STUDY OF THE PROCESS OF TURBULENT HEATING OF A PLASMA BY A LARGE-AMPLITUDE WHISTLER

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It is shown that when a whistler of large amplitude is resonantly excited in a plasma, the electrons become rapidly heated during approximately half the period of the whistler. This decreases the amplitude of the whistler by a factor of several times. With decreasing whistler amplitude, a predominant heating of the ions is observed. It is shown theoretically that heating of the electrons can be due to their scattering by turbulent pulsations of magnetized ion-acoustic oscillations propagating almost perpendicularly to the constant magnetic field. The heating of the ions, and also the heating of the electrons at small amplitudes of the high-frequency field can be attributed to a stochastic mechanism that leads to a random change of the phases of the magnetized ion-acoustic oscillations or of the hydrodynamic turbulent pulsations at the characteristic frequency and with an increment on the order of $\sqrt{\omega_{He}\omega_{Hi}}$. The predominant heating of the ions at small amplitudes is connected with the influence of inelastic collisions that prevent the heating of the electrons.

1. INTRODUCTION

It is known that various oscillations can be excited in a plasma when an electric current flows perpendicular to the external magnetic field, provided the relative velocity of the electrons and the ions $u = u_i - u_e$ exceeds a certain critical value $u_{cr} \sim \sqrt{(T_e + T_i)/m_i}$. The scattering of charged particles by turbulent pulsations of the electric field, due to the development of these instabilities, leads to an effective heating of the plasma. Anomalous absorption of high-frequency energy and heating of the plasma in strong alternating electric fields perpendicular to the external magnetic field were observed in a number of experiments (see, for example,[1-6]). In these experiments, the effective frequency of the collisions, which determines the absorption of the high-frequency field in the heating of the plasma, greatly exceeded the frequency of the paired Coulomb collisions. Nor could the observed absorption of the energy of the high-frequency field be attributed to Cerenkov or cyclotron absorption by electrons and ions of the plasma.

A small-scale hydrodynamic instability of a plasma with transverse current was observed in experiments on plasma heating by a fast magnetosonic wave (whistler) of large amplitude[7,8] in a θ pinch[9], and in experiments with plasmoids entering into an inhomogeneous magnetic field[10]. Small-scale instabilities were also observed on the front of a collisionless shock wave.

In these investigations we observed predominant heating of either electrons or ions, or else the electrons and ions were heated equally in order of magnitude. The causes remained unclear. Nor is it sufficiently clear just which of the numerous current instabilities is responsible for the observed heating. Theoretical investigations show that various types of oscillations can be excited in a plasma with a transverse current[3,11-14]. The role of each of these instabilities in the problem of anomalous resistance and heating of the plasma can be clarified only after the nonlinear theory, the development of which has only been started[15-19] is completed.

In the present paper we show experimentally that in turbulent heating of a plasma by a fast magnetosonic wave (whistler) at low amplitudes, there is predominant heating of the ions, and at large amplitudes both the electrons and the ions are heated. In the latter case, the heating of the electrons occurs within a short time interval (during half the period of the oscillations of the whistler, $\pi/\Omega \approx 7 \times 10^{-8}$ sec). A theoretical analysis shows that the observed fast heating of the electrons may be due to plasma instability resulting from the excitation of "magnetized" ion-acoustic oscillations when the electrons move relative to the ions with velocity $u \gg u_{cr}$ perpendicular to the magnetic field. The characteristic frequencies and the growth increment of these oscillations are of the order of

$$\mathrm{Re}\,\omega \sim kv_s, \quad \mathrm{Im}\,\omega \sim \sqrt{\omega_{He}\omega_{Hi}} \tag{1}$$

(where $v_s = \sqrt{T_e/m_i}$ is the velocity of the ion sound), the wavelength is shorter than or of the order of the Larmor radius of the electrons $\rho_e = v_{Te}/\omega_{He}$ and is larger than the Debye radius of the electrons:

$$v_{Te}/u \gtrsim k\rho_e \gtrsim 1, \quad kr_D \lesssim 1,$$

the direction of propagation is almost perpendicular to the magnetic field:

$$(m_e/m_i)^{1/2} < \cos\theta \sim u/v_{Te} < 1,$$

the phase velocity of the oscillations along the magnetic field is of the order of the thermal velocity of the electrons ($|\omega|/k_\parallel \sim v_{Te}$).

Heating of the ions may be due to the stochastic mechanism of interaction of nonresonant ions with ion-acoustic and hydrodynamic oscillations, due to the finite time of correlation of the turbulent pulsations of the electric field[20]. For hydrodynamic oscillations, the frequency and the growth increment are of the order of $\sqrt{\omega_{He}\omega_{Hi}}$, $k\rho_e \ll 1$ and $\cos\theta \sim (m_e/m_i)^{1/2}$[13]. This mechanism is significant at small values of the current velocity and for heating of the electronic component of the plasma.

2. EXPERIMENTAL SETUP. MEASUREMENT METHODS

A detailed description of the experimental setup and the sequence of operations involved in the production of a plasma and its heating by the high-frequency field is given in[6]. The plasma was produced by a pulse discharge with oscillating electrons in hydrogen at pressures $\lesssim 10^{-3}$ mm Hg. The duration of the discharge current was 18 μsec. The inside diameter of the glass discharge tube was $2a = 6.6$ cm, the distance between the cathodes was 88 cm. The quasiconstant magnetic field H_0 with mirror geometry (mirror ratio 1.4) had a homogeneous section in the central part of length 70 cm. Depending on the initial hydrogen pressure and the intensity of the magnetic field, the plasma density n_0 obtained in the discharge, averaged over the radius, ranged from 10^{13} to 10^{14} cm^{-3}. The radial distribution of the plasma density n(r), measured with the aid of a double Langmuir probe, was well approximated by the linear function $n(r) = 2n_0(1 - r/a)$ at all the instants of time after the termination of the discharge current.

The source of the high-frequency energy was a surge circuit with a stored energy up to 10 J. The inductance of the circuit consisted of 8 individual sections enclosing the discharge tube and connected to the capacitors of the circuit pairwise in antiphase. The axial period of the electromagnetic field generated by the coil was $\Lambda = 20$ cm. The circuit oscillation frequency was $\Omega/2\pi = 7$ MHz, and the Q without the plasma was 35.

As shown in[6], at resonance of the excitation of the fast magnetosonic wave, i.e., under conditions when the longitudinal length of the fast magnetosonic wave in the plasma coincided with Λ, more than 60% of the energy stored in the circuit entered into the plasma, and the plasma (predominantly the electrons) was heated rapidly (<1 μsec) to a temperature ~ 100 eV. The alternating magnetic field on the axis of the system increased at resonance by approximately a factor of 2. As shown in[21], in the region of the resonance of the excitation of the plasma, a complicated spectrum of the oscillations was produced, connected with the nonstationary character of the process and with phenomena of nonlinear wave interaction.

We measured the electron density averaged over the radius with the aid of a microwave interferometer. As already indicated above, the local density could be determined with the aid of a double Langmuir probe, which was inserted into the plasma from the end of the discharge tube and could be moved along the axis of the system and also radially. Analogously, it was possible to introduce into the discharge a miniature high-frequency magnetic probe, with the aid of which we measured the alternating magnetic field $\widetilde{H}_z$ on the axis of the system. The time constant of the LR network made up of the probe inductance and the load resistance was 2 μsec, and ensured a signal proportional to $\widetilde{H}_z$ with sufficient accuracy.

Measurement of longitudinal energies of the electrons and the ions emerging from the plasma along the constant magnetic field was carried out with the aid of a three-grid analyzer with a retarding potential. This analyzer was placed on the axis of the system behind the magnetic mirror. An opening was made in the cathode to permit the plasma to pass into the analyzer. During the course of the measurements, we determined the dependence of the ion or electron current in the collector on the value of the retarding potential. From this dependence it was then possible to obtain, by differentiation, the distribution function with respect to the longitudinal velocities of the particles entering the analyzer, and then also the average longitudinal energy of these particles. The particle heating in the plasma could be evaluated from these data.

The density of the transverse energy of the plasma was estimated from the diamagnetic signal. The diamagnetic probe was a multiturn coil in an electrostatic screen, placed over the discharge tube between two neighboring sections of the discharge-circuit coil, at the place where the field $\widetilde{H}_z$ generated by the coil had a node. The signal induced in the probe was integrated by an LR network made up of the inductance of the probe and the probe load resistance. The time constant of this network was 80 μsec and greatly exceeded the duration of the investigated process.

The efficiency with which the HF energy was transferred from the coil into the plasma was characterized by a transfer coefficient α, representing the ratio of the energy absorbed in the circuit in the presence of the plasma, after subtracting the ohmic losses in the circuit itself, to the total energy stored in the capacitors of the circuit prior to its turning on. The value of α was determined with the aid of an instrument for the measurement of the damping decrement of the oscillations[22]. The pickup for this instrument was a high-frequency magnetic probe producing a signal proportional to the current in the circuit.

3. MEASUREMENT RESULTS

The experimental results described below were obtained by turning on the surge circuit several microseconds after the termination of the discharge current, at the instant when the decaying-plasma electron density, averaged over the radius, decreased to $n_0 \approx 3 \times 10^{13}$ cm^{-3}. The increment of the plasma density due to the additional ionization after turning on the circuit did not exceed 20% of the density at the instant when the circuit was turned on. The circuit capacitors were charged to 54 kV, corresponding to a stored energy 7.2 J. The maximum field amplitude $\widetilde{H}_z$ under the section of the coil was ~ 250 Oe in vacuum and 500 Oe in the plasma at the resonance of the excitation ($H_0 \approx 800$ Oe).

Figure 1 shows superimposed oscillograms of several quantities characterizing the heating of the plasma by the high-frequency field. The measurements were performed in the region of the excitation resonance. It follows from Fig. 1 that after the current is turned on in the circuit (Fig. 1a), the field $\widetilde{H}_z$ in the plasma (Fig. 1b) begins to increase, and the amplitude of the field in the second half-period of the oscillations is much larger than in the first, i.e., resonant buildup of oscillations in the plasma takes place. After the high-frequency field reaches the maximum value in the second half-period of the oscillations, its amplitude decreases sharply, reaching in the third half-period only $\sim 1/4$ of the maximum value. Simultaneously with the

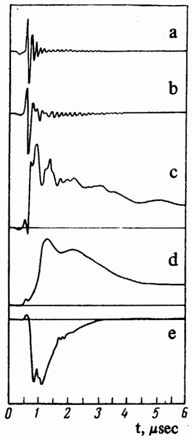

FIG. 1

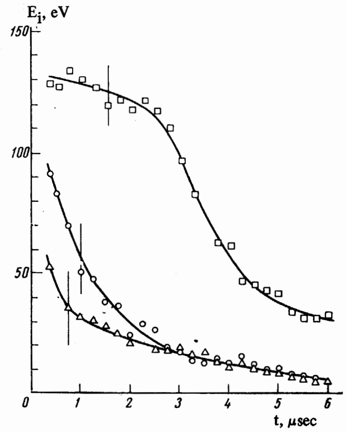

FIG. 2

FIG. 1. Oscillograms of the current in the surge circuit (a) the field in the plasma (b), the diamagnetic signal (c), the ion current (d), and the electron current (e).

FIG. 2. Dependence of the ion energy on the time elasped after turning on the high-frequency field: Δ–H_0 = 880 Oe, O–1210 Oe, □–1650 Oe.

appearance of the strong damping of the high-frequency field, there appears a rapid growth of the diamagnetic signal (Fig. 1c). (The high-frequency field generated by the surge circuit is first induced in the diamagnetic probe; this makes it possible to estimate the time of growth of the diamagnetic signal, which amounts to approximately half the oscillation period, i.e., 0.07 μsec.) The increase of the diamagnetism is due only to heating of the plasma, since no appreciable increase of the plasma density can occur within such a short time owing to the additional ionization produced by the heating of the electrons. In subsequent instants of time, there occur in the plasma, besides the continuing oscillations at the frequency Ω, also oscillations at other frequencies, which are seen not only on the oscillogram of the field $\tilde{H}_z$, but also in the diamagnetic signal. These oscillations were investigated in detail in[21]. It is interesting to note that after the sharp increase of the temperature of the plasma and the decrease of the amplitude of the oscillations at the frequency Ω, these oscillations subsequently, as follows from Fig. 1b, attenuate much more slowly.

The rapid damping of the oscillations in the plasma, accompanied by the plasma heating, can be regarded as a consequence of the sharp increase of the effective frequency of the particle collisions in the strong high-frequency field. The possible causes of such an increase of the effective collision frequency will be considered in Sec. 4.

The maximum plasma temperature, averaged over the radius and calculated from the diamagnetism signal with allowance for the distribution of the plasma density over the radius, amounted to $T_e + T_i \approx 150$ eV in our case. To determine how the ion and electron temperatures vary separately, let us turn to Figs. 1d and 1e, which show oscillograms of the ion and electron currents in the analyzer collector at zero retarding potential. It follows from these oscillograms that after the switching on of the high-frequency field there occurs a fast increase of the ion and electron currents. This increase of the current, as well as the increase of the diamagnetic signal, is due principally to the heating of the electrons and the ions. We see here that the electrons are heated simultaneously with the sharp decrease of the high-frequency field amplitude in the plasma and the increase of the diamagnetic signal. The ion current reaches a maximum value after only 1 μsec following the turning on of the high-frequency field or 0.5 μsec after the start of the electron heating. It is impossible, however, to determine the ion heating time exactly, since the time of delay of the maximum of the ion current is influenced not only by the finite heating time but also by the finite time of flight of the ions from the point of heating to the collector. (The corresponding distance is not known accurately and can be of the order of several centimeters, which at an ion velocity $\lesssim 10^7$ cm/sec, which obtained under the conditions in question, can result in an observable delay time.)

In order to estimate the ion temperature attained in this case, we used the fact that at later instants of time, when the high-frequency field in the plasma is small, the current of the electrons to the collector of the analyzer decreases more rapidly than the ion current, thus indicating a faster cooling of the electrons. By comparing the oscillograms c, d, and e we can assume that 2–3 μsec after the turning on of the high-frequency field the diamagnetism of the plasma is due mainly to the ions. Therefore the ion temperature determined from the diamagnetic signal is 30–40 eV (assuming that the plasma density is uniformly distributed over the radius and is equal to 3×10^{13} cm^{-3}; as already indicated, the true density could have been ~ 20% higher). This is in satisfactory agreement with the results presented below, obtained by directly measuring the ion energy by the method of the retarding potential.

It should be noted that turning off the magnetic mirrors did not lead to any appreciable change in the character of the oscillograms shown in Fig. 1, even at the initial instants of time, when the plasma temperature exceeded 100 eV. Such an insensitivity of the hot plasma to the magnetic mirrors is apparently due, like the fast heating, to the anomalously large effective collision frequency of the particles in the strong high-frequency field.

Let us now consider in greater detail the experimental data characterizing the heating of the ions by the high-frequency field. These data are based on measurements of the dependence of the ion current in the analyzer collector on the retarding potential at different instants of time after turning on the high-frequency field, and at different intensities of the constant magnetic field H_0. A computer reduction of the corresponding curves yielded the average longitudinal energy of the ions entering the analyzer. As is well known[23], the longitudinal-energy distribution of the ions passing through the input aperture of the analyzer is in general not identical with the distribution in the plasma. In addition, as already noted above, the distribution inside the analyzer could be greatly influenced by the finite time of flight of the ions. This influence should be particularly noticeable at the initial instants of time.

The dependence of the average ion energy on the time after turning on the high-frequency field is shown in Fig. 2 for three different values of H_0. The vertical

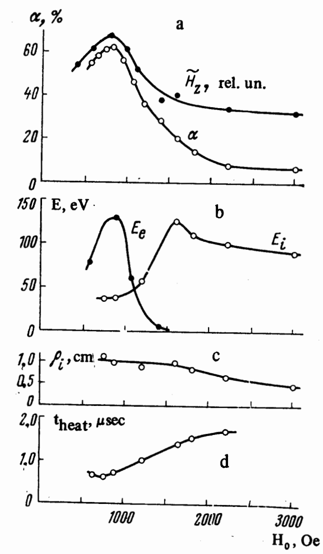

FIG. 3

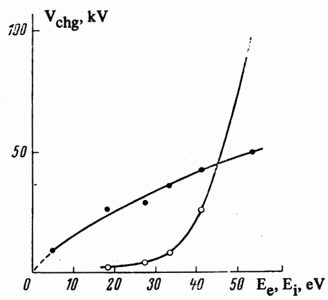

FIG. 4

FIG. 3. Transfer coefficient and amplitude of high-frequency field in the plasma (a), electron and ion energy (b), Larmor radius of the ions (c), and ion heating time (d) vs. the constant magnetic field.

FIG. 4. Electron (○) and ion (●) energy vs. charging voltage.

lines in this figure denote the instants of time when the current of the ions in the analyzer collector reaches the maximum value. It is natural to assume that the average energy of the ions entering the analyzer at this instant of time is closest to the maximum average energy acquired by the ions in the plasma by heating. The average ion energies obtained in this manner are indeed equal in order of magnitude to the ion energy in the plasma as determined from the diamagnetic signal.

The most remarkable fact concerning the heating of the ions is that the maximum of ion energy as a function of H_0 is reached not at the resonant value of H_0, when the high-frequency energy absorbed by the plasma, the amplitude of the high-frequency field in the plasma (Fig. 3a), and the electron energy (Fig. 3b) reach their largest values. As follows from Fig. 3b, the energy of the ions increases with increasing H_0 and reaches a maximum (130 eV at the maximum of the ion current) when H_0 is approximately double the resonant value. In the considered region of values of H_0, the Larmor radius of the ions heated by the high-frequency field is comparable with the radius of the discharge tube. Apparently this is precisely the circumstance that determines the upper limit of the ion energy. Indeed, as follows from Fig. 3c, the Larmor radius of the ion corresponding to the average energy plotted in Fig. 3b changes insignificantly at all the considered values of H_0, and remains at the level $\sim (0.2-0.3)$ a $\sim$ 1 cm.

Figure 3 also shows the dependence of the time of appearance of the maximum of the ion current relative to the instant of turning on the high-frequency field, as a function of H_0, for a zero retarding potential (Fig. 3d). It follows from Fig. 3d that this time increases as H_0 increases away from the resonance of the excitation (and apparently the same will also hold true with decreasing H_0). Obviously, this is due to the decrease in the heating rate, which in turn is due to the decreased amplitude of the high-frequency oscillations in the plasma (see Fig. 3a, which shows a plot of the amplitude $\tilde{H}_z$ in the second half-cycle as a function of H_0). As follows from Fig. 3c, the ions nevertheless have time to acquire an appreciable energy in the high-frequency field, for with increasing distance from resonance the damping coefficient of the oscillations in the resonant circuit decreases, and therefore the effective duration of the high-frequency pulse increases.

At the same time, as shown by measurements (Fig. 3b), the energy of the electrons, being largest at the excitation resonance, decreases rapidly with increasing distance from resonance, to a value on the order of several eV. Thus, at large H_0 practically the entire high-frequency energy absorbed by the plasma ($\alpha \sim 10\%$) goes to heating of the ions. It can be assumed that the rate of heating of the electrons, and consequently also the maximum electron temperature, depend strongly on the amplitude of the high-frequency field in the plasma. This assumption (as well as the hypothesis advanced above that the ion heating rate depends on the amplitude of the high-frequency field) is confirmed by the results of direct measurements of the dependence of the electron and ion energies on the charging voltage of the resonant circuit at the fixed value H_0 = 900 Oe, which exceeds somewhat the resonant value (Fig. 4). Under these conditions, the amplitude of the high-frequency field in the plasma during the first period of the oscillations is proportional to the charging voltage. As expected, with decreasing amplitude of the high-frequency field in the plasma, the electron energy decreases much more rapidly than that of the ions.

In light of the foregoing we can understand why appreciable heating of the electrons begins under conditions of excitation resonance (Fig. 1) only after the second half-cycle of the oscillations, namely, the amplitude of the high-frequency field in the first half-cycle is insufficiently large to ensure a large electron heating rate. Indeed, as follows from Fig. 1b, the amplitude of the high-frequency field in the second half-cycle exceeds the amplitude in the first half-cycle by 1.5 times. It follows from Fig. 4 that the corresponding electron energies can differ by one order of magnitude.

4. ION-ACOUSTIC TURBULENCE OF A PLASMA WITH TRANSVERSE CURRENT

Let us consider the turbulence of a plasma with developed ion-acoustic oscillations having a frequency and increment given by (1) and propagating almost perpendicular to the magnetic field ($\cos\theta \sim u/v_{Te}$), which may be responsible in our experiment for the turbulent heating of the plasma and for the rapid damping of the fast magnetosonic wave. When $k\rho_e \ll v_{Te}/u$, such oscillations are built up by the electrons under conditions of Cerenkov resonance, and when $k\rho_e \sim v_{Te}/u \gg 1$ excitation by electrons under conditions of cyclotron resonance becomes appreciable. Since $\cos\theta \sim u/v_{Te}$, it follows that the velocity of the resonance electrons responsible for the buildup of the oscillations is of the order of the thermal velocity, so that practically all the electrons take part in the buildup of the oscillations.

Let us estimate first the level of the turbulent noise established as a result of the development of oscillations with frequency and growth increment (1). We shall assume that the stabilization of the oscillations is due to nonlinear interaction of the waves with participation of the electrons. Nonlinear oscillations of the electrons are described by the kinetic equation for the oscillating part of the electron distribution function $\tilde{f} = f_e - f_0^{(e)}$:

$$\frac{\partial \tilde{f}}{\partial t} + \mathbf{v}\frac{\partial \tilde{f}}{\partial \mathbf{r}} - \frac{e}{m_e c}[\mathbf{v}\mathbf{H}_0]\frac{\partial \tilde{f}}{\partial \mathbf{v}} + \frac{e}{m_e}\nabla\varphi\frac{\partial f_0^{(e)}}{\partial \mathbf{v}} + \frac{e}{m_e}\left(\nabla\varphi\frac{\partial \tilde{f}}{\partial \mathbf{v}} - \left\langle \nabla\varphi\frac{\partial \tilde{f}}{\partial \mathbf{v}}\right\rangle\right) = 0; \quad (2)^*$$

$\langle\,\rangle$ denotes statistical averaging. The nonlinear term

$$\frac{e}{m_e}\nabla\varphi\frac{\partial \tilde{f}}{\partial \mathbf{v}} \sim \frac{ek_\perp}{m_e v_{Te}}|\varphi\tilde{f}|$$

in the strongly-nonlinear regime becomes of the same order as the linear term $\partial\tilde{f}/\partial t \sim \mathbf{k}\cdot\mathbf{u}\tilde{f}$, which is responsible for the Cerenkov absorption of the oscillations, when

$$\tilde{E} \sim k_\perp\varphi \sim (m_e/e)kuv_{Te} \sim (u/c)H_0 k\rho_e, \quad (3)$$

where $\mathbf{k}$ is the characteristic wave vector. In this case the pulsations of the electron and ion densities $\tilde{n}_{e,i}$ and of their average velocities $u_{e,i}$ are determined by the expressions

$$\tilde{n}_e \sim \tilde{n}_i \sim n(u/v_{Te}), \quad \tilde{u}_e \sim u^2/v_{Te}, \quad \tilde{u}_i \sim u(m_e/m_i)^{1/2}. \quad (4)$$

We note that the estimate (3) can also be obtained on the basis of the weak-turbulence theory[24,25], by equating the nonlinear damping decrement to the linear growth increment (1).

Let us consider electron heating in the field of the turbulent pulsations (3). To estimate the rate of heating of the electrons, we use the quasilinear equation for the function $f_0^{(e)}$. Taking the estimate (3) into account, we obtain for the characteristic time of variation of $f_0^{(e)}$ (the heating time) the following estimate:

$$1/\tau_e \sim \omega_{He}(u/v_{Te})^3. \quad (5)$$

The scattering by the oscillations leads to diffusion both parallel and perpendicular to the magnetic field, and the estimate (5) determines the increase of both the longitudinal and the transverse electron temperatures.

Using the quasilinear theory and the estimate (3), we can also easily obtain an estimate of the effective electron collision frequency, which determines the rate of change of the momentum of the electrons as they are scattered by the ion-acoustic oscillations (3):

$$\nu_{\text{eff}} \sim (u/v_{Te})\omega_{He}. \quad (6)$$

Let us now investigate the heating of the ions. Since the phase velocity of the ion-acoustic oscillations is much larger than the thermal velocity of the ions, the heating of the bulk of the ions is impossible under conditions of Cerenkov resonance. Owing to the presence of a strong nonlinear wave interaction leading to a "collapse" of the phases of the Fourier components of the field and to a weakening of their temporal correlation, the mechanism of stochastic heating of the ions becomes effective. Since $\tilde{u}_i \ll v_{Ti}$, the oscillating part of the distribution function of the ions is much smaller than their background distribution function $f_0^{(i)}$, which in this case is determined by an equation of the diffusion type[20]. Using this equation and recognizing that in the considered case of strong turbulence the characteristic correlation time of the quantities $\varphi_{\mathbf{k}}(t)$ will be of the order of the reciprocal nonlinear damping decrement, i.e., $\tau_{\text{corr}} \sim 1/\sqrt{\omega_{He}\omega_{Hi}}$, we obtain, taking (3) into account, the following estimate for the characteristic time of variation of the function $f_0^{(i)}$ (the ion heating time):

$$\frac{1}{\tau_i} \sim \frac{e^2}{m_i^2}\frac{\tilde{E}^2}{v_{Ti}^2\tau_{\text{corr}}\,\omega^2} \sim \omega_{Hi}\left(\frac{m_e}{m_i}\right)^{1/2}\left(\frac{u}{v_{Ti}}\right)^2. \quad (7)$$

In addition to ion-acoustic oscillations propagating almost perpendicular to the magnetic field ($\cos\theta \sim u/v_{Ti}$), there can be excited in this case also hydrodynamic oscillations and non-magnetized ion-acoustic oscillations (for the latter $\omega \sim \omega_{pi}$, $kr_D \sim 1$, $\cos\theta \sim 1$). The rate of the stochastic heating of the electrons and ions due to the finite correlation time of the hydrodynamic oscillations is also determined by the relation (7).

If the stabilization of the non-magnetized ion-acoustic oscillations is due to scattering of these oscillations by electrons, then their noise level is of the order of[19]

$$W = \tilde{E}^2/8\pi \sim (2m_e/m_i)nT_e.$$

In this case the electron heating time due to their scattering by the non-magnetized ion-acoustic oscillations is of the order of

$$\frac{1}{\tau_e^{(s)}} \sim \omega_{pe}\left(\frac{u}{v_{Te}}\right)^2\frac{W}{nT_e} \sim \omega_{pe}\left(\frac{u}{v_s}\right)^2,$$

and the rate of the stochastic heating of the ions is of the order of

$$\frac{1}{\tau_i^{(s)}} \sim \omega_{pi}\frac{u}{v_{Te}}\frac{W}{nT_i}.$$

Under the conditions of our experiments $\tau_{e,i}^{(s)} \gg \tau_{e,i}$, and the non-magnetized ion-acoustic oscillations could play a role only at $W/nT_e \sim 10^{-2}$.

We note that if stabilization of the ion-acoustic instability were to be due to nonlinear scattering of the oscillations by the ions, then[26]

$$W/nT_e \sim (u/v_{Te})(T_e/T_i).$$

This stabilization mechanism, however, is too weak and leads to values of ν_{eff} that exceed by two orders of magnitude the values of ν_{eff} calculated from the measured values of the electric conductivity of the plasma[27]

If the non-magnetized ion-acoustic oscillations make no appreciable contribution to the heating of the electrons and of the ions, and if the energy loss from the plasma can be neglected, then the picture of the plasma heating by an electromagnetic wave of large amplitude with frequency $\Omega \ll \sqrt{\omega_{He}\omega_{Hi}}$ can be represented in the following manner.

If at the initial instant of time the parameter

$$\varkappa = \frac{\tau_e}{\tau_i} \sim \frac{v_s}{u}\frac{T_e}{T_i} \quad (8)$$

is much smaller than unity, then the electrons are predominantly heated and the ratio T_e/T_i increases. As the plasma becomes heated, the quantity $v_s(T_e/T_i)$ increases, the parameter κ becomes larger than unity, and a stage of stochastic heating of the ions and electrons sets in, during which the temperature of the ions is "drawn" to the electron temperature. When the plasma temperature becomes high enough to make the

*$[\mathbf{v}\mathbf{H}_0] \equiv \mathbf{v}\times\mathbf{H}_0$.

thermal velocity of the order of the current velocity, the hydrodynamic stage goes over into the kinetic stage; the heating stops when the thermal velocity of the ions becomes smaller than the limiting velocity $u < u_{cr}$, and the instability terminates.

5. DISCUSSION OF RESULTS

The observed rapid heating of electrons can be explained as being due to their scattering by turbulent short-wave ion-acoustic oscillations propagating almost perpendicular to the magnetic field. Under the experimental conditions ($n_0 \approx 3 \times 10^{13}$ cm^{-3}, $\Omega = 4.4 \times 10^7$ sec^{-1}, $K_\parallel = 2\pi/\Lambda = 0.3$ cm, $K_r \approx 1$ cm^{-1}) with resonant excitation of the oscillations ($\tilde{H}_z \sim 500$ Oe, $u \approx u_e \sim 5 \times 10^7$ cm/sec at the point $r = a/2 = 1.6$ cm, $u_i \sim 6 \times 10^6$ cm/sec $\ll u_e$), the latter have, for a typical temperature $T_e \sim 100$ eV, a characteristic frequency and growth increment (1) equal to $\omega \sim 3 \times 10^9$ sec^{-1} and $\gamma \sim 3 \times 10^8$ sec^{-1}, a wavelength $\lambda = 2\pi u/\omega_{He} \sim 0.02$ cm ($kr_D \sim 1$) and a longitudinal wavelength $\lambda_\parallel \sim 2\pi\rho_e \sim 0.2$ cm.

A theoretical estimate of the time of electron heating by these oscillations (5) gives a value $\tau_e \sim 4 \times 10^{-8}$ sec, which is close to the observed heating time 7×10^{-8} sec. Since the estimate (5) is sensitive to the exact value of the ratio u/v_{Te}, one must not attach particular significance to the numerical equality of the estimated and measured times τ_e.

The whistler damping time τ_w can be estimated if it is recognized that the energy goes to electron heating:

$$\frac{1}{\tau_w} = \frac{8\pi}{\tilde{H}_z^2}\frac{3n}{2}\frac{dT_e}{dt} \sim \frac{u}{v_{Te}}\Omega\sqrt{\left(\frac{K_r}{K_\parallel}\right)^2+1}. \tag{9}$$

The wave-amplitude damping time estimate from (9), $\sim 6 \times 10^{-8}$ sec, is close to the observed value. We note that the time of the whistler damping due to the Cerenkov absorption is larger by one order of magnitude than the observed value.

The time of electron heating as a result of scattering by non-magnetized ion-acoustic oscillations, $\tau_e^{(s)}$, is also larger by one order than the measured value when $W/nT_e \sim 10^{-3}$. (These oscillations can account for the observed electron heating rate only at a level on the order of $W/nT_e \sim 10^{-2}$.)

At large amplitudes of the high-frequency field ($\tilde{H}_z \sim 500$ Oe), the plasma heating occurs within a time much shorter than the time of the electron energy loss due to inelastic collisions. Indeed, as seen from Fig. 1e, the time of cooling of the electrons is ~ 1 μsec. At small amplitudes of the high-frequency field, the rate of heating of the electrons decreases strongly ($\tau_e > 10^{-6}$ sec at $\tilde{H}_z \lesssim 200$ Oe), so that an increase of the electron temperature becomes impossible because of strong losses in inelastic collisions. (We note that the strong influence of inelastic collisions on the electron temperature may be the cause of the rapid decrease of T_e with increasing pressure of the neutral gas, even at relatively large values of the amplitude of the high frequency field $\tilde{H}_z = 350$ Oe[28].)

Let us discuss the results on ion heating. Recognizing that the heating of the ions became limited when their Larmor radius became comparable with the radius of the chamber ($\rho_i = v_{Ti}/\omega_{Hi} = 0.2$–$0.3a$), and using relation (7), we obtain the following estimate for the time of the stochastic heating of the ions on magnetized ion-acoustic and hydrodynamic pulsations:

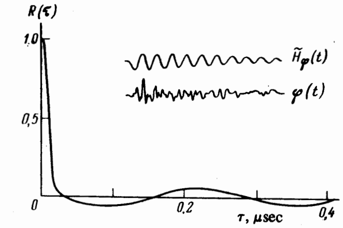

FIG. 5

$$\Delta t \sim \frac{\omega_{Hi}}{\bar{u}^2}\left(\frac{m_i}{m_e}\right)^{1/2}\left(\frac{a}{3}\right)^2, \tag{10}$$

where $\bar{u}$ is the mean value of u(t) during the heating time. Stochastic heating of ions in our experiments can take place under non-resonant conditions ($\tilde{H}_z \sim 200$ Oe), and also under resonance conditions, after the whistler amplitude decreases to approximately one-quarter. The estimate (10) yields in this case $\Delta t = 1.5 \times 10^{-6}$ sec, which does not contradict the experimental data.

Relations (7) and (10) were obtained under the assumption that the correlation time of the turbulent pulsations is of the order of $1/\sqrt{\omega_{He}\omega_{Hi}}$. This assumption was verified by us for the hydrodynamic pulsations[7,8]. The data of[8] were used to construct the autocorrelation function R (see Fig. 5) for one of the realizations of the oscillation potential $\varphi(t)$, shown on the same figure together with an oscillogram of the azimuthal component of the magnetic field of the whistler. For these data, the potential $\varphi(t)$ had a maximum of the spectral intensity at $\omega \approx 1.6 \times 10^8$ sec$^{-1} \approx 0.5\sqrt{\omega_{He}\omega_{Hi}}$ ($H_0 \approx 900$ Oe, $n_0 = 10^{13}$ cm^{-3}). It follows from Fig. 5 that the correlation time is equal to approximately 10^{-8} sec, i.e., it is indeed of the order of $1/\sqrt{\omega_{He}\omega_{Hi}}$.

Thus, the predominant rapid heating of the electrons at large amplitudes of the high-frequency field can be attributed, in accordance with (8), to the fact that in this case the electron heating by scattering by the magnetized ion-acoustic pulsations is faster than the stochastic heating of the ions by the magnetized ion-acoustic and hydrodynamic pulsations. At low amplitudes of the high-frequency field, the electron heating should occur, according to the previously developed theory, at an approximately equal rate. On the other hand, the predominant heating of the ions at low amplitudes of the high-frequency field is due to the fact that the electrons lose a large fraction of the energy acquired in the high-frequency field to inelastic collisions, whereas the role of the inelastic collisions for the ions is not large in the considered range of pressures.

In conclusion we note that the concepts presented above can also be used to explain the results of experiments on turbulent plasma heating[3,4].

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