Measurements of the dependence of the degree of screening η (defined by us as the ratio of the incident flux I_0 to the transmitted flux I in the time interval in which the transmitted flux is minimal) on the density of the radiation power incident on the target are shown in Fig. 3. The transmitted flux is expressed in terms of the incident flux and the power density W on the target by $I_0/I = A \exp(W/W_0)$, where A and W_0 are constants.

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DAMPING AND AMPLIFICATION OF LANGMUIR SOLITON

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Submitted 14 February 1973
ZhETF Pis. Red. 17, No. 7, 344 - 348 (5 April 1973)

The damping of a high-frequency Langmuir soliton in a Maxwellian plasma and the amplification of such a soliton by an electron beam were investigated theoretically and experimentally.

In a plasma placed in a magnetic field there can be produced a high-frequency Langmuir soliton propagating at an angle to the field ($k_\perp \neq 0$) with phase velocity $v_{ph} > v_{ph}^{cr} = \omega_{pe}/k_\perp(1 + \omega_{pe}^2/\omega_{He}^2)$ [1 - 3].

We have investigated experimentally the resonant interaction of such a soliton with plasma particles, whereby the soliton amplitude is adiabatically altered. In a Maxwellian plasma, interaction with the particles reflected from the "hump" of the potential causes a power-law damping of the soliton amplitude [4].

$$\alpha(z) = \frac{\alpha(0)}{\left(1 + \sqrt{\frac{\alpha(0)}{6\pi}}\,\frac{\gamma_L z}{v_{ph}}\right)^2}, \quad \gamma_L = -\frac{\omega_{pe}}{n_0} v_{ph}^2 \frac{\partial f_0}{\partial v_{ph}}, \tag{1}$$

$f_0(v)$ is the equilibrium distribution function of the plasma particles, $\alpha = e|\phi_0|/mv_{ph}^2$, and ϕ_0 is the amplitude of the potential; the conditions for the applicability of formula (1) are

$$\sqrt{\frac{e|\phi_0|}{m}}\, v_{ph} << v_{Te}^2 \text{ and } \omega_{pe} << \omega_{He}.$$

If the plasma contains a beam moving faster than the soliton, $v_0 > v_{ph}$, then the same mechanism of interaction with the resonant particles leads to amplification of the soliton. For a beam with a smeared-out velocity distribution $\Delta v/v >> (n_1/n_0)^{1/3}$ (n_1 and n_0 are respectively the beam and plasma densities), we have adiabatic amplification of the soliton without a change in its shape. The growth of the soliton amplitude is then given by

$$\frac{1}{2\sqrt{\Delta}}\ln\left[\frac{\sqrt{\Delta}-\sqrt{\alpha(0)}}{\sqrt{\Delta}-\sqrt{\alpha}}\;\frac{\sqrt{\Delta}+\sqrt{\alpha}}{\sqrt{\Delta}+\sqrt{\alpha(0)}}\right]+\frac{1}{\sqrt{\alpha(0)}}-\frac{1}{\sqrt{\alpha}}=$$

$$=\frac{16}{3^{3/2}\pi^2}\,\frac{\omega_{pe}z}{(\Delta v)^3}\,\frac{n_1}{n_0}\,v_{ph}^2,\quad \Delta=\frac{3\pi}{4}\,\frac{v_0-v_{ph}(0)}{v_{ph}(0)} \tag{2}$$

The amplitude of the potential in the soliton increases to a value

$$\phi_0^{max}=\frac{mv_{ph}^2}{e}\Delta \lesssim \frac{3\pi}{4}\sqrt{\frac{mv_{ph}^2}{e}\phi_0(0)}, \tag{3}$$

corresponding to an electric field with an energy density

$$\frac{E_{\perp max}^2}{4\pi}\approx\frac{9\pi^2}{16}n_0 m(v_0-v_{ph}(0))^2 \tag{4}$$

which is $(n_0/n_1)^{2/3}$ times larger than in the case of excitation of a monochromatic wave by an electron beam. The increase is due to the fact that in the case of the soliton there is no mechanism for nonlinear stabilization of the instability of phase "mixing" of the resonant particles (see [4]).

We note that capture of ions by the field of a high-frequency Langmuir soliton leads to formation of charge bunches that are stable in the radial and longitudinal directions, and can be accelerated together with the soliton. Such an acceleration method was proposed already in [5]. It is possible that the accelerated-ion pulse observed in the experiment of [6] is produced by such a mechanism.

The experiments were performed with a proton plasma of density $n_0 = (1 - 2) \times 10^7$ cm^{-3} and electron temperature $T_e \lesssim 4$ eV; the plasma was placed in a strong magnetic field of intensity 1 - 3 kOe ($\omega_{He} >> \omega_{pe}$). Just as in [2], the high-frequency Langmuir soliton was excited with the aid of a grid, to which a negative potential with a rise time ~10 nsec was applied, and we investigated the propagation of such a soliton in a plasma waveguide in which the section with homogeneous magnetic field was 40 cm long. The diverging magnetic fields at the ends of the setup reduced considerably the wave reflection from the ends. The wave propagation and plasma drift directions were opposite, with $v_{dr} < v_{ph}$, so that the presence of the drift did not lead to a noticeable change in the dispersion characteristics of the plasma waves. An investigation of these characteristics at small exciting-signal amplitudes has shown that the wave propagating in the system is the first radial harmonic of the Langmuir oscillations

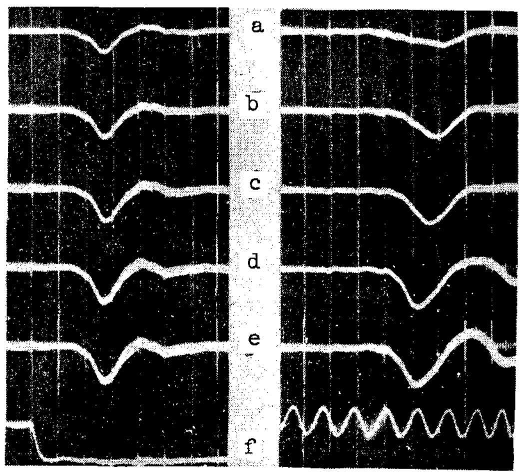

Fig. 1. Oscillograms of the adiabatic damping of a Langmuir soliton in a plasma and of its amplification by an electron beam: a) longitudinal soliton damping in length: z = 0, z = 20 cm. b - d) longitudinal soliton amplification: z = 0, z = 20 cm (b, c, d, and e correspond to $v_0 = 1.9 \times 10^8$, 2.3×10^8, 3×10^8, and 4×10^8 cm/sec, respectively; f) negative grid potential and 20 MHz time markers).

of a magnetized plasma with maximum phase velocity

$$v_{ph}^{cr} = \frac{\omega_{pe}}{k_\perp} \approx 1.2\cdot 10^8 \text{ cm/sec.}$$

When a negative step potential with voltage larger than 5 V was applied to the exciting grid, a single pulse of negative polarity was excited in the system; the width of this pulse, in accordance with the theoretical estimates [1 - 4], is given by

$$\frac{\Delta k_\perp}{2} = (1 - M^{-2})^{-1/2}, \quad M = \frac{v_{ph} k_\perp}{\omega_{pe}}, \quad M \text{ – Mach number.} \tag{5}$$

As a result of resonant interaction with the plasma particles, the soliton becomes damped as it propagates into the system. The damping is adiabatic, so that relation (5) is approximately satisfied at all values of the amplitude. The dynamics of the soliton during the course of its propagation in the system was investigated with high-frequency capacitive probes, the signals from which were fed to an oscilloscope through an integrating network and a broadband amplifier. We used for this purpose matched capacitive probes spaced 10 cm apart along the setup and oriented parallel to the transverse component of the high-frequency magnetic field. Figure 1 shows a series of oscillograms for the case of soliton damping (oscillograms a) and its amplification by an electron beam (oscillograms b - e). A beam of 10 μA current and energy up to 100 eV was produced with a two-electrode electron gun located several cm away from the exciting grid. The left-hand series of oscillograms shows the soliton registered by a probe located 5 cm away from the exciting grid. The right-hand series shows the oscillograms of the soliton after it passes 20 cm along the system. The parameter of oscillograms (b) - (e) is the electron-beam velocity v_0. From oscillograms of this type we were able to determine the spatial distribution of the soliton amplitude. Figure 2 shows the corresponding plot for the case of soliton damping, compared with the theoretical dependence obtained from formula (1) (the normalization is at the point z = 0).

The presence of an electron beam in the plasma causes amplification of the soliton. The soliton becomes amplified until its phase velocity exceeds the beam velocity. Accordingly, the maximum Mach number in the experiment M_{max} was proportional to the beam velocity (see Fig. 3). We also obtained plots of the spatial distribution of the amplitude of the amplified soliton at different values of the beam density and at a fixed beam energy. At low beam densities, $n_1 \ll n_0$, these curves agree with the theoretical ones obtained from formula (2). At high beam densities, $n_1 \sim n_0$, the beam alters the parameters of the plasma in which the soliton propagates, and the spatial growth increment of the amplitude

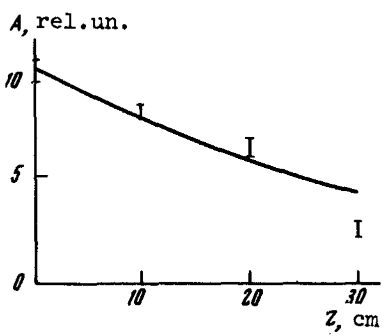

Fig. 2

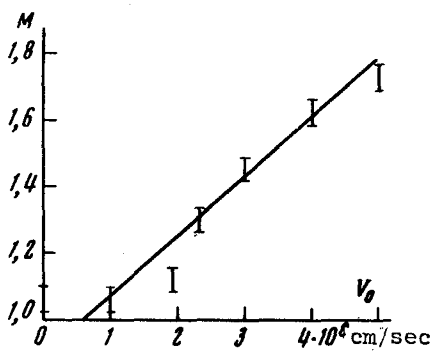

Fig. 3

Fig. 2. Variation of soliton amplitude with length.

Fig. 3. Dependence of solition velocity on the beam velocity

increases with increasing beam density more slowly than in (2).

The authors thank A.N. Antonov for help with the experiment.

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FLUSHING OUT THE LIQUID PHASE - A NEW MECHANISM OF PRODUCING A CRATER IN PLANAR FULLY DEVELOPED EVAPORATION OF A METALLIC TARGET BY A LASER BEAM

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Submitted 21 February 1973
ZhETF Pis. Red. 17, No. 7, 348 - 351 (5 April 1973)

We examine a new physical mechanism of crater development in a metal acted upon by laser radiation, wherein the liquid metal is ejected from the irradiated region by the reaction vapor-pressure gradient. The mechanism in question is decisive at not too high a radiation intensity. The obtained estimates agree with the experimental data.

In the case of fully developed metal evaporation by a laser beam, the entire energy of the radiation absorbed in the surface metal layer goes over into the work required to evaporate the material from this layer. The simplest to interpret are experiments in which the depth h of the crater produced by the radiation in the target is small in comparison with its diameter d, which is determined by the dimension of the irradiated area on the target surface. In this case we can neglect effects on the edges of the shallow crater and assume approximately that the evaporation is planar and one-dimensional. The conditions under which such a process can be observed are given by the inequalities [1]

$$d > ut > \sqrt{\chi t} > 1/\alpha. \qquad (1)$$

Here α is the metal absorption coefficient, χ is the temperature conductivity, t is time of action of the laser beam, and u is the velocity of the evaporation front and is determined from energy considerations.

$$I(1 - R) = \rho\lambda u. \qquad (2)$$

In Eq. (2), ρ is the density of the condensed phase, R is the reflection coefficient, and I is the beam intensity. According to (1) and (2), the crater depth $h = ut$ is determined by the mass of the removed vapor and depends linearly on the intensity, $h \propto I$.

It was shown in [2] that the metal disintegrates via a liquid-vapor phase transition, unlike in [1], where a solid-vapor transition was considered. This difference is fundamental and leads to new effects, one of which is discussed in the present paper.