A COLLISIONLESS THERMAL WAVE IN A PLASMA

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Experiments are described on the collisionless propagation of heat in a plasma along the magnetic field. Thermal waves can be propagated in a medium whose thermal conductivity is a power function of the temperature. In a collisionless plasma, where the mean free path of particles is much greater than all typical lengths, in particular the length of the equipment, the heat must be propagated by a different method. Experiments to study this phenomenon showed that heat is propagated along the magnetic field with velocity exceeding that of ion-acoustic velocity (I. A. velocity), and that the spatial width of the thermal disturbance is much less than the mean free path. Heat is propagated because hot electrons are replaced by cold [1]. Noise was observed experimentally in the vicinity of the ion plasma frequency and an estimate of its intensity was obtained. Theoretical discussion showed that the I. A. velocity instability which develops at the wave front leads to collective friction of the cold electrons with the ions and makes it possible to find the effective collision frequency. It was also shown theoretically that, in accordance with experiment, noise is localized near ω_{pi}, and the level agrees with that obtained experimentally. The phenomenon can be pictured as follows: hot electrons expanding into the region occupied by cold electrons and ions create an electric field. Cold electrons, accelerating in this field, oscillate the I. A. velocity. This instability leads to heating of the electrons and the appearance of collective friction which forms the heat front.

1. To explain the mechanism of the propagation of heat along the magnetic field, we first consider the following formulation of the problem.

Suppose that at time $t = 0$ the left half-space $x < 0$ is occupied by a "hot" plasma in which the transverse $T_\perp$ and longitudinal $T_\parallel = T$ temperatures are of the same order, while the half-space $x > 0$ is occupied by a "cold" plasma with temperature $T_x \ll T$. The temperature of the ions is assumed to be zero, and the magnetic field is directed along the x axis. At later moments of time the hot particles enter the half-space $x > 0$ because of the pressure gradient. If there were neutral particles, collisions would cause free dispersion. In the case under consideration there is no free dispersion. As soon as the hot particles move to the right, an electric field is created which must slow down the hot electrons. If the cold electrons could, for some reason, remain motionless (for example, if there were a large friction force for cold electrons due to Coulomb collisions with the ions), there would be formed an electric field which would balance the pressure gradient,and the hot electrons would stop moving.

Since the cold electrons can move, they will accelerate in the electric field which is formed, and move to the left. After the cold electrons have moved to the left, the hot electrons can move farther, since the electric field becomes smaller.

Thus, a peculiar wave of displacements, or a collisionless thermal wave, is formed. The velocity of such a wave is defined by the flow rate of the cold electrons. We must make the following observation here: the velocity of the cold electrons is by no means $v_{Tx} \sim \sqrt{T_x/m}$. It is considerably greater and is determined by the potential difference which the cold electron passes through. If the potential difference is of the order T_e^{-1}, the wave velocity must be of the order $\sqrt{T/m} \gg C_S$.

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We seek a wave with a typical scale length much greater than the Debye radius and we consider its evolution with time much greater than ω_{pe}. To describe the steady-state process we can use the kinetic equation without the collision term for the hot electrons and the equations of hydrodynamics for the cold electrons. Since we are considering scales much greater than the Debye radius, instead of Poisson's equation, we can use the quasi-neutrality equation. If we take the following as the initial conditions: at time $t = 0$ the hot electrons are in the region $x < 0$ with concentration n_0, while in $x > 0$ the cold electrons have the same concentration, the system of equations, as in the case of dispersion of the plasma in a vacuum [2], have a self-similar solution. This is because the equations do not contain parameters with the dimensions of length or time.

We introduce the variables

$$\tau = \frac{x}{t}\left(\frac{m}{2T}\right)^{1/2}, \qquad v = v\left(\frac{m}{2T}\right)^{1/2}, \qquad v_x = v_x\left(\frac{m}{2T}\right)^{1/2}$$
$$\varphi = \frac{e\varphi}{2T}, \qquad n = \frac{n}{n_0}, \qquad f = \frac{1}{n_0}\left(\frac{m}{2T}\right)\sqrt{\pi}$$

If we eliminate the variables φ, v_x, and n by using the condition $\varphi = 0$, $v_x = 0$ at $\tau = -\infty$, we obtain the equation

$$(v-\tau)\frac{\partial f}{\partial \tau} - \frac{1}{n}\frac{\partial n}{\partial \tau}\left(\frac{1}{n}\int_{-\infty}^{\tau} n\,d\tau\right)^2 \frac{\partial f}{\partial v} = 0 \tag{1.1}$$
$$n = 1 - \frac{1}{\sqrt{\pi}}\int_{-\infty}^{\infty} f\,dv$$

Denoting the coefficient of $\partial f/\partial v$ in (1.1) by $F(\tau)$, we obtain the equation for the characteristics

$$\frac{dv}{d\tau} = -\frac{F(\tau)}{v-\tau} \tag{1.2}$$

Along the lines $v(\tau)$, defined by (1.2), f has a constant value. The function $F(\tau)$ is nonnegative, $\partial \ln n/\partial \tau \geq 0$, which is confirmed by numerical computation. Hence it follows from (1.2) that when $v > \tau$, the derivative $dv/d\tau < 0$, while when $v < \tau$, we have $dv/d\tau > 0$.

We can say that when $v > \tau$ the hot particles slow down, providing their longitudinal energy to create the potential. The hot particles, reflected from the wave front when $v = \tau$ accelerate in the matched field. Already this indicates that the pattern of the phenomenon is more complicated than that which was first described. In particular, the hot particles may be reflected by the potential created by the same particles, as happens, for example, when the electrons in the field of a heavy ion are in a Boltzmann distribution. Hence, we cannot say that the velocity of motion of the front is of the order $(T/m)^{\frac{1}{2}}$, since not only displacement of the cold electrons by hot ones, but also the reflection of hot electrons in the matched field takes place. The line $v = \tau$ in the $v\tau$ plane is peculiar. We shall investigate this important moment in more detail. We begin with the case when $F(\tau) = 0$, which corresponds to free dispersion. Then the equation for the characteristics (1.2) takes the form

$$\frac{dv}{d\tau}(v-\tau) = 0 \tag{1.3}$$

The solution of (1.3) is trivial; either $v = \tau$, or $v = \text{const}$ when $v \neq \tau$. If $f = f_M$ when $\tau = -\infty$,

$$f(v,\tau) = f_M(v)\,\theta(v-\tau), \quad \theta(v-\tau) = \begin{cases} 1 \text{ for } & v > \tau \\ 0 \text{ for } & v < \tau \end{cases} \tag{1.4}$$

since when $\tau = -\infty$, $f_M(v)$ lies above the line $v = \tau$ (there are no particles with velocity $v = \pm\infty$). Thus, the characteristics do not approach the region $v < \tau$. Now we consider the behavior of the characteristics when $F(\tau) \neq 0$ in the neighborhood of the point $v = \tau$. Let some characteristic intersect the line $v = \tau$ in the neighborhood of the point $\tau = \tau_0$. We replace $F(\tau)$ by $F(\tau_0)$ approximately. Then the equation for the characteristics takes the form

$$\frac{dv}{d\tau} = -\frac{F(\tau_0)}{v-\tau}, \quad \text{or} \quad \frac{d(v-\tau)}{d\tau} = -1 - \frac{F(\tau_0)}{(v-\tau)}$$

For the initial condition $v = \tau_0$ when $\tau = \tau_0$, the solution is as follows:

$$\tau = v + F(\tau_0)\{1 - \exp[(v - \tau_0)/F(\tau_0)]\} \tag{1.5}$$

i.e.,

$$\frac{dv}{d\tau} = -\frac{F(\tau_0)}{v - \tau} = \frac{1}{1 - \exp[(v - \tau_0)/F(\tau_0)]}$$

It follows from this equation that when $F(\tau_0) \to 0$, we have

$$dv/d\tau \to 1 \quad \text{for} \quad v < \tau_0, \qquad dv/d\tau \to 0 \quad \text{for} \quad v > \tau_0.$$

Then, as must be the case, as $F(\tau_0) \to 0$, we obtain the case of free dispersion.

Now let $F(\tau_0) \neq 0$, but $F(\tau_0) < \tau_0$, which holds at the leading and trailing parts of the wave, where the electric field is small. If the point τ_0 is left in such a way that $|\tau - \tau_0| > F(\tau_0)$ (but, of course, $|\tau - \tau_0| < |\tau_0|$), then it follows from (1.5) that

$$\tau = v + F(\tau_0)$$

i.e., the characteristic approaches a line parallel to the straight line $v = \tau$ and passing below the latter by an amount $F(\tau_0)$ (Fig. 1). When $\tau \to -\infty$, we have $F(\tau_0) \to 0$, and so the characteristics approach the line $v = \tau$. Thus, the characteristics intersecting the line $v = \tau$ at distinct points τ, have the asymptote $v = \tau$. The characteristics need not intersect at any point since $dv/d\tau$ increases as $|v - \tau|$ increases. It is interesting to note that when $\tau = -\infty$, accelerated particles, moving to the left, must appear. Since all the characteristics approach the one line $v = \tau$, the number $\int f dv$ of such particles is negligible.

Numerical computation is necessary to obtain the exact solution of the problem. We note that, as distinct from [1], the computation is more complex, since for a given value of τ, v has two values (Fig. 1), due to the reflection of particles from the front. When $\tau = -\infty$ (it is sufficient that $\tau = -2$) the distributions of the particle velocities and the angles of the characteristics to the τ axis are Maxwellian. When using the fact that the distribution function does not change along the characteristic, we can find $F(\tau)$, and thus $dv/d\tau$ at the point $\tau = -2 + h$, where h is the step length. It should be noted that to each of the n characteristics leaving the region $v > \tau$ there corresponds its continuation in the region $v < \tau$, in reverse order with respect to the line $v = \tau$ (which is their asymptote) (Fig. 1). The form of the distribution function for the hot electrons and various values of τ is shown in Fig. 2. We see that there can be free dispersion only for $\tau \gtrsim 0.5$, and that less than 10% of the particles disperse freely. The greater part of the hot electrons are trapped by the matched potential. Approximately half the hot electrons disperse with velocity $0.2(2T/m)^{\frac{1}{2}}$, ($\tau = 0.2$), i.e., as already noted at the outset, the wave velocity is significantly less than v_{Te}, but nevertheless greater than C_S.

The distribution function we have obtained for the electron velocity is unstable both with respect to the buildup of Langmuir oscillations and with respect to the buildup of I. A. velocity since, first, the cold and hot electrons move with respect to each other and, second, the electrons move with respect to the ions.

If we consider the one-dimensional case ($\omega_{He} \gg \omega_{pe}$), the effect of such instabilities on the velocity of motion of the wave is insignificant. Indeed, from one-dimensional quasi-linear theory [3] in the region between v_x and $\langle v^\circ \rangle = \int f(v)dv$ the distribution function for the hot electrons has a plateau due to the buildup of plasma oscillations, while for $v_x \leq v \leq 0$ the distribution function for the cold electrons has a plateau due to the buildup of I. A. velocity. Then as before, $T > T_x$. In this region of τ, where there is no free dispersion and where the particles are basically concentrated, this does not lead to a marked change in the velocity of the wave. For those values of τ for which between the distribution functions for the hot and cold electrons there is a region of values of v where $f = 0$ (for $\tau > 0.3$), the establishment of a plateau for the distribution function of the hot electrons leads to a reduction in the average velocity of the group of hot electrons dispersing freely, by approximately a factor of 2. However, because there are very few freely dispersing particles, this has no effect on the velocity of the fundamental mass of electrons and so no effect on the nature of the wave motion. The scattering of cold electrons by I. A. velocity oscillations when $\omega_{He} < \omega_{pe}$ has a much more significant effect. We know that it leads to a restriction in the velocity of the cold electrons to an amount $\alpha\sqrt{T_x/M}$, where α is a constant, greater than unity [4]. At the same time the cold electrons are heated, i.e., T_x increases. We are not interested in analyzing the dispersion relation. As for

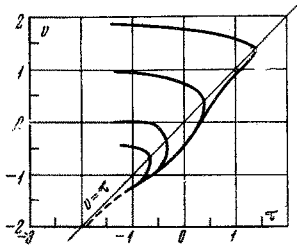

Fig. 1

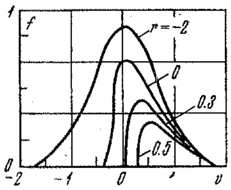

Fig. 2

the hot electrons, even without taking the instabilities into account, their distribution is almost Maxwellian, and so their density is related to the matched potential as follows:

$$n^{\circ} \sim \exp\,(e\varphi\,/\,T)$$

The collective friction of the cold electrons with the ions leads to an increase in the absolute value of the electric field and so of the variable F in (2.2). The greater the value of F, the greater the slope of the characteristics $dv/d\tau$ for $v > \tau$ and the fewer the freely dispersing particles, and so the nearer the density distribution of the hot electrons to the Boltzmann distribution.

Now we consider a formulation of the problem which is more in accord with experiment by taking account of collective effects. Suppose that at the initial moment of time the region $-L/2 < x < L/2$ is occupied by hot electrons (L is the width of the region near the boundary, ~10 cm) of concentration n_0 equal to the concentration of the ions. The remaining part of the space $|x| > L/2$ is occupied by a cold plasma of the same concentration. To determine the motion of the wave we use the following equation of motion for the cold electrons:

$$v_x = \alpha \sqrt{T_x/M} \qquad (1.6)$$

(which is equivalent to the equation $v_x = eE/m\nu_{eff}$, where ν_{eff} is the effective collision frequency), the equation of continuity, and also the equation defining the temperature increase

$$\partial T_x\,/\,\partial t = ev_x\,\partial\varphi\,/\,\partial x \qquad (1.7)$$

It follows from Eqs. (1.6) and (1.7), and the determination of the frequency of effective collisions that

$$\nu_{eff} \sim \frac{1}{t}\,\frac{M}{m\alpha^2}$$

i.e., the effective frequency decreases with time, which is a consequence of the increase in the temperature of the cold electrons T_x. The hot electrons can be assumed to have a Boltzmann distribution

$$n^{\circ} = n_0 \exp\,[e\varphi\,/\,T\,(x)] \qquad (1.8)$$

where T(x) is defined by the adiabatic equation

$$T\,(x) = T\,(L\,/\,x)^{a-1}$$

where a is the adiabatic index. If we assume that the energy of the hot electrons is not transformed from transverse into longitudinal motion, then $a = 3$ (one-dimensional adiabatics).

Indeed, the longitudinal temperature $T_{\parallel}^{\circ}$ decreases since the energy of the longitudinal thermal motion is expended on creating the electric field and the acceleration of the cold electrons, and hence quite rapidly $T_{\parallel}^{\circ}$ becomes smaller than the transverse temperature of the hot electrons $T_{\perp}^{\circ}$. In these conditions ($T_{\parallel}^{\circ} < T_{\perp}^{\circ}$) instabilities may develop at the electron cyclotron frequency ω_{He}, leading to equality of the transverse and longitudinal temperatures [5]. If the increment in this instability is greater than the typical inverse times for the motion of the wave front, $T_{\parallel}^{\circ} \approx T_{\perp}^{\circ}$ and the adiabatic index $a = 5/3$. We need one further equation for the conservation of the hot electrons [1]

$$\int_{-\infty}^{\infty} n^{\circ} dx = n_0 L \qquad (1.9)$$

From this it follows that

$$n^{\circ} \ll n_x, \quad n_x \sim n_0$$

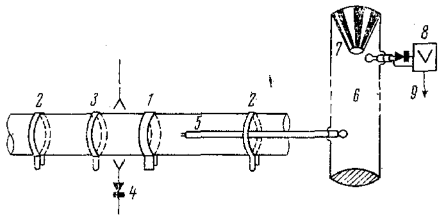

Fig. 3

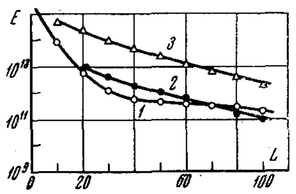

Fig. 4

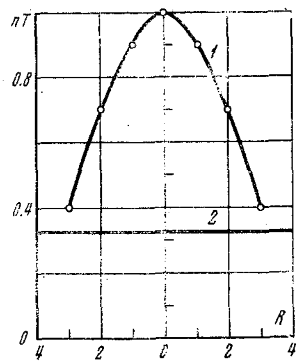

Fig. 5

for $x \gg L$, since

$$n^\circ + n_x = n_0 \tag{1.10}$$

To determine the law of motion we have to express n_x and v_x in terms of n° in the equation of continuity (3) of [1]. We seek the solution for $x \gg L$. Since the velocity of the wave x/t is the velocity of the hot electrons, $v_x = (n^\circ/n_x)x/t$, and so $tv_x/x = n^\circ/n_x \ll 1$. Hence the equation of continuity can be transformed as follows:

$$\frac{\partial n_x}{\partial t} + \frac{\partial}{\partial x}(n_x v_x) = -\frac{\partial n^\circ}{\partial t} + \frac{\partial}{\partial x}(n_0 v_x) - \frac{\partial n^\circ}{\partial x} v_x - n^\circ \frac{\partial v_x}{\partial x} = 0 \tag{1.11}$$

The third and fourth terms are less than the first and second, respectively, by a factor of $(n^\circ/n_0)^{-1}$, and so they can be neglected. If we now express v_x in terms of n° in (1.11), we obtain

$$\frac{\partial^2 n^\circ}{\partial t^2} = \frac{\partial}{\partial x}(n_0 \alpha C_{s_x}^{\cdot}) = \frac{\partial}{\partial x}\left\{\frac{n_0 \alpha^2}{M} \frac{\partial}{\partial x}\left[T\left(\frac{L}{x}\right)^{a-1} \ln \frac{n^\circ}{n_0}\right]\right\} \tag{1.12}$$

If we now take account of (1.9), from which it follows that $n^\circ \sim L/x$, we can find the relation between x and t, or the self-similar variable

$$x \sim L\left(\frac{\alpha\sqrt{T/M}\,t}{L}\right)^{2/a} \tag{1.13}$$

If $a = 5/3$, $x/t \sim t^{1/5}$ remains almost constant with time, while for $a = 3$, $x/t \sim t^{1/3}$. In the theory of turbulent heating α is not defined precisely. We know only that it must be greater than unity and less than $\sqrt{M/m}$ [4]. Putting $\alpha = 6$, $a = 5/3$, we can obtain satisfactory agreement with experiment.

Considering the effect of instabilities on the nature of the wave motion, we dwell only on the I. A. noise and ignore the effect of plasma oscillations with frequency $\omega = \omega_{pe}$ and increment

$$\gamma \sim \omega_p \frac{n^\circ}{n_0}\left(\frac{x}{t v_{Te}}\right)^3$$

At the very beginning of the process the increment is less than the maximum increment in the I. A. velocity instability and so the plasma oscillations can be ignored. After effective collisions between the fundamental mass of electrons and ions of frequency ν_{eff} appear in the system the plasma waves are not built up if their increment $\gamma < \nu_{eff}$. It is easy to see that this condition holds during the time defined by the inequality

$$\gamma < \nu_{eff}$$

i.e.,

$$t\omega_{pe}\left(\frac{1}{L}\, t\left(\frac{T}{M}\right)^{1/2}\alpha\right)^{1/a} < \left(\frac{M}{m}\frac{1}{\alpha^2}\right)^{1/2}$$

This relation holds for typical values of the parameters in a laboratory experiment.

2. The apparatus for studying thermal waves (Fig. 3) consisted of a cylindrical evacuated volume – a glass tube of diameter 8 cm and length 300 cm in a homogeneous magnetic field. The length of the solenoid in which the quasi-constant field was created is 120 cm, the duration of magnetic field pulse of semi-sinusoidal form was 15–30 msec. The intensity of the field at the moment of generating the thermal wave was 1 kG (in the homogeneous part).

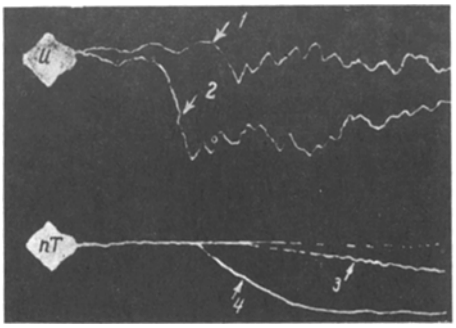

Fig. 6

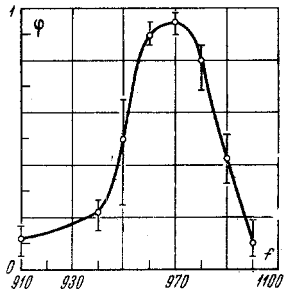

Fig. 7

The evacuated volume was pumped out to a limiting pressure of $5 \cdot 10^{-7}$ torr. The working pressure of the gas (hydrogen or argon) was $2\text{-}4 \cdot 10^{-4}$ torr. The volume was initially ionized by two high frequency generators (2) of frequency 20 MHz and pulse power 100 kW. In this way the plasma was established with concentration $n = 2 \cdot 10^{12}\text{-}7 \cdot 10^{1}$ cm^{-3}, electron temperature T_{e0} = 0.5-10 eV, and ion temperature T_i of the order 2-3 eV, the degree of ionization reaching 70% .

At the center of the evacuated volume (Fig. 3) was a shock ring in the form of a narrow coil (diameter 8 cm, length 2.5 cm), on the tube and supplied from a charged condenser through a low inductance discharger.

The current in the shock circuit had the form of a damped sinusoid with characteristic frequency 11 MHz and duration 1 μsec. The maximum value of the magnetic field intensity due to the current was of the order of 1800 G at the center of the coil.

The pulsed magnetic field excited in the plasma column an oblique nonlinear magnetosonic wave. The characteristics of this wave are described in detail in [6]. Here it is essential only to note that intensive dissipation of this wave leads to heating of the electron component of the plasma in a quite restricted segment near the shock ring. The duration of the heating of the region near the coil is very small and lasts for only a few tens of nanoseconds.

The following apparatus was used to control the parameters:

1) the concentration, averaged over the diameter, was determined by a high-speed interferometer at a wavelength of λ = 8 mm;

2) the magnetic field of the wave was studied by a magnetic probe of small dimensions introduced into the plasma;

3) the plasma pressure nT was determined either by an external diamagnetic sensor averaging the pressure over the cross section and fed to the discharge chamber from outside or by an internal sensor introduced into the plasma;

4) the noise was measured by twin electric probes;

5) to determine the radiation of the plasma an antenna system in the microwave frequency range (4-16 mm) was used.

The sequence of operations in the experiments was as follows: first the quasi-static magnetic field H_0 and at the same time the high-frequency generators were switched on after preliminary ionization of the gas; some tens of μsec later, after the generators had been switched on, there was a breakdown and a build-up of the concentration which was fixed by the 8-mm diagnostic apparatus (4). Then, at a given moment of time which was chosen on the basis of the requirement on the initial parameters n and H_0, the shock circuit was switched on and the plasma heated to a sufficiently high temperature (to several hundreds of eV) by the magnetosonic wave in the region near the circuit.

The plasma pressure was measured by an external or internal diamagnetic sensor (3) which could move along the discharge volume. Comparison of the signals from the diamagnetic sensor and the magnetic probe made it possible to determine the true behavior of nT as a function of the distance from the central plane of the shock circuit. The distributions of the energy of the wave $H^2/8\pi$, the noise $E^2/8\pi$, and the plasma nT, along the length of the plasma column are shown in Fig. 4 (1, 2, and 3, respectively). The energy (E) is in eV/cm^3, and the length (L) in cm. At 20 cm from the central plane of the shock circuit nT is greater than the wave pressure $H^2/8\pi$ and is equal to 10^{16} eV/cm^3. The radial distribution (Fig. 5) obtained from the internal diamagnetic sensor with an annular coil has its maximum on the chamber axis and is close to the function $J_0(kr)$; the accuracy with which the results coincide with measurements from the internal (1) and external (2) sensors when averaged was not worse than 5%. The internal diameter of the hole in the coil of the sensor introduced into the plasma was not less than 12 mm, which is an order of magnitude greater than the Larmor diameter of the electrons for the given experiment. In Fig. 5, nT is in relative units, R is in cm.

By moving the sensors along the tube a series of oscillograms was obtained, making it possible to estimate the velocity of propagation of the pressure along the magnetic field and the typical scale of the nT front.

Control experiments in which the concentration was fixed on the pressure front using a high-speed interferometer, $\lambda = 8$ mm, made it possible to prove that the perturbation in the concentration which occurs there is less than 10% of the original value. Hence, nT is transported by the thermal wave, moving with velocity several times greater than that of I. A. velocity $\sqrt{T_e/M}$, and approaching the thermal velocity of the electrons (cf. [1]). The physical nature of this phenomenon can be explained by investigating the noise at the thermal front. To this end a series of measurements were made in the range of the plasma electron frequency ω_{pe}. The measurements were made both inside and outside the plasma column. In the former case a sensitive pin dielectric antenna was introduced from the end of the chamber into the plasma volume. This series of experiments showed that no significant epithermal radiation was observed in the frequency range ω_{pe}-$2\omega_{pe}$.

On the other hand, a direct experiment with twin electric probes showed that there is intense noise in the region of ω_{pi}. In this case the form of the experiment is as follows (Fig. 3). The signal from the twin electric probes (5) was fed through a cable to a loop which excited a wave-guide filter (6) with critical wavelength 30 cm. The wave-guide filter was a copper tube of diameter 18 cm, closed at both ends by copper flanges. From one side a matched load (7) in the form of an aquadag-coated truncated cone was introduced into the volume to ensure the propagation of the running waves. The excitation loop which was loaded by the twin electric probes had no galvanic contact with the wave-guide. The construction of the receiving loop was similar, the signal from it being passed through a detector to an amplifying circuit (8, 9) (the band of the circuit was $\Delta f = 80$ MHz).

It was thus possible to search for noise in a wide frequency range ($\omega > 6 \cdot 10^9$ sec^{-1}) and reliably screen the receiving circuit from noise generated by the shock ring and also by the nonlinear waves in the plasma. Since it has previously been proved that there is no marked level of radiation in the region of ω_{pe}, to a high probability the noise extracted by the twin electric probes has to be associated with the frequency range ω_{pi}.

Figure 6 shows oscillograms of the noise (u) obtained at distances of 30 (2) and 60 (1) cm from the central plane of the coil. The duration of the scan was 50 nsec/cm. The velocity of the front of the noise signal corresponds to the velocity nT (3, 4), i.e., noise appears at a given point only with the arrival of the thermal front. The intensity of the noise falls as the wave amplitude decreases.

The noise spectrum was investigated in more detail using a narrow-band tunable resonator with $\Delta f =$ 20 MHz (Fig. 7). The plasma density was $n \sim 1.7 \cdot 10^{13}$ cm^{-3}, the amplitude of the oscillations φ was in relative units, the frequency in MHz. In this case it was possible to show that the frequency range in which noise occurs is quite narrow, $\Delta\omega \ll \omega_{pi}$. But the frequency dependence of the noise can only give in particular qualitative information about the spectrum since the point $\omega \sim \omega_{pi}$ is a boundary point at which the static volt-ampere characteristic can be used to describe the dynamic properties of the probe [7, 8].

An absolute estimate of the spectral density of the noise can be obtained by other, more graphic, representations if the coefficient connecting the probe with the potential wave is characterized by the geometrical coefficient $k^2 l^2$, in accordance with [9], where l is the length of the probe, and k is a typical wave number of the potential oscillation. Taking account of this coefficient, we find that the noise energy $W = E^2/8\pi$ is $\sim 10^{13}$ eV $\cdot$ cm^{-3} (Fig. 4).

Control experiments with argon showed that the typical noise frequency shifts approximately in accordance with $\sqrt{M(H)/M(Ar)}$. In addition, a magnetic probe showed that in the region of ω_{pi} there was no marked noise level of a magnetic nature.

Thus, it can be assumed that at the thermal front there is an intense background of ionic-acoustic noise which ensures effective friction between the electrons and the ions.

From this we can obtain an experimental value for the degree of turbulence at the wave front as the ratio W/nT; in the experimental conditions it was equal to 0.01.

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