Anomalous electron heating in the field of oblique Langmuir waves in induced *ls* scattering

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The anomalous heating of electrons by high-frequency field near the hybrid resonance in a plasma with ion-sound instability is considered. The analysis takes into account the increase of the high-frequency wave energy due to scattering of the pumping wave, as well as the effect of variation of the plasma potential on the electron energy losses along the magnetic field. It is shown that the electron heating can be explained by introducing an effective collision frequency that is close to the plasma frequency and exceeds the pair-collision frequency by an order of magnitude. The thermal and electric conductivities of the plasma are found to decrease along the magnetic field. A correlation is observed between the electron heating and the high-frequency field excitation of ion-sound waves. It is shown that the excitation of ion-sound waves can be attributed to induced *ls* scattering of waves in a current-carrying plasma. By comparing the measurement results with those of the theory of induced *ls* scattering it is concluded that the weak turbulence theory can explain only the energy gained by the electrons in a high-frequency field, but not the observed loss of the momentum component along the magnetic field.

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1. The mechanisms that heat the electrons when parametric instability of RF waves is produced in a plasma near the lower hybrid resonance remain unclear to this day. Since the phase velocities of the Langmuir waves excited by the RF field exceed substantially the thermal velocities of the electrons, the appearance of "tails" of fast electrons is natural. This result is qualitatively confirmed in a number of experiments.[1–4] At the same time, heating of the entire mass of the electron gas was observed in Refs. 5–9. It is usually assumed that the observed effect cannot be attributed to wave dissipation by collisions. Such statements, however, call for additional corroboration, since the analysis of the measurement results has so far not included discussions of essential factors that influence the RF heating in particle collisions. These include the decrease of the energy lost by the electrons because of the increase of the plasma potential relative to the walls, the lower RF-field energy density in the plasma compared with the energy density of the pump wave on account of its parametric instability, the existence in the pump-wave spectrum of spatial modes of oscillations with phase velocity close to the electron thermal velocity, and other. All these circumstances call for performance of experiments in which these factors are monitored. On the other hand, it was pointed out in Ref. 10 that electron heating near the lower hybrid resonance can be due to induced scattering of a Langmuir wave by electrons with transformation of the scattered wave into an ion-sound one (*ls* scattering). Such a process can be substantially more effective in a magnetoactive plasma than in a plasma without a magnetic field. Indeed, the condition that phase velocity of the beats of the Langmuir and sound waves be in synchronism with the thermal electron takes the form where ω_0 and ω_s are the cyclic frequency of the pump wave and of the ion-sound wave, $K_{0\parallel}$ and $k_{s\parallel}$ are the components of the pump- and sound-wave vectors along the magnetic field, and $v_{Te} = (T_e/m)^{1/2}$ is the thermal velocity of the electrons. In the case of lower-hybrid and Langmuir waves in a magnetoactive plasma the pump frequency is much lower than the Langmuir electron frequency ω_{pe}, therefore

$$k_{s\parallel} \approx \omega_0/v_{Te} \ll \omega_{pe}/v_{Te} = r_{De}^{-1},$$

where $k_{s\parallel} > k_{0\parallel}$. This means that, in contrast to a nonmagnetized wave, long-wave ion-sound noise participates in the stimulated scattering, and not solely oscillations with $kr_{De} \gtrsim 1$. The result should be an increase in the energy density of the acoustic noise that participates in the heating, and consequently an increased heating efficiency.

All these considerations prompted us to search for a correlation between the electron-heating conditions and the characteristics of the ion-sound noise excited by the Langmuir waves. The heating and the oscillation excitation can be compared in terms of several parameters. First, it must be verified that the heating is simultaneously accompanied by excitation of oscillations with wave vectors defined by the condition (1). Second, the corresponding calculations (see the Appendix) provide estimates of the effective frequencies of the variation of the energy and longitudinal momentum of the electrons at a given energy density of the ion-sound noise and of the pump waves. This permits the change of the electron temperature to be set in correspondence with the calculated quantities. Third, it is known[11,12] that the onset of sufficiently effective *ls* scattering calls for the presence of asymmetry of the distribution function, for example for the existence of directional electron velocity. A control experiment is therefore necessary. In addition, as already noted, when the measurements are made it is necessary to monitor the conditions that contribute to electron heating via collisional dissipation of the RF waves.

2. The experiments were performed on an argon plasma

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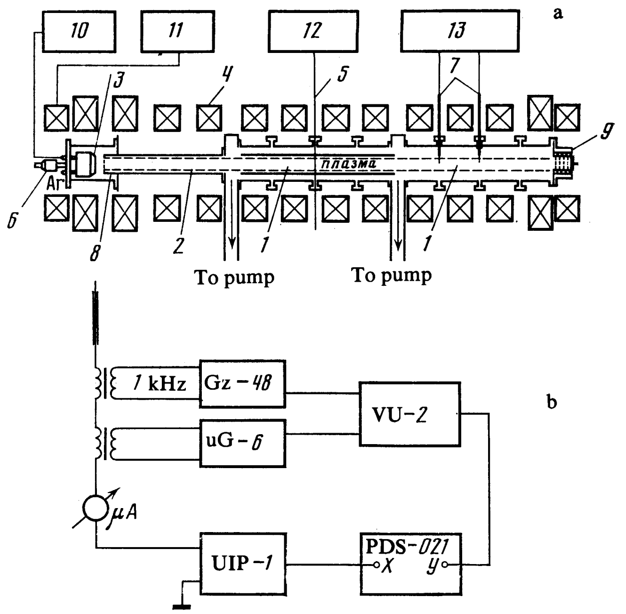

FIG. 1. a.—Experimental setup: 1—working chamber, 2—differential evacuation section, 3—cathode unit, 4—magnetic-field coil, 5—plates for excitation of the RF field, 6—gas leak valve, 7—probes, 8—grid anode, 9—multigrid probe, 10—plasma supply source, 11—magnetic-field supply, 12—RF oscillator, 13—diagnostic apparatus; b—circuit for electric differentiation of probe characteristic.

column produced in a uniform magnetic field by a low-energy beam of electrons emitted from a hot cathode (Fig. 1a). The cathode region was separated from the region where the plasma interacts with the RF field by a drift tube that ensured a 20-fold pressure drop between these regions. The cathode potential was varied in a range from − 60 to − 150 V. The longitudinal magnetic field strength was 600–800 Oe. The argon pressure in the interaction region was $(2-6)\times10^{-4}$ Torr. The plasma density was $(1-2)\times10^{10}$ cm^{-3}. The initial electron temperature was varied in the range from 3 to 6 eV. The plasma density was uniform along the column over the entire length of the interaction region (100 cm) and had a bell-shaped radial distribution with a characteristic inhomogeneity scale 2 cm.

The pump wave was an axisymmetric oblique Langmuir wave of frequency 24 MHz (Trivelpiece-Gould mode). The interaction region was spanned by half the oscillation wavelength. The frequency relation was $\Omega_i \ll \omega_s \ll \omega_0 \approx 5\omega_{\rm lh} \ll \omega_{pe} \lesssim \Omega_e$ (Ω_i, and Ω_e are the gyrofrequencies of the ions and electrons, and $\omega_{\rm lh}$ is the lower hybrid frequency).

The plasma density was determined from the saturation ion current of the Langmuir probes, while the ion-sound waves were determined from the ac component of the ion saturation current. The high-frequency field near the pump-wave frequency was measured with a double Langmuir probe loaded by an RF transformer. The probe was calibrated beforehand in a capacitor field.

The wave numbers of the ion-sound and RF oscillations near the pump frequency were determined by measuring the spatial mutual correlation function of the signals from two probes that were moved relative to each other, with the signal produced at the specified frequency.[13]

The electron distribution function was measured with single-electrode flat probes introduced into the chamber through lateral stubs and displaced along the chamber radius. The probes were constructed of two electrodes measuring 2×3 mm, separated by a ceramic insulator, so that currents could be measured simultaneously in two opposite directions. The electron distribution function in longitudinal velocity was obtained by electric differentiation of the currents as the retarding potential was varied. For electric differentiation, an alternating voltage of 0.1 V amplitude and 1 kHz frequency was applied to the probe electrode (Fig. 1b). Synchronous detection was used to increase the circuit sensitivity.

It is known that a magnetic field influences substantially the electron current of a flat probe if the radius r_0 of the latter exceeds the electron gyroradius ρ_e (Refs. 14, 15). As a result, the determination of the electron distribution function by differentiating the dependence of the probe current on the retarding potential (of the current-voltage characteristic) calls for special elucidation. It is known that the variation of the electron current to a flat probe, with allowance for the influence of the magnetic field, is given by the expression[14,15]

$$i_{\uparrow\uparrow}=env_{Te}S_0\{{}^1/_4E_1+{}^1/_2\varkappa^{1/2}[1-\Phi(x_1)]\}\times\frac{8r_0(D_{e\perp}D_{e\parallel})^{1/2}}{\{{}^1/_4E_1+{}^1/_2\varkappa^{1/2}[1-\Phi(x_1)]\}S_0v_{Te}+8r_0(D_{e\perp}D_{e\parallel})^{1/2}} \quad (2a)$$

$$i_{\uparrow\downarrow}=env_{Te}S_0\{{}^1/_4E_2-{}^1/_2\varkappa^{1/2}[1-\Phi(x_2)]\}\cdot\times\frac{8r_0(D_{e\perp}D_{e\parallel})^{1/2}}{\{{}^1/_4E_2-{}^1/_2\varkappa^{1/2}[1-\Phi(x_2)]\}S_0v_{Te}+8r_0(D_{e\perp}D_{e\parallel})^{1/2}}\cdot \quad (2b)$$

Here $i_{\uparrow\uparrow}$ is the value of the current of a probe turned towards the cathode (0°), $i_{\uparrow\downarrow}$ is the current of a probe turned towards the end face (180°), $\Phi(x)$ is the probability integral,

$$x_{1,2}=(e\varphi/T_e)^{1/2}\mp(\mathscr{E}_0/T_e)^{1/2},$$

φ is the retarding potential of the probe relative to the plasma, $\mathscr{E}_0 = mv_0^2/2$ is the energy of the directional motion of the electrons, $E_{1,2} = \exp(-x_{1,2}^2)$, $S_0 = \pi r_0^2$, and $\varkappa = \mathscr{E}_0/T_e$.

It is easily seen from (2) that a magnetic field can strongly distort the current-voltage characteristics at low retarding potentials by decreasing the electron current. With increasing retarding potential, however, this influence weakens. Measurements (Fig. 2) show that at $\varphi = 0$ the probe currents are 1.5–2 times smaller than the values given by Eq. (2), if the quantities $D_{e\perp}$ and $D_{e\parallel}$ are taken to mean the classical coefficients of diffusion across and along the magnetic field ($D_{e\perp} = \rho_e{}^2\nu_{en}$, $D_{e\parallel} = v_{Te}{}^2/\nu_{en}$, ν_{en} is the electron-neutral collision frequency). Thus, the model considered in Refs. 14 and 15 describes correctly the influence of the magnetic field on the probe current. This is confirmed also by measurements of the ratio of the electron and ion currents when the probe is oriented towards the end face (180°). For an ion Larmor radius $\rho_i > r_0$, Ref. 14 gives the relation

$$\frac{i_e}{i_i}\approx4\left(\frac{M}{m}\right)^{1/2}\frac{\rho_e}{r_0}. \quad (3)$$

The difference between the measured and calculated current ratios is not more than 50%.

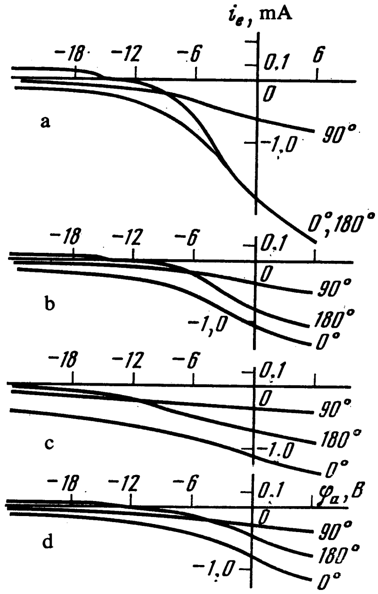

FIG. 2. Current-voltage characteristics of plane one-sided probes at three orientations relative to the electron beam: a, c—distance of probe from end 80 cm, b, d—20 cm; a, b—$E_0 = 0$; c, d—$E_0 = 9$ V/cm (φ_a is the potential of the probe relative to ground).

From an analysis of the curves of Fig. 2a it can also be seen that even at $|e\varphi|/T_e \approx 2.3$ the probe-current decrease due to the influence of the magnetic field is not more than 25%. Thus, the first derivative of the dependence of the probe current on the retarding voltage is indeed the electron longitudinal-velocity distribution function at retarding-potential values exceeding by 2–2.5 times the electron temperature. This is clearly seen from the curves of Fig. 3: a deviation of the distribution function from exponential is observed only at low energies, not higher than 7–12 eV.

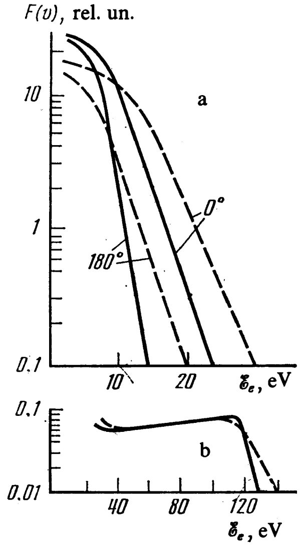

FIG. 3. Change of electron longitudinal velocity distribution function upon action of an RF field at $n = 10^{10}$ cm^{-3}, $p = 3\times10^{-4}$ Torr and $H = 600$ Oe. Solid curve—$E_0 = 0$, dashed—$E_0 = 9$ V/cm; a—low-energy part, b—high energy part.

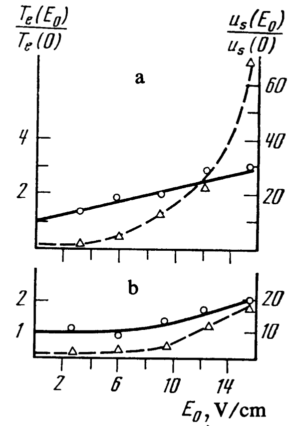

FIG. 4. Relative change of electron temperature and ion-sound-oscillation intensity for two source regimes: a) $T_e = 4.2$ eV, b) $T_e = 6.5$ eV; the level of the ion-sound noise is 6 times lower in regime a than in regime b.

We note now a number of features of the current-voltage curves and of the distribution function. When the probe faces the cathode (0°) the current-voltage characteristics show clearly the current of the beam of electrons that ionize the gas. From a comparison of the distribution functions measured along and counter to the direction of the ionizing beam (Fig. 3) it can be seen that they are asymmetric, meaning the presence of an electron drift velocity, with $v_0 = 0.1\, v_{Te}$.

Application of the RF field deforms the current-voltage curves. First, the saturation electron current is decreased. Second, a difference arises between the saturation currents measured along and counter to the ionizing-beam direction. This difference increases with increasing RF field. These features of the behavior of the current in an RF field will be discussed later on.

3. The dependences of the electron temperature on the RF field were measured at various initial temperatures and energy densities of the ion-sound noise (Fig. 4). The initial electron temperature and ion-sound noise density decreased with increasing argon pressure and with decreasing drift current. The measurements showed (Fig. 4) that when the RF field was increased to 16 V/cm the electron temperature at the center of the plasma column increased to 12–16 eV, with simultaneous increase of the ion-sound noise energy density to $(2–4)\cdot10^{-4} nT_e$, and was practically independent of the initial values of the temperature and noise energy density. The relative changes of these quantities were larger the lower their initial values.

Simultaneously with measurements of the temperature and intensity of the ion-sound fluctuations, we measured the distribution of the plasma potential over the plasma-column cross section (Fig. 5). It can be seen that the potential did not change substantially over the cross section of the plasma column.

It is easily seen that the recorded rise of the electron temperature cannot be attributed to the overheating of the electrons in the RF field by the collisions. Indeed, the overheating due to collisions can be estimated by recognizing that the electrons are cooled by departure to the end face of

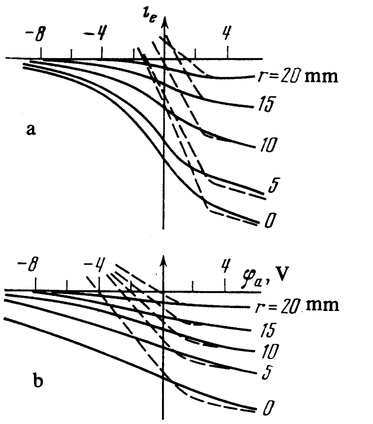

FIG. 5. Characteristics of "cold" (solid lines) and "hot" (dashed) probe for different probe positions on the radius at $E_0 = 0$ (a) and $E_0 = 9$ eV (b).

the chamber along the magnetic field. Assuming that the heat transport is determined by the longitudinal heat conductivity of the electrons and that the energy release is uniform along the column, solution of the heat conduction equation yields the following estimate of the temperature T_{e0} near the end of the electrode and for the temperature drop ΔT_e along the column:

$$T_{e0}=\left[\frac{(2\pi)^{1/2}\exp(e\varphi_0/T_{e0})}{3+2e\varphi_0/T_{e0}}\,m^{1/2}L\left(\frac{W_l}{n}+mv_0^2\right)\nu_{en}\right]^{2/3}, \qquad (4a)$$

$$\Delta T_e=\left[\,{}^2/{}_5\,\nu_{en}^2L^2m\left(\frac{W_l}{n}+mv_0^2\right)\right]^{1/2}. \qquad (4b)$$

Here φ_0 is the plasma potential near the end-face electrode. In the solution of the heat-conduction equation it was assumed that the heat flux to the end face is equal to the flux of the energy released in the plasma volume by dissipation of RF waves with energy density W and by Joule heat of a current with drift velocity v_0. Comparison of the RF-wave energy density at $E_0 = 15$ V/cm with the energy density of the Joule losses shows that the former amounts to only 0.2 of the latter. Even this circumstance alone indicates that the electron heating has a very complicated dependence on the pump field. Estimates by Eqs. (4) at $e\varphi_0/T_{e0} \approx 1$ and $v_0 = 0.1$ $v_{Te} = 10^7$ cm/sec yield values $T_{e0} = 0.81$ eV, $\Delta T_e = 1.23$ eV and $T_{e\,max} = T_{e0} + \Delta T_e \approx 2$ eV, close to the initial electron temperature and lower by 6–8 times that the value measured at $E_0 = 16$ V/cm. We now take into account the fact that the RF-field energy density in the plasma is five times larger than the energy of pump wave, owing to the waves scattered by the ion-sound oscillations due to the drift dissipative instability of the plasma.[16] This increases T_{e0} to 1.13 eV and ΔT_e to 1.6 eV, which is again one-fifth the measured value. Agreement between the calculated quantities and the measured ones can be obtained by introducing an effective collision frequency $\nu_{\rm eff}$ that exceeds by 5–10 times the frequencies of the pair collisions. This assumption leads to two consequences. First, an appreciable temperature drop should occur along the column. Second, the increase of the collision field should cause a decrease in the voltage drop on the plasma column, owing to the flow of longitudinal current. It turns out that these two consequences indeed take place, as shown by measurements of the electron temperature and the plasma potential, in two sections of the column at distances $z = 20$ and 70 cm from the end face (see Table I).

An estimate of the effective collision frequency, obtained from the voltage drop along the plasma column, shows that at $E_0 = 12$ V/cm the value of $\nu_{\rm eff}$ is indeed higher by a decade than the frequency of electron-neutron collisions.[1)] Additional evidence in favor of the increase of $\nu_{\rm eff}$ in an RF field is the decrease of the saturation electron current of a flat Langmuir probe (see Fig. 2). Indeed, this phenomenon can be attributed both qualitatively and quantitatively to a decrease in the value of $D_{\rm eff} = v_{Te}^2/\nu_{\rm eff}$ in accord with (2). It must be assumed here that the collision frequency increases only for electrons moving along the magnetic field. Agreement between the measured decrease of the current and the calculated value is obtained if it is recognized that $\nu_{\rm eff}$ increases by an order of magnitude, and if account is taken of the measured increase of the electron temperature and of the decrease of the average plasma density.

In the upshot, all the measurements of the electron heating, of the voltage drop along the plasma column, and of the decrease of the saturation electron current of flat probes leads us to the conclusion that $\nu_{\rm eff}$ increases in an RF field to values of the order of the ion plasma frequency ω_{pi}.

4. One can attempt to attribute the effective collision frequency to the process of induced scattering of Langmuir waves by electrons. An increase of the intensity of this process calls for asymmetry of the electron-longitudinal-velocity distribution function (see Eq. (A.1) in the Appendix), and this is satisfied in the experiment because of the presence of the longitudinal current. The induced ls scattering should excite ion-sound oscillations with longitudinal wave vectors $k_{s\parallel} \geqslant (\omega_0/\omega_{pe})r_{De}^{-1}$. Measurements of the spatial-correlation function for different frequencies of the ion-sound frequencies reveal a substantial variation in their longitudinal-wave-number spectrum in an RF field (Fig. 6). It was found that

TABLE I. Change of electron temperature and of plasma potential with changing pump field strength in two sections of the plasma column.

z, cm	E_0, V/cm: 0		3		6		9		12		16	
	T_e, eV	φ, V	T_e, eV	φ, V	T_e, eV	φ, V	T_e, eV	φ, V	T_e, eV	φ, V	T_e, eV	φ, V
20	2.6	4.8	3.2	6.8	3.5	7.7	4.8	8	6.4	10	7.3	12.3
70	4.6	4.2	5.7	4.2	6.6	3.8	7.8	3.3	10.2	2.6	12.1	2.4

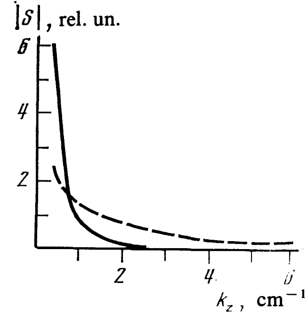

FIG. 6. Spectral energy density vs the wave-vector component along the magnetic field for a LF field at frequency 1 MHz: solid line—$E_0 = 0$, dashed—$E_0 = 10$ V/cm.

the ion-sound oscillations excited by the drift current have in the initial state a substantial anisotropy of the transverse and longitudinal wave vectors.[17] The action of the pump field increases the intensity mainly of the oscillations having large longitudinal wave-vector components. Indeed, as follows from Fig. 6, whereas in the initial plasma at a frequency 1 MHz more than 90% of the energy goes into waves with

$$k_{s\parallel} < (\omega_0/\omega_{pe}) r_{De}^{-1} \approx 1\,\text{cm}^{-1},$$

at $E_0 = 10$ kV/cm more than 70% of the energy goes to waves with $k_{e\parallel} > 1$ cm^{-1}. Thus, the increase of the energy density of the ion-sound noise can indeed be attributed to induced *ls* scattering.

As a control experiment we measured the intensity of ion-sound noise and of the electron temperature following pulsed turning-on of the cathode voltage, which decreased the electron-current drift to zero. From the oscillograms of Fig. 7 it can be seen that if the cathode voltage is turned off completely and the drift current decreases to zero, the ion-sound noise is damped out 5–7 μsec at the maximum pump field (Fig. 7a). The damping of the sound oscillation is not

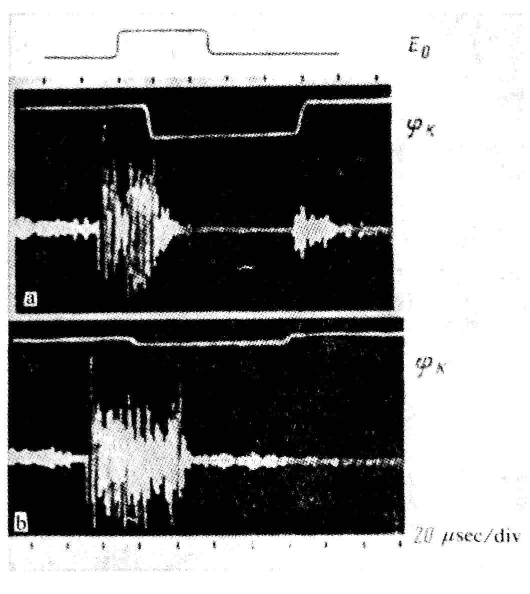

FIG. 7. Oscillograms of ion-sound noise following pulsed turning-on of the RF field at $E_0 = 15$ V/cm and pulsed turning-off of the cathode voltage: a—complete turnoff, b—incomplete.

due to a decrease of the plasma density by turning off the cathode voltage, since the plasma decay has an approximate time constant 20 μsec. At the same time if the cathode voltage is not completely turned off (e.g., it is decreased to $\varphi_c = -80$ V (Fig. 7b), when the drift current still flows but the plasma density does decrease), the same enhancement of the ion-sound noise is observed as prior to turning off the cathode voltage. The excitation of the ion-sound oscillations correlates thus with the presence of the electron drift current.

We compare now the measured values of ν_{eff} with their estimates obtained in the theory of weak turbulence of the induced *ls* scattering (see the Appendix). Measurements of the energy density of the sound noise with Langmuir probes at the total saturation current yield values $(n\sim/n)^2 = u_s \approx 10^{-5}$ without RF and $(4\text{–}10)\times 10^{-4}$ at $E_0 \approx 16$ V/cm. Therefore, according to Eqs. (A.5) and (A.6) of the Appendix we can obtain the estimates $v^{p}_{\text{eff}} = (1-3)10^5$ sec^{-1} and $v^{\mathscr{E}}_{\text{eff}} = (1.6-5)10^7$ sec^{-1}. In expression (4a) for the electron temperature T_{e0} of the end-face electrode, ν_{en} stands in fact for the quantity $\nu^{\mathscr{E}}_{\text{eff}}$ determined from rate at which the electrons acquire energy from the RF field. Substituting in (4a) the value $v^{\mathscr{E}}_{\text{eff}} = (1.6-5)10^7$ sec^{-1} calculated above, we obtain $4.9 \leqslant T_{e0} \leqslant 9.6$ eV which is in satisfactory agreement with the measured $T_{e0} = 7.3$ eV. The decrease of the longitudinal thermal and electric conductivities, as well as of the electron diffusion coefficient, which manifests itself in the temperature and potential drops and in a decrease of the saturation electron current in the probes, cannot be attributed to the increase in the electron scattering frequency in the *ls* process, since the value of ν^{p}_{eff} is lower than the required one by almost two orders. One of the sources of the discrepancy between theory and experiment may be the underestimate of the energy density of the short-wave ion-sound noise with $k_s r_{De} \gtrsim 1$, measured with probes larger than the oscillation wavelength. In addition, the cause of the observed discrepancy may be that the weak-turbulence theory is not applicable at the large experimental values of u_s.

5. The analysis of the electron heating in the field of oblique Langmuir waves convinces us that the phenomenon considered is essentially nonlinear and cannot be attributed to electron heating by collisions, even if it becomes possible to take into account the effect of the change of the plasma potential on the heating condition and of the general increase of the Langmuir-wave energy density, compared with the pump wave energy as a result of scattering and nonlinear excitation. It turns out that to explain the heating it is necessary to introduce an effective collision frequency of the order of the ion plasma frequency. That the effective collision frequency exceeds the pair-collision frequency is confirmed also by measurements of the decrease of the electric conductivity and of the electron diffusion along the magnetic field with increasing RF field.

Measurements show also that when the pump-wave field is increased from 3 to 16 kV/cm the energy of the ion-sound noise is increased by two orders. It is found here that the longitudinal wave vectors of the excited noise are typical of the induces *ls*-scattering process, and the noise is excited only in the presence of an electron drift current, as called for by the *ls*-scattering theory. The measurement results permit a correlation to be established between the heating process and the excitation of ion-sound noise. Thus, both phenomena can be attributed to the occurrence of induced *ls* scattering.

Comparison of the measured values of $\nu_{\rm eff}^{p}$ with those calculated in weak-turbulence theory shows that they differ in order of magnitude. It appears that in this case we are dealing with a strong plasma turbulence.

Since the role of the *ls* scattering in the interaction of RF waves with a nonisothermal plasma should increase when the pump frequency is made lower than the Langmuir frequency, one can expect this scattering to influence strongly the lower-hybrid plasma heating, including the transport processes accompanying this heating. Nor can one exclude a manifestation of induced scattering of microwaves by electrons, with excitation of ion sound, also in the case of electron cyclotron heating, when $|\omega_0 - \Omega_e| \approx k_s v_{Te}$.

The authors take the opportunity to express deep gratitude to the late Professor M. S. Rabinovich for support, constant attention, and interest. The authors thank also L. V. Kolik for developing the measurements.

APPENDIX

Change of electron energy and momentum in induced ls scattering

Let us determine the change of the electron energy and momentum in a magnetized plasma as a result of induced *ls* scattering by electrons. We direct the magnetic field along **z**. We put $\Omega_e \gg \omega_{pe}$ and $\Omega_i \ll \omega_s$. An oblique Langmuir wave with $k_0 \ll k_s$, $k_{0\parallel} \ll k_{s\parallel}$, and frequency $\omega_0 = \omega_{pe}\cos\theta_0 \gg \omega_s$ propagates at an angle θ_0 to the **z** axis ($\cos\theta_0 \ll 1$). From the nonlinear theory of plasma oscillations (see, e.g., Refs. 18 and 19) it follows that at a beat phase velocity

$$v_\parallel = (\omega_0 \pm \omega_s)/(k_{0\parallel} \pm k_{s\parallel}) \approx \pm \omega_0/k_{s\parallel} \leqslant v_{Te}$$

the *ls* scattering leads to the following growth rates of the ion-sound and Langmuir waves:

$$\gamma_N{}^s = \frac{\pi}{4} u_l \omega_s \,\mathrm{sign}(k_{s\parallel}) \frac{\omega_{pe}^2}{\omega_0^2} \frac{(k_s r_{De})^2}{1+(k_s r_{De})^2} \times \left(\cos^2\theta + \sin^2\theta \frac{\omega_{pe}^2}{\Omega_e^2}\sin^2\varphi\right)\left[F'\left(\frac{\omega_0+\omega_s}{k_{s\parallel}}\right) + F'\left(\frac{\omega_s-\omega_0}{k_{s\parallel}}\right)\right], \tag{A.1}$$

$$\gamma_N{}^l = \frac{\pi}{4}\omega_0 \frac{\omega_{pe}^2}{\omega_0^2}\int d^3\mathbf{k}_s W_s(\mathbf{k}_s) \frac{(k_s r_{De})^2}{1+(k_s r_{De})^2} \times \left(\cos^2\theta + \sin^2\theta \frac{\omega_{pe}^2}{\Omega_e^2}\sin^2\varphi\right) \times \left[F'\left(\frac{\omega_0+\omega_s}{k_{s\parallel}}\right) - F'\left(\frac{\omega_s-\omega_0}{k_{s\parallel}}\right)\right], \tag{A.2}$$

where $\cos\theta = k_{s\parallel}/k_s$, φ is the angle between the planes ($\mathbf{k}_0\mathbf{z}$) and ($\mathbf{k}_s\mathbf{z}$), u_l and u_s stand the total energy density and W_s is the spectral energy density, referred to nT_e ($u = \int W(\mathbf{k}) d^3\mathbf{k}$), $F(v_\parallel/v_{Te})$ is the one-dimensional electron-distribution function normalized to n/v_{Te}, and the prime denotes differentiation with respect to the argument.

The energy input to the plasma is determined by the growth rate of the Langmuir wave:

$$\frac{1}{nT_e}\frac{d}{dt} nT_e = \nu_{eff}^{g} u_l = -2\gamma_N{}^l u_l. \tag{A.3}$$

The total loss of electron momentum due to induced *ls* scattering is

$$\frac{d}{dt}P_\parallel = -\nu_{eff}^{p} m v_0 n = -\int 2\gamma_N{}^s \frac{W_s}{\omega_s} k_{s\parallel} d^3\mathbf{k}_s, \tag{A.4}$$

where v_0 denotes the averaged directional velocity of the electrons.

To estimate the possible role of the induced *ls* scattering in the decrease of the longitudinal electric conductivity of the plasma we put

$$F(\xi) = (2\pi)^{-1/2}\exp[-(\xi-\xi_0)^2/2], \quad \xi_0 = v_0/v_{Te} \ll 1.$$

We assume also that the total energy of the ion-sound oscillations is concentrated in the region of the wave vectors $k_{s\parallel} \approx k_{s\parallel}{}^0 = \omega_0/v_{Te}$ corresponding to nonlinear interaction of the waves with the thermal electrons[2)] (we assume here for simplicity that $W_s \sim \delta(k_{s\parallel} - k_{s\parallel}{}^0)$).

For the distribution of the acoustic-oscillation energy in k_s, we use (at $k_s r_{De} < 1$) the approximation

$$du_s/d\omega_s \sim \frac{1}{k_s}\ln\frac{1}{k_s r_{De}}$$

(with formation of the spectrum at $k_{s\,max} r_{De} \sim 1/10$); this corresponds approximately to the experimentally observed spectrum of the sound oscillations. Then, assuming that $k_s r_{De} < 1$ and $v_0 > v_{Te}\omega_s/\omega_0$, we find from (A.1) and (A.4) that the increment of the effective collision frequency that determines the longitudinal conductivity of the plasma is:

$$\nu_{eff}^{p} \approx \frac{1}{10}\omega_0 u_s u_l \frac{\omega_{pe}^2}{\omega_0^2}\frac{\omega_{pe}^2}{\Omega_e^2}. \tag{A.5}$$

We note that the estimate (A.5) depends little on the concrete form of the frequency distribution $du_s/d\omega_s$ of the ion-sound energy.

For the same spectrum of the ion-sound oscillations we obtain from (A.2) and (A.3)

$$\nu_{eff}^{g} \approx \frac{1}{10}\omega_0 u_s \frac{\omega_{pe}^2}{\omega_0^2}\frac{\omega_{pe}^2}{\Omega_e^2} = \frac{\omega_{eff}^{p}}{u_l}.$$

1)We do not take into account in this estimate the voltage drop along the column as a result of the temperature difference. It is easily seen that allowance for this effect increases $\nu_{\rm eff}$. It turns out here that in the absence of an RF field $\nu_{\rm eff}$ is 2–3 times higher than the electron-neutral collision frequency.

2)Apparently, however, since their number is relatively low, the superthermal electrons make no substantial contribution to the total current, and hence also to the change of the longitudinal electric conductivity.

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