at $P = 16$ Torr: a) 12 nsec, b) 20 nsec, c) 48 nsec, d) 92 nsec, e) 136 nsec, f) 226 nsec, g) 316 nsec, h) relative positions of spherical (1) and cylindrical (2) shock waves, 3—annular condensation region.

Nonetheless, the sensitivity of this method was sufficient to register an annular increas-density region that appeared as a result of interaction of the cylindrical shock wave with the spherical wave produced in the isotropic expansion of the plasma (Fig. 3, b and g). On the photographs, this region was elliptic in shape with an eccentricity and position that depended on the angle between the directions of observation and propagation of the jet. This has made it possible to determine both the spatial orientation of the jets (in contrast to the interoferometric measurements, which registered only the projection of the jet on the plane of the object), as well as the diagram of the expansion of the cylindrical shock wave. The motion of the latter (see Fig. 4) is satisfactorily approximated by the model of instantaneous cylindrical explosion[11]:

$$r=[\alpha(\gamma)W]^{1/4}\rho^{-1/4}t^{1/2},$$

where W is the explosion energy, ρ is the gas density, and $\alpha(\gamma)$ is the coefficient that depends on the adiabatic exponent γ. This made it possible to determine accurate to ~100% the value of W, which usually amounted to ~0.01 J/cm. The inaccuracy in the determination of W is due to some uncertainty in the choice of the value of γ, which depends on the temperature behind the front and can vary with time.

In the experiments performed in vacuum ($P \lesssim 10^{-5}$ Torr), the plasma expansion was investigated also with the aid of ion collectors with negative bias,[12] placed in the vacuum chamber at various angles.

It turned out that most experiments the oscillograms of the ion current registered only ions with thermal velocities ($v \lesssim 10^8$ cm/sec) that agreed with the data of the interferometric measurements. The number of ions per unit solid angle in the investigated direction differed by less than a factor 1.5, which barely exceeds the measurement error. Thus, a comparison of the results of the collector and interferometry measurements leads to the conclusion that the scattering of the thermal ions is isotropic.

In some experiments, however, a group of ion was observed with velocities $v > 10^8$ cm/sec, and in contrast to the thermal ions their scattering was anisotropic. Figure 5 shows the results of the reduction of the

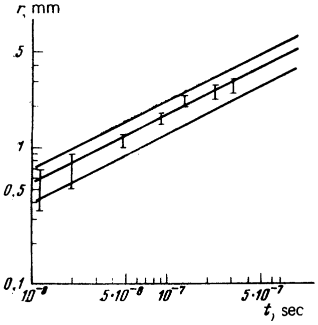

FIG. 4. The experimental r-t diagram of motion of a cylindrical shock wave is represented by the points. The curves corresponds to the model of instantaneous point explosion at energies W =5, 1.5, and 0.5 J/m, respectively for the upper, middle, and lower curves.

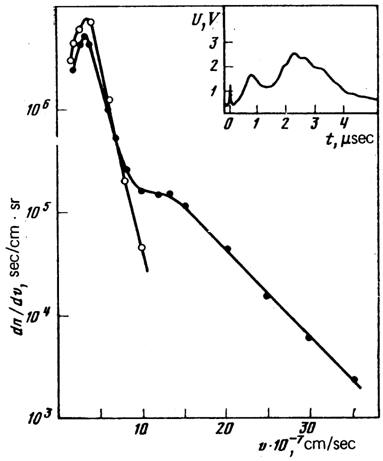

FIG. 5. Velocity distribution of the number of ions in an experiment where the fast ions are registered with one collector (●) and with the other (○). The inset shows an oscillogram of the signal of the first collector (distance from target L =110 cm).

oscillograms obtained with two collectors located at an angle ~90° to each other, when one of the detectors registered an appreciable number of fast ions ($v > 10^8$ cm/sec). The number of these ions is ~10^{10}, meaning ~5% of the total number of plasma ions entering the aperture of this collector (solid angle ~10^{-3} sr), and the energy (in the observation direction) is up to 40% of the total ion energy, with the average fast-ion energy ~280 keV. It must be emphasized that an isotropic spreading of fast ions was observed only in a small number of experiments and is apparently due to the small angle aperture of the employed system of ion collectors.

Even though the collector measurements in vacuum have confirmed the narrow directivity of the spreading and have established the velocity distribution of the fast ions, they cannot be used to estimate the total number of ions with velocities $v > 10^8$ cm/sec, since their directivity pattern remains unknown.

An estimate of the total number of fast ions can be made by using the results of high-speed optical methods. According to Ref. 13, the energy lost by an ion of velocity $v \approx 2.5\times10^8$ cm/sec when slowed down in the gas surrounding the target at a pressure $P \approx 5$ Torr is $dE/dx \sim 20$ keV/cm. The ratio of the energy W of the cylindrical shock wave to the energy loss per ion dE/dx is an approximate estimate of the number of fast ions in the jet, the characteristic value of which ~5×10^{12} ions, or 0.1% of the total number of ions in the target. The fraction of the energy of the fast ions of the jet amounts in this case to $\lesssim 3\%$ of the absorbed laser energy. We emphasize that this is a highly approximate estimate.

Simultaneously with the procedure described above, we photographed the plasma with a spatial resolution $\delta \lesssim 10$ am in its own x radiation with the aid of pinpoint cameras (quantum energy $\gtrsim 1-2$ keV). There was no correlation between the appearance of the jets and the changes in the spatial distribution of the x-ray luminosity of the plasma corona was observed. At the same time, a connection was observed between the spatial position, on the target, of the regions of jet generation, on the one hand, and the localization of the corona regions in which the plasma luminosity at the second-harmonic frequency of the heating radiation increased noticeably, on the other. To investigate the plasma luminosity at the second-harmonic wavelength, the image of the target was projected through interference filters by a lens on photographic film. The dimensions of the increased-luminosity regions (Fig. 6) were comparable with the spatial resolution (~15 μm) of the registration system. Simultaneously with photography of the plasma in the second-harmonic light, we registered, in the same direction, its spectrum with spatial resolution. The second-harmonic spectrum consists of two components: basic, shifted towards the "red" side by several angstroms, and a long-wave pedestal. This line shape was interpreted in Ref. 14, where a connection was obtained between the appearance of the basic component and the linear transformation of the incident light wave into plasma waves in a spatially inhomogeneous plasma, while the long-wave component was attributed to effects of parametric turbulence. Investigations of the contour of the spectrum of the second-harmonic line over the target surface have shown that in the region of increased luminosity there is a noticeable increase in the intensity of the basic component compared with the other sections of the surface (Fig. 7).

3. DISCUSSION OF EXPERIMENTAL RESULTS

To interpret the described experiments it is necessary first to find a particle-acceleration mechanism capable of explaining the following effects:

1. The fast particles are not produced on the entire surface of the target, but in a limited number of lo-

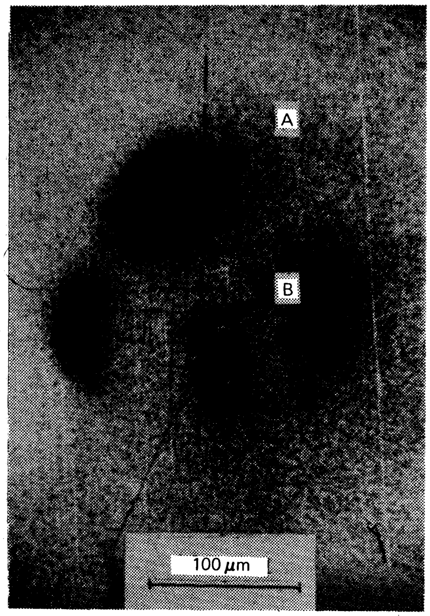

FIG. 6. Photographs of plasma corona in the light of the second harmonic of the heating radiation. The dashed line shows the position of the spectrograph slit.

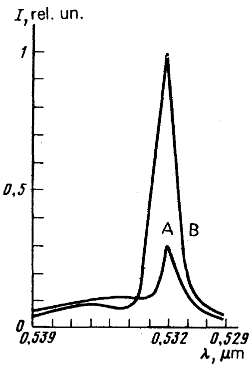

FIG. 7. Spectral distribution of the intensity of the 2ω harmonic at the two plasma-corona points A and B indicated by the arrows in Fig. 6.

calized zones.

2. The particles that ionize the gas around the target are emitted in almost parallel streams in directions that do not coincide with any of the heating beams.

3. The fast particles of high velocities exceeding by one order of magnitude the average thermal velocity of the spreading ions.

4. The group (jet) of accelerated particles carries an appreciable fraction of the energy, sufficient to produce a cylindrical shock wave in the gas surrounding the target.

5. A correlation was observed between the target regions from which the accelerated ions emerge and the regions where the second-harmonic of the heating radiation is generated most intensively.

We note first that the mechanism most frequently used to interpret laser experiments, namely electrostatic acceleration of the ions in the laser-plasma corona by the ambipolar potential of the "hot" electrons,[15,16] is incapable of explaining the observed anisotropic spreading of the ions. For the same reason we must reject also the assumption that the ions are accelerated on the plasma boundary[17]—it is seen from Fig. 2 that the plasma expands with spherical symmetry.

The narrowly directed group of fast ions could be produced by a cumulative jet formed when two shock waves converge in a dense plasma. However, strong anisotropy of the jet is preserved only over scales comparable with the dimensions of the converging shock waves. In our case, however, the jet is 100 times longer than the target dimension. In addition, the energy contained in a cumulative jet is much less than the energy of the converging waves. To explain the observed effect we must therefore assume that the energy of the colliding shock waves is of the order of the energy input to the target. Having monitored in the experiment the homogeneity of the target irradiation with a calorimetric system and x-ray pinpoint cameras, we can state that the possibility of formation of such powerful shock waves is low. In experiments with hollow shell targets, cumulative jets might be produced because of the asymmetry of the compression. However, the fact that the jets were observed also in experiments with solid spherical and flat targets allows us to reject this assumption, too. One more possible mechanism of formation of a jet of accelerated ions is connected with the presence of the gas surrounding the target. A powerful electron beam, produced for some reason in the corona of the laser plasma and entering the neutral gas, ionizes the latter and some of the ions acquire in this case velocities of the order of the ionization-wave propagation velocity. This effect is used for collective acceleration of ions.[18] In our case, however, this effect cannot occur for two reasons. First, anisotropic plasma expansion was registered in experiments in vacuum by time-of-flight collector measurements. Second, the necessary strong-current electron beam cannot be produced in the target because the needed number of electrons exceeds the total number of electrons in the target. In fact, in a glass shell of radius ~50 μm and a wall thickness ~1.5 μm contains $3\cdot 10^{15}$ atoms with average charge $z\sim 10$, i.e., $\sim 3\times 10^{16}$ electrons. On the other hand, knowing the energy released in a cylindrical shock wave and the energy lost per unit length by D^+ ions with velocity $(1.5-3.5)\times 10^8$ cm/sec (energy 25–100 keV), we find that the number of ions necessary to produce such a shock wave is 10^{13}. Recognizing that the electrons in the laser plasma can have an energy not higher than 100 keV, and that the coefficient of conversion of electron-beam energy into ion-beam energy, even at optimal pressure (0.1–0.5 Torr for deuterium) does not exceed 10^{-4}, we find that the minimum necessary number of electrons in the beam (10^{17}) exceeds the total number of electrons in the target.

Finally, the spontaneous appearance, in a laser plasma, of localized regions with a strong magnetic field (magnetic "islands") can also be the cause of ion acceleration. The collision of two such islands with oppositely directed magnetic force lines converts the magnetic energy into kinetic energy of the medium. Such an induction mechanism (reconnection of magnetic force lines) was used to explain the acceleration of the plasma in the atmosphere of the sun and the stars).[19] For a laser plasma this effect is insignificant, for according to Ref. 19 the reclosing of the magnetic force lines accelerates the plasma to velocities of the order of the Alfven velocities:

$$v_i\sim v_A\sim B/(4\pi n_i M_i)^{1/2}.$$

The same estimate can be obtained by equating the energy of the magnetic field of the island to the kinetic energy contained in this volume of plasma. Recognizing that the most probable location for the formation of the magnetic island is the region of critical density,[20] we find from the foregoing formula that the magnetic field needed to attain velocities $v_i\sim 3\times 10^8$ is $B\sim 100$ MG. It is not very likely that magnetic fields of such strength will be produced in a laser plasma.

Besides the ion-acceleration mechanisms listed above, under laser-plasma conditions there is another possible mechanism connected with resonant amplification of the electric field of the light wave in the vicinity of the critical density. It will be shown below that it is precisely this mechanism that can explain the experimentally observed jet acceleration of ions.

4. RESONANT ACCELERATION OF IONS IN A LASER PLASMA

The effect of resonant acceleration of ions in the vicinity of the region of the critical density, in the linear approximation in the amplitude of the electric field of the pump radiation, has been discussed in the literature. The main theoretical premises were formulated by Silin.[4] The energy distribution of the fast ions in a laser plasma, at not too large pump-flux energies, was obtained in Ref. 21. Physically, the cause of the resonant acceleration is that in the case of oblique incidence of the electromagnetic wave on an inhomogeneous plasma layer the wave electric-field component in the incident-wave plane (the p component) increases sharply at the plasma resonance point $\omega_0=\omega_{Le}$. This singularity is due to the resonant excitation of longitudinal Langmuir oscillations of the plasma by the p component of the pump field.[22]

In the linear approximation, the maximum field amplitude E_* at resonance is determined by the relation[22,23]

$$E_*\approx E_0\Phi(\theta)(2\pi k_0 L)^{-1/2}(\omega_0/\nu_{\text{eff}}), \tag{1}$$

where E_0 is the amplitude of p-polarized pump wave in vacuum, Φ is a resonant function of the wave incidence angle θ and reaches a maximum $\Phi(\theta_r)\approx 1$ at $\theta\sim\theta_r\approx(k_0L)^{-1/2}$, $k_0=\omega_0/c=2\pi/\lambda_0$, L is the characteristic scale of the inhomogeneity of the density in the vicinity of the critical point. The quantity ν_{eff} is determined either by the frequency of the Coulomb collisions ν_{ei}, or by the drift of the plasma waves out of the resonance region:

$$\nu_{\text{eff}}=\max\{\nu_{ei},\ \omega_0(r_{De}/L)^{2/3}\}, \tag{2}$$

where r_{De} is the Debye radius of the electrons.

The width Δx of the resonant peak is determined by the effective collision frequency

$$\Delta x\approx L\nu_{\text{eff}}/\omega_0. \tag{3}$$

In the case of a sufficiently hot plasma, however, when

$$\nu_{ei}<\omega_0(r_{De}/L)^{2/3}, \tag{4}$$

the region of existence of intensive Langmuir oscillations with wavelength $\sim\Delta x$ is much wider than (3) since the Langmuir waves that go out of resonance can remove energy to a distance

$$\Delta x_L\approx 3(k_{st}r_{De})^2 L \tag{5}$$

into less dense plasma layers, where they are absorbed by the electrons on account of the Cerenkov effect. For a Maxwellian electron-velocity distribution we have $k_{st}\approx 0.3\,r_{De}^{-1}$.

The excess pressure of the electric field of the Langmuir oscillations causes the plasma to be ejected from the region of the critical density. The maximum energy acquired by the accelerated ions is given by the relation[4]

$$n_i M_i v_i^2\approx E_*^2/4\pi. \tag{6}$$

The region of applicability of the presented formulas of the linear theory can be obtained by recognizing that the resonantly excited peak of the Langmuir waves can be parametrically unstable. A stability check on the peak shows that the nonlinear restructuring of the field can be neglected if

$$\frac{E_0^2}{4\pi n_c T_e} \leqslant 2\pi k_0 L \Phi^{-2}(\theta) \left(\frac{\nu_{\text{eff}}}{\omega_0}\right)^3 \max\left\{1, \frac{\nu_{\text{eff}}}{\omega_0}\left(\frac{\omega_{Le}}{\omega_{Li}}\right)^2\right\}. \qquad (7)$$

We note that allowance for the electronic nonlinearity (self-crossing of the trajectories)[24] leads to a greater restriction on the applicability of the linear theory than relation (7) only in a strongly inhomogeneous plasma, when $L < r_{De}\omega_{Le}/\omega_{Li}$. This inequality under real conditions of a plasma produced by a neodymium laser ($\lambda_0 = 1$ μm) is not satisfied at fluxes $q < 10^{16}$ W/cm^2.

The development of an aperiodic parametric instability in the resonant region leads to further increase of the amplitude of the electric field and by the same token to an increase of the energy of the accelerated ions (6). To determine the maximum electric field intensity at nonlinear resonance it is necessary to trace the evolution of the nonlinear stage of the aperiodic instability. We determine first the scale that characterizes the produced aperiodic instability. We confine ourselves here only to the case of sufficiently steep density profiles

$$L < r_{De}\omega_{Le}^2/\omega_{Li}\nu_{\text{eff}},$$

when $\Delta x/r_{De} < \omega_{Le}/\omega_{Li}$. In this case, when the threshold (7) is exceeded, a hydrodynamic aperiodic instability is excited in the resonant region Δx_L, with a characteristic scale $k_a - 1$ determined from the relation[25]

$$k_a r_{De} \approx \frac{\omega_{Li}}{\omega_{Le}} \frac{\Delta x}{r_{De}} \frac{E_*}{(4\pi n_c T_e)^{1/2}}. \qquad (8)$$

We note that the characteristic scale of the parametric instability k_a^{-1} (8) is less than the width Δx (3) of the peaks. As a result of the hydrodynamic aperiodic instability the resonant peaks therefore break up into several density wells in which the Langmuir oscillations are trapped.

The subsequent evolution of the cavitons produced by the instability can be considered with the interaction between them neglected. To solve our problem, that of determining the maximum energy of the accelerated ions, we do not have to trace the entire dynamics of the self-contraction of the cavitons, and it suffices only to obtain an estimate of the maximum energy of the Langmuir field in the caviton. We assume, following the results of Refs. 26 and 27, that the total energy of the trapped Langmuir field changes little in the course of the nonlinear evolution of the caviton.

The minimum size $\Delta x_{\min}$ to which the caviton is compressed is determined from the condition that the trapped Langmuir field not be subject to Landau damping. Consequently, when the caviton reaches the size $\Delta x_{\min} \sim k_{st}^{-1}$ the self-compression processes stops and effective absorption of the Langmuir oscillations takes place. From the condition for energy conservation in the caviton we obtain the maximum field intensity in the nonlinear regime:

$$E_{max} \approx E_* (k_{st}/k_a)^{3/2}. \qquad (9)$$

Under the high-temperature plasma conditions of interest to us, when the inequality (4) is satisfied, we obtain from (9) the following estimate of the maximum energy density of the Langmuir field in the caviton:

$$\frac{E_{max}^2}{4\pi} \sim \frac{2\pi}{\Phi(\theta)} (k_{st} r_{De})^3 \frac{(n_c T_e)^{1/2}}{E_0} \left(\frac{\omega_{Le}}{\omega_{Li}}\right)^3 \left(\frac{v_{Te}}{c}\right)^{5/3} (k_0 L)^{-1/6}. \qquad (10)$$

The total number of cavitons $N \sim \Delta x_L k_a$ can be quite large under the conditions (4) because of the large dimension of the region Δx_L:

$$N \sim k_a \Delta x_L \sim \frac{3(k_{st} r_{De})^2}{(2\pi)^{1/2}} \Phi(\theta) \left(\frac{c}{v_{Te}}\right)^2 (k_0 L)^{1/2} \frac{\omega_{Li}}{\omega_{Le}} \frac{E_0}{(4\pi n_c T_e)^{1/2}}. \qquad (11)$$

The caviton lifetime $\tau \sim (k_a v_s)^{-1}$ is determined by the self-compression time, since the "burnup" of the high-frequency content of the caviton occurs within a time $\Delta t \ll \tau$ (here v_s is the speed of sound).

We determine now the energy that can be acquired by an ion passing through such a structure consisting of N nonstationary cavitons. The energy determined by (6) is acquired only by ions initially present in the caviton. On passing through the caviton, which can be regarded as a nonstationary potential barrier, the ion acquires an energy increment $\Delta\varepsilon_i \equiv \Delta(M_i v_i^2)$ determined by the energy increment of the barrier during the time of flight $\Delta t_i \sim (k_a v_i)^{-1}$, i.e.,

$$\Delta\varepsilon_i \sim (\Delta t_i/\tau)(E_{max}^2/4\pi n_i) \sim (v_s/v_i)(E_{max}^2/4\pi n_i).$$

Thus, on passing through one caviton the ion draws from the caviton only a small fraction of its energy, but the total energy of the ion passing through all N cavitons turns out to be quite large:

$$\varepsilon_i \sim z T_e (N E_{max}^2/4\pi n_c T_e)^{2/3}. \qquad (12)$$

If the process of nonlinear amplification of the field (aperiodic instability) develops simultaneously in the entire region of x_L (5), then an ion with velocity $v_i = (\varepsilon_i/M_i)^{1/2}$ increases its energy during the time $\tau \sim (k_a v_s)^{-1}$ of the caviton collapse not in all the N cavitons (11), but in smaller number

$$N_{\text{eff}} \sim v_i/v_s \sim (\varepsilon_i/zT_e)^{1/2} \sim E_{max}(4\pi n_c T_e)^{-1/2}.$$

However, the group velocity of the Langmuir waves that leave the plasma-resonance region Δx (3) turns out to be comparable with the velocity of the fast ions accelerated in the first resonant peak. Therefore the compression of the cavitons in more rarefied plasma layers (at distances $\sim\Delta x_L$ from the resonance) should set in later than at resonance, for example at the time when ions accelerated in the first cavitons arrive in this region. The delay of the development of the aperiodic instability in the rarefied regions of the plasma then ensure an increase of the ion energy in all N [see (11)] cavitons.

The considered ion-acceleration mechanism results in a sharply anisotropic distribution of the accelerated particles, due to the spatial structure of the nonlinear field in the caviton. Indeed, according to the results of Ref. 28, a contracting caviton is a disk-shaped capacitor with rounded edges, the distance between whose "plates" is substantially less (by a factor 3–5) than their transverse dimensions. The orientation of the cavitons in space is determined by the Langmuir-wave polarizatin vector along which the growth rate of the aperiodic instability is maximal. Since the Langmuir waves propagate strictly along the density gradient, the planes of the plates of all the cavitons are parallel to one another and the ions accelerated in the direction of the plasma inhomogeneity have a small angle spread determined by the ratio of the ion temperature to their maximum energy:

$$\Delta\theta\sim(T_i/\varepsilon_i)^{1/2}.$$

Some of the ions are accelerated and escape through the lateral surfaces of the cavitons. By virtue of the caviton geometry, however, their number is smaller by a factor 3–5 than the number of ions accelerated in the inhomogeneity direction, and the angle spread is quite small.

The acceleration takes place within a very short time, of the order of the caviton collapse time. There is little likelihood that this process will be repeated, since the striction forces that arise in the resonant region should lead to a strong deformation of the density profile and by the same token to a complete restructuring of the electric field. We note that the obtained expression (12) for ε_i depends little on the pump-wave field intensity, since $N \propto E_0$ and $E^2_{max} \propto E_0^{-1}$.

To find the total number of ions that take part in the acceleration process it is necessary to determine the volume of the plasma in which the cavitons exist. The longitudinal dimension (along the inhomogeneity direction) is determined by the quantity of Δx_L (5). The area S of that spot on the critical surface at all points of which the incidence angle θ of the radiation is close to the resonant value $\theta_r \simeq (k_0L)^{-1/3}$, under conditions when the diameter of the laser-beam cross section in the target plane exceeds the diameter of the target and $k_0L \gg 1$, can be set equal to $S\approx (R_c\theta_r)^2$, where R_c is the curvature radius of the critical surface. The total number of accelerated ions in the volume is

$$N_i\approx\frac{3}{z}(k_{st}r_{De})^2 n_c L R_c^2 (k_0L)^{-2/3}. \qquad (13)$$

This formula should be regarded as an upper-bound estimate of the number of fast ions, since not all the ions acquire the maximum energy (12). The product $N_i\varepsilon_i$ yields the upper-bound estimate of the energy carried away by the jet:

$$N_i\varepsilon_i\sim T_e n_c R_c^2 \lambda_0 (k_{st}r_{De})^{16/3}(c/v_{Te})^{2/9}(\omega_{Le}/\omega_{Li})^{4/3}(k_0L)^{5/9}. \qquad (14)$$

The dependence of the jet energy on the pump-field intensity enters formula (14) implicitly, only via the hydrodynamic parameters: $N_i\varepsilon_i \propto R_c^2 L^{5/9} T_e^{8/9}$.

For the plasma parameters in the discussed experiment ($T_e \approx 0.6$–1 keV, $L\sim 20$ μm, $R_c \sim 100$ μm, $z\sim 10$, $A\approx 20$), according to formula (7), the nonlinear processes in the plasma resonance should set in at $q_0 > 2\cdot 10^{13}$ W/cm^2. The fluxes $q\sim 10^{14}$ W/cm^2 used in the experiment exceed this threshold value. According to (12), the maximum energy of the resonantly accelerated ions under the conditions of the discussed experiment is

$$\varepsilon_i\sim 20 z^{1/3} A^{4/3} T_e^{8/9} (L/\lambda_0)^{2/9}\,\text{keV}, \qquad (15)$$

(T_e is in keV), and at the parameters quoted above it yields a value $\varepsilon_i \sim 1$ MeV, which agrees with the measurement data. We note that according to formula (15) the ion energy shou'd increase with increasing atomic weight A of the target material and with increasing charge z.

The number of accelerated ions (13) $N_i \sim 10^8(L^{1/3}R_c^2/z\lambda_0^{4/3})$ (L, R_c, λ_0 are in microns) under the conditions of the experiment in question is ~10^{12}, which agrees within the limits of experimental error with the experimentally obtained value of N_i.

5. CONCLUSION

The mechanism proposed by us for the acceleration of ions in a laser plasma can explain all the phenomena listed above (see Sec. 3): 1) The onset of intense electric fields is limited to small sections of the critical surface, where the angle of incidence of the radiation on the plasma is close to optimal. 2) The accelerated ions are emitted in a direction normal to the target surface, and the angle spread turns out to be quite small—either of the order of the ratio of the velocity of the thermal ion to their maximum velocity, or of the order of the angle of curvature of the critical surface in the resonant region

$$\Delta\theta\sim\max\{(T_i/\varepsilon_i)^{1/2},\ (k_0L)^{-1/3}\}.$$

3) The energy of accelerated ions turns out to be quite high, larger by almost two orders of magnitude than the average thermal energy of the expanding plasma. 4) The energy of the flux of the accelerated ions is sufficient to ionize the gas around the target. The resonant acceleration of the ions by the ponderomotive force of the high-frequency electric field should be accompanied by intense generation, by the plasma, of harmonics of the frequency of the heating radiation, not of only the second harmonic observed in this experiment, but also of higher harmonics: third, fourth, etc.

It should be noted that with increasing laser-radiation flux density the energy (15) of the ions accelerated in the plasma resonance should not increase substantially. However, the directivity pattern of the fast ions should broaden because the density profile of the plasma becomes steeper in the critical region. Fast C VI ions with energy ~0.5 MeV were observed in experiment[29] at neodymium-laser radiation fluxes ~10^{15}–10^{16} W/cm^2. The energy and emission direction of these ions can be explained within the framework of the acceleration mechanism considered above.

[1]G. H. McCall, F. Young, A. W. Ehler, I. F. Kephart, and R. P. Godwin, Phys. Rev. Lett. 30, 1116 (1973).

[2]N. G. Basov, V. A. Boĭko, S. M. Zakharov, O. N. Krokhin, Yu. A. Mikhaĭlov, G. V. Sklizkov, and S. I. Fedotov, Pis'ma Zh. Eksp. Teor. Fiz. 18, 314 (1973) [JETP Lett. 18, 184 (1973)].

[3]J. Martineau, M. Rabeau, J. L. Bocher, J. P. Elie, and C. Patou, Opt. Commun. 18, 347 (1976).

[4]V. P. Silin, Pis'ma Zh. Eksp. Teor. Fiz. 21, 333 (1975) [JETP Lett. 21, 152 (1975)].
[5]N. G. Basov, Yu. A. Zakharenkov, N. N. Zorev, A. A. Kologrivov, O. N. Krokhin, A. A. Rupasov, G. V. Sklizkov, and A. S. Shikanov, Zh. Eksp. Teor. Fiz. 71, 1788 (1976) [Sov. Phys. JETP 44, 44, 938 (1976)].
[6]N. G. Basov, A. A. Kologrivov, O. N. Krokhin, A. A. Rupasov, A. S. Shikanov, G. V. Sklizkov, Yu. A. Zakharenkov, and N. N. Zorev, in: Laser Interaction and Related Plasma Phenomena, Plenum Press, New York, 1976, Vol. 4.
[7]Yu. V. Afanas'ev, N. G. Basov, O. N. Krokhin, V. V. Pustovalov, V. P. Silin, G. V. Sklizkov, V. T. Tikhonchuk, and A. S. Shikanov, Vzaimodeĭstvie moshchnogo lazernogo izlucheniya s plazmoĭ (Interaction of High-Power Laser Radiation with a Plasma), Itogi Nauki, Ser. Radiotekh., 17, Izd. VINITI, 1978, Chap. VII.
[8]Yu. A. Zakharenkov, A. V. Rode, G. V. Sklizkov, S. I. Fedotov, and A. S. Shikanov, Kvantovaya Elektron. (Moscow) 4, 815 (1977) [Sov. J. Quantum Electron 7, 453 (1977)].
[9]A. A. Erokhin, Yu. A. Zakharenkov, N. N. Zorev, G. V. Sklizkov, and A. S. Shikanov, Fiz. Plazmy 4, No. 3 (1978) [Sov. J. Plasma Phys. 4, No. 3 (1978)].
[10]O. N. Krokhin, A. S. Shikanov, G. V. Sklizkov, and Yu. A. Zakharenkov, Proc. of the Thirteenth Internat. Conf. on Phenomena in Ionized Gases, Sept. 1977, Berlin, DDR, p. 883.
[11]L. I. Sedov, Metody podobiya i razmernosti v mekhanike (Similarity and Dimensionality Methods in Mechanics), Gostekhizdat, 1954.
[12]N. G. Basov, E. Volovski, E. Voryna, S. Denis, Yu. A. Zakharenkov, S. Kaliski, G. V. Sklizkov, J. Farny, and A. S. Shikanov, Preprint Fiz. Inst. Akad. Nauk, No. 194, 1978.
[13]O. F. Nemets and Yu. V. Gofman, Spravochnik po yadernoĭ fizike (Nuclear Physics Handbook), Naukova dumka, Kiev, 1975.
[14]N. G. Basov, V. Yu. Bychenkov, O. N. Krokhin, A. A. Rupasov, V. P. Silin, G. V. Sklizkov, A. N. Starodub, V. T. Tikhonchuk, and A. S. Shikanov, Generatsiya vtoroĭ garmoniki v lazernoĭ plazme (Generation of Second Harmonic in a Laser Plasma), Preprint Phys. Inst. Acad. Sci. No. 196, 1978.
[15]E. J. Valeo and I. B. Bernstein, Phys. Fluids 19, 1348 (1976).
[16]K. A. Brueckner and R. S. Janda, Nuclear Fusion 17, 1265 (1977).
[17]A. V. Gurevich, L. V. Pariĭskaya, and L. P. Pitaevskiĭ, Zh. Eksp. Teor. Fiz. 49, 647 (1965) [Sov. Phys. JETP 22, 449 (1966)].
[18]A. A. Kolomenskiĭ, V. M. Likhachev, I. V. Sinil'shikova, O. A. Smit, and V. N. Ivanov, Preprint Phys. Inst. Acad. Sci. No. 44, 1975.
[19]S. I. Syrovatskiĭ, Izv. Akad. Nauk SSSR Ser. Fiz. 39, 359 (1975).
[20]K. Estabrook, Phys. Fluids 19, 1733 (1976).
[21]D. Baboneau, G. Di Bona, P. Chelle, M. Decroisett, and J. Martineau, Phys. Lett. A 57, 247 (1976).
[22]N. G. Denisov, Zh. Eksp. Teor. Fiz. 31, 609 (1956) [Sov. Phys. JETP 4, 544 (1957)].
[23]V. L. Ginzburg, Rasprostranenie elektromagnitnykh voln v plazme (Propagation of Electromagnetic Waves in a Plasma), Nauka, 1967. [Pergamon].
[24]N. S. Erokhin, V. E. Zakharov, and S. S. Moiseev, Zh. Eksp. Teor. Fiz. 56, 179 (1969) [Sov. Phys. JETP 29, 101 (1969)].
[25]T. A. Davydova and K. P. Shamraĭ, Preprint IFT Akad. Nauk Ukr. SSR, ITF-77-140R, Kiev, 1978.
[26]V. V. Gorev, A. S. Kingsep, and L. I. Rudakov, Izv. Vyssh. Uchebn. Zaved. Radiofiz. 19, 691 (1976).
[27]A. A. Galeev, R. Z. Sagdeev, V. D. Shapiro, and V. I. Shevchenko, Zh. Eksp. Teor. Fiz. 73, 1352 (1977) [Sov. Phys. JETP 46, 711 (1977)].
[28]L. M. Detyarev, V. E. Zakharov, and L. I. Rudakov, Zh. Eksp. Teor. Fiz. 68, 115 (1975) [Sov. Phys. JETP 41, 57 (1975)].
[29]E. A. McLean, R. Decoster, B. H. Ripin, J. A. Stamper, and S. E. Bodner, Appl. Phys. Lett. 31, 9 (1977).

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