Strong Langmuir turbulence excited by an electron beam in a plasma

D. M. Karfidov, A. M. Rubenchik, K. F. Sergeĭchev, and I. A. Sychev

General Physics Institute, Academy of Sciences of the USSR
(Submitted 21 February 1990; resubmitted 17 July 1990)
Zh. Eksp. Teor. Fiz. **98**, 1592–1604 (November 1990)

We study the excitation of strong Langmuir turbulence by injection of an electron beam into a laboratory plasma, by investigating the consequences such as the generation of short electromagnetic radiation bursts, the generation of short-wavelength ion sound with scales of a few Debye lengths, and the growth of accelerated electron "tails" in the plasma electron distribution function. We show the connection between those processes and the ion dynamics in cavitons and we establish a threshold for the transition to strong turbulence.

INTRODUCTION

By now, the following picture of strong Langmuir turbulence has evolved.[1,2] Density wells (cavitons) filled with Langmuir waves, occur spontaneously in the plasma as a result of the modulational instability. These cavitons are compressed after a finite time and the Langmuir waves remain trapped in them during the compression. At the end of the collapse stage the Langmuir field in the cavitons "burns out" and the energy of the oscillations is transferred to the electrons. As a whole, turbulence is an ensemble of collapsing cavitons which are at different stages of evolution. The elementary entity of the turbulence, the collapsing Langmuir caviton, has been well studied. The geometry of a caviton and the way it is compressed have been explained analytically and through numerical simulation of the averaged hydrodynamic equations of a simplified description variant.[1,2]

It has been possible to carry out for the concluding stage of the caviton compression a numerical simulation which requires a complete kinetic description.[3] The process of energy transfer to the electrons has been studied; the minimum size l of the caviton has been determined and turns out to be close to $l \sim 15 r_{De}$. Finally, in Ref. 4 it was possible to study the collapse process of an isolated caviton experimentally.

Much worse is the situation regarding the study of turbulence itself. Existing papers[2,5] have been based on the use of the above mentioned dynamic equations, which do not describe the concluding stage of the collapse or the particle acceleration process. It is unclear how adequately these effects have been taken into account through a series of modifications of the hydrodynamic equations. A complete kinetic simulation of turbulence by the particle method lies beyond the limits of possibilities of present-day computers. It is therefore of great interest to study strong plasma turbulence experimentally.

The transition from weak to strong turbulence leads to a significant change in its macroscopic manifestations. For instance, one assumes that in a strongly turbulent plasma the dissipation of the oscillations leads not to plasma heating, but only to the acceleration of a small group of electrons. This fact is extremely important, for instance, for the interpretation of experimental data about the interaction of strong laser radiation and matter. However, this statement has not been confirmed either experimentally or theoretically.

In theoretical papers the main attention has been paid to a study of the structure of the electric field of the caviton. This kind of experimental study is singularly difficult since the cavitons are very small and the collapse process is fast. As a rule, the electric field is measured by probes. If there are many probes, the picture of the electric field is strongly distorted. For a study of the structure of the field of a caviton by a single probe it is necessary that the picture of the electric field be reproduced from pulse to pulse.

In the successful experiments of Ref. 4 (see also Ref. 6) the picture of the field is fixed by a selection of the current density of the electron beam and by boundary conditions such that in the whole volume of the plasma a well reproducible caviton is produced, the structure of which was studied. It is clear that as a matter of principle this situation does not enable us to study the properties of the plasma turbulence.

Let us dwell upon yet another aspect of the study of strong plasma turbulence. After the experiments of Ref. 4 one might gather the impression that Langmuir collapse is an exotic phenomenon which occurs only under special conditions. On the other hand, in accordance of the theoretical ideas of Ref. 7, there should be, in fact, strong turbulence in the majority of experimental situations.

In the present paper we show that strong Langmuir turbulence is, indeed, realized in a rather ordinary plasma-beam system and that its properties agree with existing theoretical ideas.

The transition to a strong-turbulence regime was fixed by the appearance of electromagnetic radiation bursts in the plasma. One should state that such bursts had also been observed earlier in a number of experiments.[8,9] We carried out correlation measurements showing that the radiation comes from localized plasma regions which were interpreted as cavitons filled with Langmuir oscillations.

Another method consisted in using probes to measure the evolution of the electron distribution function. Experimentally we observed the process of the formation of tails on the distribution functions. The picture obtained confirms and complements measurements of the spectra of ion-sound turbulence and of the distribution function of the beam electrons which pass through collapsing cavitons.

The paper is constructed as follows. In the first section we give a description of the experimental setup and the plasma and beam parameters. We discuss the experimental methods used. In the second section we describe the results of the measurements of the electromagnetic radiation and their interpretation. The third section is devoted to probe measurements of the ion-sound oscillations of the turbulent

 0038-5646/90/110892-07$03.00

plasma. In the fourth section we discuss the evolution of the distribution function of the electrons in the beam and in the plasma.

1. EXPERIMENTAL CONDITIONS AND METHODS OF THE MEASUREMENTS

The experiments were carried out with a collapsing-plasma column produced by a pulsed plasma-beam discharge (Fig. 1a) in a weak longitudinal static magnetic field $H_0 \sim 100$ G in an argon or xenon medium with a pressure of $p_{\mathrm{Ar,Xe}} \approx 5 \times 10^{-4}$ Torr. The discharge arose as the result of the injection into the gas of a strong electron beam with a current of ≈ 12 A, generated by the thermocathode 1 with an emission layer of lanthanum hexaboride LaB_6 to which a negative voltage pulse with an amplitude of 3 kV and a length of up to 100 μs is applied. The launching anode was the tantalum grid 2 positioned 5 mm from the cathode. After the end of the pulse of the discharge current the plasma broke up as the result of diffusion and recombination processes. Parallel with this there was a "cooling" of the electrons and a drop in the amplitude of the plasma-density fluctuations. After approximately 100 μs the plasma density at the center of the column cross section was lowered approximately by a factor two and the temperature by approximately a factor five. The plasma density profile had for the discharge in argon (Fig. 1b) the shape of a plateau within a radius of 6 cm, beyond which there followed a rather steep drop, reaching its maximum slope at a radius of ≈ 10 cm. The drop in the electron temperature with radius is smoother: by a factor two at a radius of ≈ 15 cm. The electron density and temperature distributions along the axis of the plasma column were uniform.

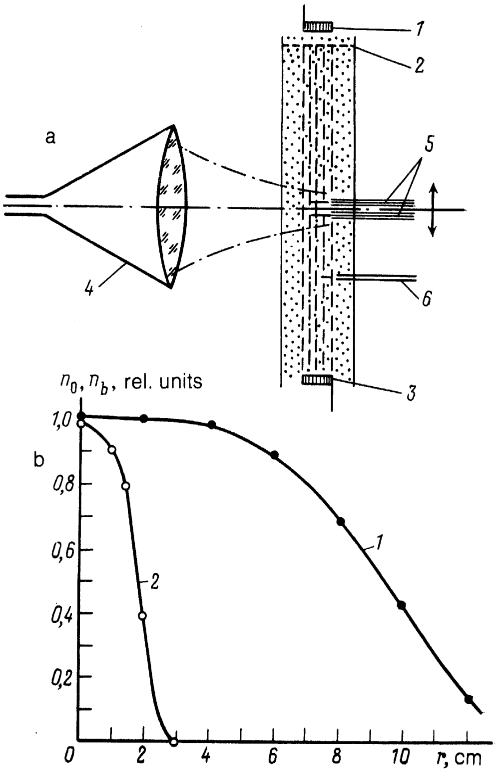

FIG. 1. a: Setup of the experiment: 1: thermoemission cathode for producing a plasma; 2: anode for the starting discharge current; 3: thermoemission cathode for injecting an electron beam into the plasma; 4: horn-lens receiving antenna for measuring radiation from the plasma; 5: loop antennae registering plasma fields; 6: Langmuir plasma probe. b: plasma density profile along the radius of the column (1) and electron beam density profile (2).

The excitation of Langmuir oscillations in the plasma was accomplished by pulsed injection of a low-voltage electron beam with an energy up to 300 V and a current up to 15 A from the auxiliary emission cathode 3 (Fig. 1a). The excitation of Langmuir oscillations was due to beam instability. At the time of the injection of the electron beam ($t = 100\,\mu$s) the plasma parameters were the following: electron density on the column axis $n_0 \leqslant 10^{12}$ cm^{-2}, electron temperature 2.5 eV, nonisothermy of electrons and ions $T_i/T_e \lesssim 0.1$, electron-density fluctuations $\delta n/n_0 \leqslant 5\cdot 10^{-3}$. The characteristic diameter of the electron beam corresponded approximately to the thermocathode diameter 3 cm. The plasma density within the limits of the beam cross section was therefore deemed to be homogeneous within 1–2%.

The electron pulse-duration 5 μs was chosen such that during the pump-pulse duration no noticeable ionization of the neutral gas was produced: $\Delta n \lesssim 10^{-2} n_0$.

The plasma characteristics were measured with ordinary Langmuir probes 6 (Fig. 1a). The probe wire diameter was 0.2 mm and the length 5 mm.

The distribution function of the beam electrons or of the accelerated-electron elelectrons was investigated with a multigrid probe of 10 mm diameter. The current of the electrons separated from the ions by the positively charged grid was further energy-analyzed by the retarding-potential method. To obtain the distribution function, the delay curves of the current to the collector probe were differentiated. The delay curve was recorded using a memory unit by applying to the analyzing grid of the probe a sawtooth voltage pulse duration 0.5 μs and amplitude up to -450 V. The distribution function of the bulk of the plasma electrons was determined from measurements of the electron branch of the current-voltage characteristic of the Langmuir probe, likewise by differentiation. The analysis of the thermal electrons of the plasma in the energy region from 0 to 10 V with the aid of a multigrid probe was incorrect because of the influence of the probe geometry.

The electromagnetic radiation of the plasma was recorded with a horn–lens antenna[4] (Fig. 1a) collecting the radiation from the focal region with a characteristic diameter of 5 cm into a waveguide with cross-section 1.5×3.5 cm^2, corresponding to the entrance to the receiver. The diameter of the pupil of the lens of was ≈ 30 cm and the distance from the focal plane to the lens ≈ 20 cm. The sensitivity of the receiver was $\sim 10^{-10}$ W and the amplifier bandwidth ≈ 10 MHz.

The horn-lens antenna did not, however, have a sufficient spatial resolution to determine the dimensions of the localized sources of the radiation from the plasma. To do this we used loop antennae inserted into the plasma. Each of these antennae is a circular thin-wire loop connecting the central conductor with the screen of a coaxial feeder with a wave resistance of 50 Ω. The radius R of the ring satisfies the relation $l \ll R\lambda_0/2\pi$, where $\lambda_0 = 2\pi c/\omega_{pe}$ is the wavelength of the emitted transverse wave, and l the size of the emitting source. Taking into account that such an antenna operates in the near zone of the local source one must consider it rather as a coupling element with the source of the plasma oscillations.

Regarding this source as a weakly radiating dipole with an azimuthal magnetic field component, one can show that the emf $\mathscr{E}$, produced in the contour of the loop by the magnetic flux intersecting it, depends logarithmically on the ratio of the loop radius R and the distance between the point dipole and the conductor of the loop: $\mathscr{E} \propto \ln(R/|\Delta\mathbf{r}|)$ when $\Delta\mathbf{r}|/R \ll 1$. The singularity as $\Delta\mathbf{r} \to 0$ is removed by the fact that the conductor has a finite thickness and by the perturbing action of the loop on the dipole field when $\Delta\mathbf{r}| \approx l$. The symmetric disposition of the two identical loop antennae, one parallel to the other, enables us to establish for small distances between them a correlation of the emf signals produced in the loops from electromagnetic radiation bursts and to assess the characteristic dimensions of the radiation sources from the signal correlation length.

The measurements of the spectral function of the plasma radio-emission near the plasma frequency were carried out by means of a SKCh-64 type spectral analyzer together with a S4-80, which enabled us to obtain a scan of the spectrum of a single pulse. The signal selection was varied with time within the range 0 to 0.5 μs. The spectra of the ion-sound oscillations of the plasma were measured both with an S4-80 spectral analyzer, and also by means of a tunable frequency filter.

2. ELECTROMAGNETIC RADIATION FROM THE PLASMA

The injection of an electron beam into the plasma is accompanied by the appearance of nonthermal electromagnetic radiation at a frequency close to the plasma frequency ω_{pe}. It is caused by the conversion of plasma oscillations into electromagnetic oscillations. In a uniform plasma this radiation may be connected with decay (see, e.g., Ref. 10) of a plasma wave into an electromagnetic and a sound wave:

$$\omega_k \to \omega_t + \omega_s. \tag{2.1}$$

Moreover, one can have conversion on density fluctuations which arise under the action of ponderomotive forces connected with the bunching of Langmuir oscillations, for instance, in collapsing cavitons. Registration of the radiation by means of the horn-lens receiving antenna showed that it depends significantly on the density n_b of the electrons in the beam. For small values of n_b there is no electromagnetic radiation, i.e., its intensity is lower than the sensitivity of the receiver, but for $n_b > n_{b\,\mathrm{thr}}$ there appear bursts with a frequency which increases with increasing beam density. We note that this kind of radiation bursts have been observed in a number of experiments.[8,9]

It is natural to identify these radiation bursts with their generation in collapsing cavitons. However, the nonstationarity of the radiation can be connected with the relaxation oscillations of the beam, its bunching, and the passage of electron bunches through the observation region. To elucidate the mechanism of the bursts we carried out correlation measurements by means of the loop probe antennae described earlier.

Measurements showed that the radiation sources are localized. For large distances between the probe-antennae, $L > 0.1$ cm, we did not observe correlations between the bursts. For small distances between the antennae we registered a mutual correlation for a significant number of bursts. Figure 2 illustrates the correlation of signals for two radiation bursts. The size of the sources is, clearly, less than the distance between the antennae for which the correlation between bursts disappeared. The results of the measurements showed that the size of the source is $l < 0.1$ cm ($100 r_{De}$) and depends weakly on the beam parameters. The length of the bursts also depends weakly on the beam parameters.

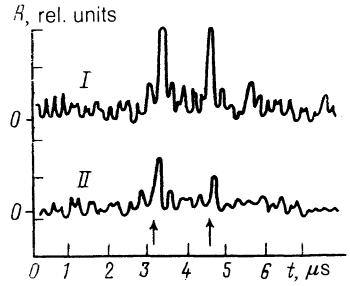

FIG. 2. Oscillograms of plasma field "bursts" at a frequency 4.8 GHz, obtained from two identical loop antennae with their loop planes separated by a distance of 0.1 cm.

The recording probe-antennae are in the near zone of the source at distances much less than the wavelength. The magnetic and electric fields of the source (in the dipole approximation) decrease here proportional to $|\Delta\mathbf{r}|^{-3}$, where $|\Delta\mathbf{r}|$ is the distance between the conductor of the probe loop and the source. The scatter of the signals from the probe in Fig. 2 is connected both with the scatter of the distance of the radiation sources, and with the random orientation of their dipoles.

We compared the length of single bursts for different gases, in particular, argon and xenon. It turned out that the length of the bursts increases with increasing ion mass. The length of a radiation burst, measured by the half-height, was $\tau = 8 \times 10^{-8}$ s in an argon plasma and $\tau = 1.7 \times 10^{-7}$ s for xenon, i.e., approximately inversely proportional to the ion plasma frequency.

We established that a stationary external magnetic field, which in the experiments was varied within the limits $60 \leqslant H_0 \leqslant 300$ G, did not show any appreciably effect on the burst dynamics. The experiment thus shows that the burst sources which are localized in space have small dimensions, less than 100 Debye radii, and are caused by the ion motion in the plasma. It is difficult to explain all these results by anything else but emission by collapsing cavitons. We now discuss the results obtained from this point of view.

The intensity of the dipole radiation generated by a caviton at a frequency $\omega \approx \omega_{pe}$ is equal to (see, e.g., Ref. 11)

$$I_\omega \sim \varepsilon(\omega) \mathscr{D}^2 / c^3,$$

where

$$\mathscr{D} = e \int \delta n_i v \, d^3 r. \tag{2.2}$$

Here $v = eE/m\omega_{pe}$ is the oscillatory electron velocity, δn_i the low-frequency variation in the ion density, and $\varepsilon = 1 - \omega_{pe}^2/\omega^2$ the plasma permittivity. Since ε is close to zero the intensity of the radiation depends significantly on the plasma density distribution around the caviton and it is therefore difficult to estimate the absolute value of the quantity I_ω. However, (2.2) enables us to give a comparative estimate of the contributions to the radiation from the different stages of the collapse.

Numerical calculations (see, e.g., Ref. 11) show that collapsing cavitons soon reach a self-similar compression regime. The variations in the ion density and in the electric field in the caviton then change as follows:

$$\delta n_i/n_i=(t_0-t)^{-4/3}V[2(t_0-t)^{-2/3}], \quad E=(t_0-t)^{-1}f[\mathbf{r}/(t_0-t)^{2/3}], \tag{2.3}$$

which after substitution into (2.2) leads to

$$I_\omega \propto E^2(\delta n_i)^2 \propto l^{-1} \propto (t_0-t)^{-2/3}, \tag{2.4}$$

i.e., the radiation appears mainly in the concluding, very short stage of the process. This also explains the "burstlike" nature of the radiation observed in the experiment. It is necessary to bear in mind that Eqs. (2.3) and (2.4) are inapplicable when we take dissipation in the concluding stage of the collapse into account for $t \approx t_0$. Simulation by the particle method[3] shows that the length of the bursts is $\tau \approx (5\text{–}10)\omega_{pi}^{-1}$, which for an argon plasma is close to 40 to 80 ns. This agrees well with the experimental value $\tau \approx 80$ ns, given above. The increase of the burst length, $\tau \propto M^{1/2}$, with increasing ion mass also agrees with the experiment.

We note that one sometimes observed bursts also after the beam was switched off with a delay up to 0.5 μs when the residual beam current at the moment the bursts appeared was already well below the threshold value of the modulational instability. This agrees with the time for developing the collapse $\tau \approx 50\omega_{pi}^{-1}$ and with ideas about the autonomous nature of this process in relation to the pump: the caviton is created at the end of the pump pulse and later goes over into the collapse phase. The time for the dissipation of turbulence is the time for the formation, collapse, and burning-out of the caviton, which is filled with oscillations. Since the length of the initial stage of the process is considerably longer than the burn-out time, the delay of the burst is considerably longer than its length. Unfortunately the absence of reliable estimates of the level of turbulence makes it impossible to determine more precisely the agreement between the time for the formation of a caviton with the calculated time.

An estimate of the intensity I_ω of the radio-emission bursts can be obtained directly from the experiment. Without having exact data about the orientation of the dipoles of the radiating cavitons we assume that the angular diagram of the radiation is spherically symmetric and in that case

$$I_\omega \approx \frac{4\pi}{\alpha} P_\omega, \tag{2.5}$$

where P_ω is the power of the radiation received by the horn-lens antenna from the solid angle α. The estimate of P_ω for maximum burst amplitudes gives a value of $\sim 10^{-6}$ W. The radiated fraction of energy from the Langmuir field captured by the caviton reaches, according to the present experiment, an order of a few percent.

The absence of any apprecable effect of the magnetic field on the burst dynamics can be explained by the fact that in the experiment the condition for the magnetic field to influence the collapse process,[2]

$$3k^2 r_{De}^2 < \omega_H^2/\omega_{pe}^2,$$

is not satisfied; here $k = \omega_{pe}/v_b$ is the wave number of the plasma oscillations excited by the beam and ω_H the gyrofrequency. This condition takes the form

$$kr_H < 0{,}6,$$

from which it follows that the effect of the magnetic field on the collapse process appears when the wavelength of the excited plasma waves has an order of magnitude which is larger than the magnitude of the Larmor radius r_H of electrons with a thermal velocity. In our experiments the opposite condition, $kr_H > 0.6$, held in the whole range over which the beam electron energy varied, for the indicated range of values of the magnetic field.

3. ION-SOUND PLASMA OSCILLATIONS

The deformation of the plasma density under the action of ponderomotive forces means that the caviton generates ion-sound oscillations. It follows from (2.3) that

$$\delta n_i/n_i \propto (r_{De}/l)^2, \tag{3.1}$$

i.e., sound is generated mainly in the concluding stage of the collapse so that the spectral density of the sound oscillation energy, $W_s \sim nT_e(\delta n/n_i)^2$, must have a maximum for $k \sim l_{\min}^{-1}$ and drops rapidly when k decreases, $W_s \propto k^4$.

In the framework of weak turbulence theory the wave numbers of the ion-sound oscillations cannot be much larger than the wave numbers of the oscillations excited by the beam. The ion-sound spectrum must decrease rapidly with increasing k. The transition from a weak- to a strong-turbulence regime must thus be accompanied by a qualitative restructuring of the ion-sound spectra with the formation of separate frequency peaks in their short-wavelength part. Such a picture was, in fact, observed in the experiment.

We show in Fig. 3 the spectral density of the ocillations in the probe-saturation ion current measured by means of a

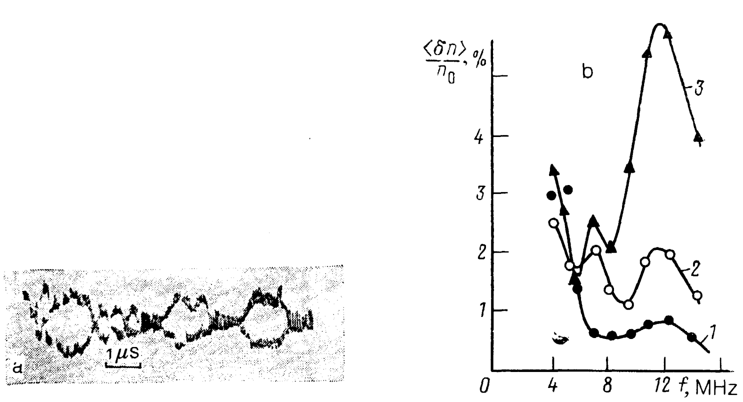

FIG. 3. Ion-sound oscillation spectra. a: oscillogram of the ion current of the saturation Langmuir probe passing through a resonance filter tuned at a frequency of 12 MHz; b: oscillation spectra for a beam energy $U_b = 300$ eV for three values of the electron beam density parameter: 1: $(n_b/n_0)_{\text{thr}} = 3.2\times10^{-3}$; 2: $2(n_b/n_0)_{\text{thr}}$; 3: $12(n_b/n_0)_{\text{thr}}$ (the Langmuir ion frequency is $f_{pi} \approx 16$ MHz).

frequency-tunable filter for different values of n_b/n_0. The minimum size $l_{\min}$ of the caviton, determined from the dispersion relation $\omega_s = \omega_{pi}(1 + k^2 r_{De}^2)^{-1/2}$ to be $l_{\min} \sim 2\pi c_s \omega_{s\max}^{-1}$, was 2.5×10^{-2} cm $\sim 7 r_{De}$. Here $\omega_{s\max}$ is the frequency corresponding to the maximum of the spectrum. This value corresponds satisfactorily to $l_{\min} \sim 15 r_{De}$ obtained in a numerical simulation.[3] We note that since the caviton, owing to the generation of ions, continues to deepen for still some time in the burn-out stage of the Langmuir oscillations, the value of $l_{\min}$ determined from the ion-sound spectrum must, of course, be less than $l_{\min}$ corresponding to the final collapse and the start of the burn-out. We also note that the position of the maximum of the frequency peak is independent of how much the beam electron density is above threshold, i.e., of the turbulence level.

What we have said about the relation between the caviton size and the frequency of the excited ion-sound oscillations is qualitative in nature. However, the fact that the size and the properties of the cavitons, determined this way, agree with the results of numerical calculations, for instance, the fact that the final size of the cavitons is independent both of the turbulence level and of the ion mass, confirms the existing ideas about the strong turbulence picture.

We have shown experimentally that the threshold for the transition to the strong-turbulence regime, measured by the appearance of radiation bursts, corresponds also to a restructuring of the ion-sound oscillation spectrum. We now discuss what determines the transition threshold.

If the energy density W of Langmuir turbulence excited by an electron beam is determined by quasilinear effects, the following relation (see Ref. 7) is valid for W:

$$\frac{W}{nT_e} \sim \frac{n_b}{n_0}\left(\frac{v_b}{v_{Te}}\right)^4. \tag{3.2}$$

In the regime of weak turbulence, caused by decay processes of a plasmon into a plasmon and a phonon, the parameter W/nT_e can be estimated from the condition of equality of the growth rates of the nonlinear wave interaction

$$\gamma_{nl} \sim \frac{W}{nT_e}\frac{m}{M}(kr_{De})^2\omega_{pe}$$

and the beam instability

$$\gamma_b \sim \omega_b \frac{n_b}{n_0}.$$

Using the fact that $k \sim \omega_{pe}/v_b$ we get

$$\frac{W}{nT_e} \sim \frac{M}{m}\left(\frac{v_{Te}}{v_b}\right)^2\frac{n_b}{n_0}. \tag{3.3}$$

It follows from (3.2) and (3.3) that the turbulence level is determined by quasilinear effects if

$$\left(\frac{v_b}{v_{Te}}\right)^6 < \frac{M}{m}, \tag{3.4}$$

which is satisfied in the present experiment.

However, it is well known that for turbulence levels

$$\frac{W}{nT_e} > (kr_{De})^2 \tag{3.5}$$

modulational instability and Langmuir wave collapse must occur. We then get for the transition to a regime of strong turbulence according to (3.5), if we use (3.4),

$$\frac{n_b}{n_0} > \left(\frac{v_{Te}}{v_b}\right)^6 > \frac{m}{M}. \tag{3.6}$$

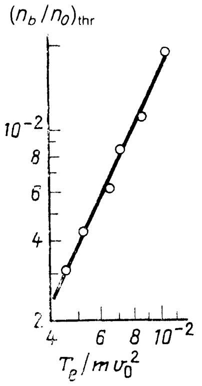

FIG. 4. The threshold value of the electron beam density parameter $(n_b/n_0)_{thr}$ as function of the energy parameter T_e/mv_b^2 for the transition to strong Langmuir turbulence.

Indeed, unstable density n_b obtained on the left-hand side of (3.6) is an overestimate. If the oscillation spectrum is narrow with a width $\Delta k \ll k$, the instability threshold is lowered to

$$\frac{W_{thr}}{nT_e} \sim (\Delta k r_{De})^2 \tag{3.7}$$

and we then have from (3.7) for the threshold beam density

$$\frac{n_{b\,thr}}{n_0} \sim \left(\frac{v_{Te}}{v_b}\right)^4 (\Delta k r_{De})^2. \tag{3.8}$$

In the quasilinear regime the spectrum of the excited oscillations is suffiently narrow, and the threshold dependence found in the experiment

$$\frac{n_{b\,thr}}{n_0} \propto \left(\frac{v_{Te}}{v_b}\right)^4$$

(see Fig. 4) is thus a natural one. At the same time this means that there is no weak-turbulence regime in the framework of the present experiment. The fact that the threshold for the transition to a strong-turbulence regime is independent of the ion mass also gives indications in favor of strong turbulence.

4. EVOLUTION OF THE ELECTRON DISTRIBUTION FUNCTION

The relaxation of the electron beam depends significantly on the nature of the turbulence. In the quasilinear regime the beam relaxes as the result of retardation by the plasma oscillations excited by it to a "plateau" state, but the main plasma electron distribution function is then unchanged. In the weak-turbulence regime nonlinear effects limit the turbulence level and the rate at which the plateau is established is slowed down. Moreover, the transfer of the oscillation energy to the long-wavelength part of the spectrum where the phase velocity of the waves exceeds the beam velocity leads to a partial acceleration of the beam electrons, which by and large changes little the qualitative relaxation picture.

If, however, plasma collapse plays an important role in

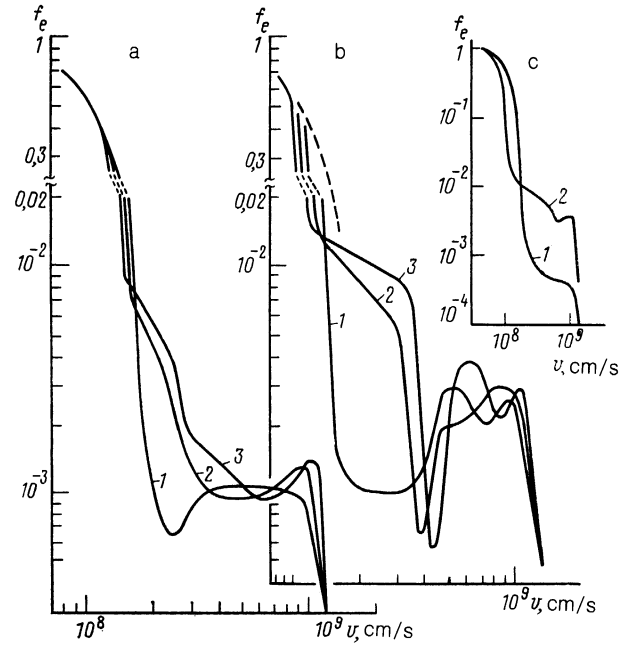

FIG. 5. Dynamics of the distribution function of the plasma and the beam electrons when the beam density parameter is three times (a) and six times (b) the threshold value $(n_b/n_0)_{\rm thr}$. The measurement were carried out at the following times: 0.2 (1); 0.5 (2); and 1 μs (3). (c) comparison of the stationary distribution functions measured 1 μs after the beam was switched on for a value of the beam density parameter close to its threshold value $\sim (n_b/n_0)_{\rm thr}$ (1) and when it is ten times that value (2).

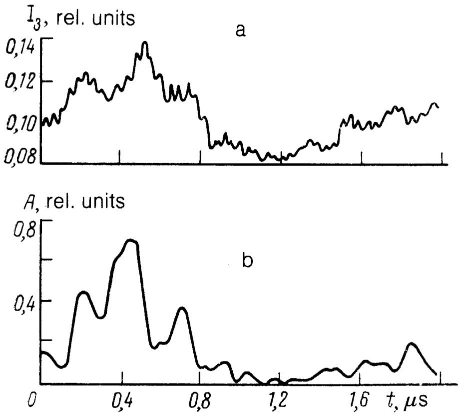

FIG. 6. Oscillograms of the accelerated beam electron current to a probe after having been retarded through a potential $U_a \gtrsim U_b$ (a) and of plasma field "bursts" near the input opening of the probe (b).

the turbulence there occurs at the same time as the beam relaxation an acceleration of the electrons in the plasma itself, a drawing out of "tails" of the main electron distribution function since only the collapse process leads to a transfer of energy to the short-wavelength part of the spectrum in the region of small phase velocities of waves which are damped by the plasma electrons.

We show in Fig. 5 the temporal evolution of the distribution function of the electrons in the plasma and beam. Under the conditions of Fig. 5b (beam current 6 A or $n_b/n_{b\,\rm thr} \approx 6$) there is a sharp separation between the distribution functions of the electrons in the plasma and in the beam, owing to the existing "gap" occurring between them, at an electron energy ≈ 30 eV. The beam relaxation is rapid and it makes it impossible to resolve it in time, whereas the opposed process of a growth of tails on the main plasma distribution function is a slow process which sets in only during the first microsecond, which in terms of ion oscillation periods corresponds to $80\omega_{pi}^{-1}$ and which as to order of magnitude coincides with the characteristic duration of the caviton collapse.[4]

When we decrease or increase the beam current (or, correspondingly, the parameter $n_b/n_{b\,\rm thr}$) (cf. Figs. 5a, c) the gap between the accelerated plasma electrons and the beam electrons is less pronounced and as time goes on ($\tau > 0.5\,\mu$s) tends to vanish, but in the small velocity region the nature of the distribution function remains different from a plateau. The beam relaxation is essentially different in the distributions given by curves 1 and 2 of Fig. 5c. Whereas in the first case for a value of the current close to threshold $I_b \lesssim 1$ A ($n_b/n_{b\,\rm thr} \lesssim 1$) there is a typical plateau for quasilinear relaxation, in the second case when $I_b \approx 10$ A ($n_b/n_{b\,\rm thr} \approx 10$) the place of the plateau is occupied by the "non-exponential" tails of the plasma electron distribution function.

The acceleration of the electrons in the plasma itself, reflected in the formation of tails in the main electron distribution function is a direct confirmation of the existence in the plasma of strong Langmuir turbulence processes leading to the transfer of plasma waves to the small-scale region. This is also indicated by the slow evolution of the tails of the plasma electron distribution corresponding to the ion dynamics of collapsing cavitons and also by the fact that the maximum increase in the energy of the accelerated plasma electrons $\Delta\mathscr{E}_{e\max}$ is close to the estimate of that quantity for the final stage of the collapse:

$$\Delta\mathscr{E}_{e\max} = eE_{\max} l_{\min} = 2^{1/2}(W_{\max}/n_0T_e)^{1/2} T_e l_{\min} r_{De}^{-1} \sim 30T_e. \tag{4.1}$$

Here $W_{\max} = E^2_{\max}/8\pi$ and $l_{\min}$ are, respectively, the energy density of the Langmuir field in the caviton and its characteristic size in the final stage of the collapse, which according to the data of the numerical experiments are equal to $W_{\max}/n_0T_e \approx 2$ and $l_{\min} \approx 15 r_{De}$.

The beam electrons also exchange energy with the caviton field. Applying to the multigrid probe a negative potential $U_a > U_b$ which retards the bulk of the beam electrons we can split off the current of the beam electrons which obtained additional energy in the plasma. It turned out that the current of the additionally accelerated beam electrons undergoes a modulation with time which is illustrated by the oscillogram of Fig. 6a. At the same time we registered, by means of a loop antenna encircling the input opening of the probe, field-oscillations bursts which we identified above with the sources of the electromagnetic radiation of the plasma (Fig. 6b). The intercorrelation of these signals makes it possible to assume that the beam electrons are accelerated in the cavitons which are formed in the immediate vicinity of the input opening of the probe. The increase in the beam electron energy must in this case be maximal also in the final stage of the collapse. The quantity $\Delta U_{b\max}$ must agree in order of magnitude with the estimate (4.1).

CONCLUSION

In contrast to well known Refs. 4 and 12, we have studied not separate localized strongly nonlinear structures, solitons, or collapsing cavitons but an ensemble of such structures which characterize strong plasma turbulence. Using various methods we have, in our opinion, amply confirmed the picture of turbulence as an ensemble of collapsing Langmuir cavitons. We were in this case able not only to confirm the qualitative picture of the turbulence but also to determine the minimum size of the cavitons, the length of the concluding stage of their collapse, and to estimate the energy density of the Langmuir field in the cavitons. These results are well correlated with data from a numerical experiment.

On the basis of the data presented here we can give a quantitative description of turbulence in terms of the average energy density of the Langmuir field in the plasma. Knowing the average number N of radiation bursts per unit time, the time τ of their burn-out, and also the volume V from which we know the method of radiation, we can estimate the caviton density in the final stage of collapse in the form

$$\eta_{\rm fin} = N\tau/V.$$

On the other hand, using the data of Ref. 9 for the final volume of the caviton, $v_{\rm fin} = l_{\parallel\min} l^2_{\perp\min} \approx 9 l^3_{\parallel\min} \approx 2.5\times 10^4 r^3_{De}$, where $l_\parallel$ and $l_\perp$ are, respectively, the longitudinal and transverse dimensions of the caviton, we can estimate the fraction of the volume occupied by cavitons in the final stage of their collapse with a Langmuir field energy density $W_{\max}/n_0 T_e$ in them of order 1:

$$\alpha = \eta_{\rm fin} v_{\rm fin}.$$

The quantity α is the coefficient determining the average value of the Langmuir field energy in the plasma: $W_{\rm av} = \alpha W_{\max}$. In actual fact, the Langmuir field energy concentrated in the cavitons at different stages of their evolution is about four times larger, since the total caviton density is correspondingly larger than the density of the radiating cavitons owing to the ratio of the time of caviton formation up to the concluding stage $\Delta t = 0.4\,\mu$s (measured by the delay in the appearance of bursts relative to the front where the beam starts) and the length of the concluding stage itself $\tau = 0.1\,\mu$s. For instance, for a beam with energy $U_b = 300$ eV and density parameter $n_b/n_0 = 7(n_b/n_0)_{\rm thr}$ the number of bursts per unit time is $N = 2\times 10^6\,{\rm s}^{-1}$ and the caviton density in the final stage is $\eta_{\rm fin} = 20\ {\rm cm}^{-3}$ which corresponds to an average Langmuir energy density $W_{\rm av}/n_0 T_e \approx 5\times 10^{-2}$.

[1] V. E. Zakharov, in *Basic Plasma Physics* (eds A. A. Galeev and R. N. Sudan), North Holland, Amsterdam (1984), Vol. II, p. 81.
[2] V. D. Shapiro and V. I. Shevchenko, in *Basic Plasma Physics* (eds. A. A. Galeev and R. N. Sudan), North Holland, Amsterdam (1984), Vol. II, p. 123.
[3] V. E. Zakharov, A. N. Pushkarev, A. M. Rubenchik *et al.*, Pis'ma Zh. Eksp. Teor. Fiz. **47**, 239 (1988) [JETP Lett. **47**, 287 (1988)].
[4] A. Y. Wong and P. Y. Cheung, Phys. Rev. Lett. **52**, 1222 (1984); P. Y. Cheung and A. Y. Wong, Phys. Fluids **28**, 1538 (1985).
[5] D. Dubois, H. Rose, and D. Russell, Phys. Rev. Lett. **61**, 2209 (1988).
[6] N. N. Astrakhantsev, O. L. Volkov, Yu. Karavaev, and G. N. Kichigin, Phys. Lett. **A110**, 129 (1985).
[7] A. A. Galeev, V. D. Shapiro, V. I. Shevchenko, and R. Z. Sagdeev, Zh. Eksp. Teor. Fiz. **72**, 507 (1977) [Sov. Phys. JETP **45**, 266 (1977)].
[8] E. D. Volkov, N. F. Perepelkin, V. A. Suprunenko, and E. A. Sukhomlin, *Collective Effects in a Current-Carrying Plasma*, Naukova Dumka, Kiev (1979).
[9] P. Y. Cheung, Thesis, UCLA (1984).
[10] J. Dawson, in *Advances in Plasma Physics*, Interscience, New York (1968), Vol. 1, p. 1.
[11] V. Yan'kov, Fiz. Plazmy **8**, 86 (1982) [Sov. J. Plasma Phys. **8**, 48 (1982)].
[12] M. V. Nezlin, *Beam Dynamics in a Plasma*, Energoatomizdat, Moscow (1982), p. 190.

Translated by D.ter Haar