Determination of plasma temperature from cyclotron absorption in an inhomogeneous magnetic field

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Cyclotron oscillations accompanying nonmonotonic variation of the magnetic field in a region occupied by a plasma are considered. It is shown that at a frequency somewhat lower than the minimal cyclotron frequency (or higher than the maximal one) the oscillation absorption coefficient depends strongly on the plasma temperature. This uncovers a possibility of determining the plasma temperature by measuring the absorption coefficient. The proposed procedure was tested on the simplest plasma object, a gas-discharge plasma. The results agree well with data obtained by the probe method. It is assumed by the authors that the proposed method of determining the plasma temperature can be used also in thermonuclear systems.

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INTRODUCTION

The known method of determining plasma temperature from the Doppler broadening of the cyclotron-absorption line (see, e.g., [1]) can be used if the inhomogeneity of the magnetic field is small, and the change $\delta\omega_e$ of the cyclotron frequency in the region occupied by the plasma is small with the characteristic Doppler shift kV_T. Here k is the wave number of the oscillations and V_T is the thermal velocity of the electrons (for the sake of argument, we consider the electronic component of the plasma).

Systems intended for magnetic containment of a plasma are as a rule characterized by the inverse ratio of $\delta\omega_e$ and kV_T ($\delta\omega_e \gg kV_T$). In this case the absorption coefficient does not depend at all on the electron temperature and is determined by the magnetic-field gradient and by the plasma density. (The absorption coefficient was calculated in[2], see also the review[3], and was experimentally determined in[4].) This statement, however, is valid only if the cyclotron-resonance point at which the condition $\omega=\omega_e(z)$ is satisfied is not an extremum point of the magnetic field.

It is shown in the present paper that if the magnetic field has an extremum within the confines of the system, then the absorption coefficient of oscillations of frequency close to an extremal cyclotron frequency depends substantially on the electron temperature. Consequently, the results of the experimental measurement of the absorption coefficient contain information on the electron temperature, and generally speaking make it possible to determine its value.

Assume, for example, that the magnetic field has a minimum at a certain point inside the region occupied by the plasma (see Fig. 1). Such a magnetic-field configuration is known to be typical in adiabatic traps. Let the oscillation frequency ω be lower than the minimal cyclotron frequency ω_{e0}, but let the difference $\omega_{e0}-\omega$ be small, so that $\omega_{e0}-\omega \approx kV_T$. Although in this case the resonance condition $\omega=\omega_e(z)$ is not satisfied anywhere, the wave will interact with electrons that move towards it and for which, owing to the Doppler shift, the wave frequency $\omega'=\omega-kV$ becomes comparable with or larger than ω_{e0}. It is obvious that with decreasing ω the number of electrons interacting with the wave, and hence also the absorption coefficient $\eta(\omega)$, decreases in accordance with a law that depends on the electron temperature. From the character of the $\eta(\omega)$ dependence we can thus assess the electron temperature. This temperature can, of course be determined also from a single measurement of the absorption coefficient at a certain value of ω. But this calls in addition for exact knowledge of ω_{e0}, which is sometimes quite difficult, for example because of the diamagnetism of the plasma.

It should be noted that in an inhomogeneous magnetic field, in addition to the Doppler effect, the smearing of the cyclotron absorption line is caused also by the finite time of the resonant interaction. In fact, at resonance in an inhomogeneous magnetic field, when the resonance condition $\omega=\omega_e(z)$ is strictly speaking satisfied at only one point, the electron interacts with the oscillations for a certain finite time δt (see, e.g., [3]). During this time the electron cannot "distinguish" between oscillations whose frequencies differ from ω_{e0} by $\delta\omega \approx \delta t^{-1}$. The effects connected with the finite time of the resonance interaction were taken into account in an analysis[5] of the stability of a plasma in an adiabatic trap, and also in an investigation[6] of the cyclotron absorption of electrons in an installation of finite dimensions.

In the theoretical part of this paper we calculate the absorption coefficient of oscillations with $\omega<\omega_{e0}$. The calculations are made under the assumption that $\eta \ll 1$,

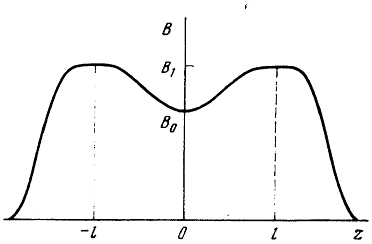

FIG. 1. Distribution of the magnetic field along the system axis. The plasma occupies the region $-l<z<l$.

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when the influence of the resonance interaction on the propagation of the oscillations can be taken into account via a small correction. We have actually calculated only the "tail" of the cyclotron absorption line at rather large differences $\Delta\omega = \omega_{e0} - \omega$. But this is fully adequate for the determination of the electron temperature.

What was determined directly in the experiment was not the absorption coefficient η, but the transmission and reflection coefficients τ and ρ, respectively. The absorption coefficient was then calculated from the formula $\eta = 1 - \tau - \rho$. In our case the reflection of the oscillations is due to the inhomogeneity of the medium in which they propagate (inhomogeneity of the density or of the magnetic field). An analysis shows that the reflection coefficient can be reliably calculated only if the condition $C = \omega_p^2 L/(\omega\Delta\omega)^{\frac{1}{2}} c \ll 1$ is satisfied, where ω_p is the plasma electron frequency and L is the characteristic dimension over which the magnetic field changes. Actually the experiments were performed under conditions when the reflection coefficient was negligibly small.

In the second part of the paper we describe experiments on the determination of the temperature of the electrons in a gas discharge in an inhomogeneous magnetic field. The magnetic field increased in the direction along the system axis from the center towards the edges of the region occupied by the plasma. The discharge tube was placed in a waveguide in which electromagnetic oscillations were excited. We measured the absorption coefficient of the oscillations as a function of the minimal value of the magnetic field. (Using the fact that the expression for the absorption coefficient contains not the frequency ω itself, but the difference $\Delta\omega = \omega_{e0} - \omega$, we varied, at a fixed value of ω, the value of the magnetic field.) From the law governing the decrease of the absorption coefficient with increasing ω_{e0} we determined the electron temperature. The obtained values of the temperature agree well with the results of probe measurements. For the experiments we chose a gas discharge, since gas-discharge plasma is the simplest and most convenient object for the investigation. At the same time, we see no objections, in principle, to the use of the proposed method in thermonuclear systems.

1. THEORY

A. Propagation of oscillations

We consider electromagnetic oscillations propagating in a plasma along a magnetic field. If the electric vector of the oscillations rotates in the same direction as the electrons in the magnetic field, but the resonance conditions $\omega = \omega_e(z)$ is not satisfied anywhere, then the wave equation takes the form (see, e.g., [2,3]):

$$\frac{d^2E_-}{dz^2} + \frac{\omega^2}{c^2}\left(1 + \frac{\omega_p^2}{\omega(\omega_e(z) - \omega)}\right)E_- = 0. \tag{1}$$

We use here a Cartesian coordinate system with the z axis directed along the magnetic field. The dependence of the cyclotron frequency on the z coordinate is chosen in the form $\omega_e(z) = \omega_{e0}(1 + z^2/L^2)$, the oscillation frequency is assumed to be close to ω_{e0}, $\omega < \omega_{e0}$, $E_- = E_x - iE_y = E_-(z)e^{-i\omega t}$. At a low plasma density, when the condition $C \ll 1$ is satisfied, the second term in the parentheses in (1) can be treated as a small correction. Neglecting this term in the first approximation, we obtain for a wave propagating from left to right $E_-(z) \approx e^{i\omega z/c}$. If the reverse condition is satisfied, $C \gg 1$, then the influence of the plasma can be taken into account in the quasiclassical approximation. In this case

$$E_-(z) \approx k^{-1/2}\exp\left(i\int^z k(z)dz\right), \quad k(z) = \frac{\omega}{c}\left(1 + \frac{\omega_p^2}{\omega(\omega_e(z) - \omega)}\right)^{1/2}. \tag{2}$$

B. Absorption of oscillations

As noted in the Introduction, the electrons can exchange energy with the oscillations even if the condition $\omega = \omega_e(z)$ is not satisfied anywhere within the confines of the system. We assume that the frequency difference $\Delta\omega = \omega_{e0} - \omega$ is large enough to make the resonance effects small. In this case we can assume in the first approximation that the oscillations take the form of a traveling wave, and the resonance effects can be taken into account by successive approximations.

We write the equation of motion of an electron in the field of a circularly-polarized electromagnetic wave in the form

$$\dot{V}_- + i\omega_e(z(t))V_- = -\frac{eE_-}{m}\exp(-i\omega t + ikz(t)). \tag{3}$$

Here $V_- = V_x - iV_y$, and the relativistic effects are not taken into account. From (3) we obtain

$$V_- = -\frac{eE_-}{m}\int^t dt' \exp\left(-i\omega t' + ikz(t') + i\int_t^{t'} \omega_e(t'')dt''\right). \tag{4}$$

Assuming the electron velocity along the magnetic field of the wave to be constant, we obtain from (4) the change δV_- after passage through the minimum of the magnetic field (see also, e.g., [7]):

$$\delta V_- = \frac{ieE_-}{m}e^{i\varphi}\frac{\pi^{1/2}}{\omega^{1/3}}\left(\frac{L}{V}\right)^{2/3}\mathrm{Ai}\left(\left(\frac{L\omega}{V}\right)^{2/3}\frac{\Delta\omega + kV}{\omega_{e0}}\right). \tag{5}$$

Here Ai is the Airy function, V is the longitudinal velocity of the electron, and φ is the phase of the oscillations at the instant of the passage through the minimum of the magnetic field. In the derivation of (5) the wave number was assumed for simplicity to be constant. To take into account the dependence of k on the coordinate in (5) it is necessary to make the substitution

$$L \to L\left(1 + \frac{V}{2c}\left(\frac{\omega_p}{\Delta\omega}\right)^2\left(1 + \frac{\omega_p^2}{\omega\Delta\omega}\right)^{-1/2}\right)^{-1/2}. \tag{6}$$

The Airy function attenuates exponentially at positive values of the argument and oscillates at negative ones, the amplitude of the oscillations decreasing in power-law fashion with increasing absolute value of the argument. If the argument of the Airy function is negative, this means that when account is taken of the Doppler shift the oscillation frequency $\omega' = \omega - kV$ exceeds ω_{e0}. It is obvious that such oscillations can interact with the

Larmor rotation of the electron. Since the plot of the cyclotron frequency against the longitudinal coordinate is a parabola, the resonance condition $\omega' = \omega_e(z)$ is satisfied at two points symmetrically located on opposite sides of the magnetic-field minimum. At definite distances between the resonance points, the phase of the cyclotron oscillations will differ by $(2k+1)\pi$ at the instants of their passage. In this case, the contributions of the resonance points to V_- will cancel each other. With increasing $\omega' = \omega - kV$, the resonance point move away from the minimum of the magnetic field, and this shortens the resonance zone. According to[3], the dimension of the resonance zone is

$$\delta z \approx \begin{cases} L^{2/3} r_e^{1/3}, & 1 < \dfrac{\omega'}{\omega_{e0}} \leqslant 1 + \left(\dfrac{r_e}{L}\right)^{2/3} \\ r_e^{1/2} \left(\dfrac{1}{B}\dfrac{dB}{dz}\right)_{z_S}^{-1/2}, & \dfrac{\omega'}{\omega_{e0}} \geqslant 1 + \left(\dfrac{r_e}{L}\right)^{2/3} \end{cases} \tag{7}$$

where z_S is the resonance point, $r_e = V_T/\omega_{e0}$, $V_T = (2T_e/m)^{1/2}$ is the average thermal velocity of the electrons along the magnetic field. Since the time of the resonance interaction becomes shorter, V_- changes by a smaller amount. This explains the decrease of the oscillation amplitude with increasing modulus of the argument of the Airy function.

At positive values of the argument of the Airy function we have $\omega' < \omega_{e0}$ and energy exchange between the oscillations and the electrons is due to effects due to the finite time of the resonance interaction. The characteristic scale of the decrease of δV_- with decreasing ω' is equal to the reciprocal of the electron flight time through the resonance zone at $\omega' = \omega_{e0}$.

When the electron passes through the minimum of the magnetic field the energy of the Larmor rotation changes by

$$\tfrac{1}{2} m \left(|\delta V_-|^2 + \mathrm{Re}\, V_- \delta V_-^*\right). \tag{8}$$

Averaging with respect to time cancels the second term in the parentheses and we obtain for the energy absorbed by the electrons per unit time the expression

$$\delta\varepsilon = \tfrac{1}{2} m \int dV\, n_0 f_0(V)\, V |\delta V_-|^2. \tag{9}$$

Owing to the rather complicated $\delta V_-(V)$ dependence (see (5)) the integral in (9) can be calculated only by making some simplifying assumptions. We assume first that the characteristic scale for the change of δV_-, namely $\omega k^{-1}(r_e/L)^{2/3}$, is small compared with the thermal velocity V_T of the electrons. Averaging over the fast oscillations of $|\delta V_-(V)|^2$, we put

$$|\delta V_-(V)|^2 \approx N\theta(\omega - \omega_{e0} - kV),$$
$$N = \frac{\pi^2}{2\cdot 3^{4/3}\Gamma^2(2/3)} \left(\frac{eE}{m}\right)^2 \frac{1}{\omega^{2/3}} \left(\frac{L}{V}\right)^{4/3}, \tag{10}$$

$\theta(x) = 1 (x > 0)$ or $\theta(x) = 0 (x < 0)$. This representation actually means that the only electrons interacting with the oscillations are those for which the oscillation frequency, with the Doppler shift taken into account, exceeds ω_{e0}. Assuming a Maxwellian electron velocity distribution, we get from (9)

$$\delta\varepsilon \approx \frac{\pi^{5/2}}{6^{8/3}\Gamma^2(2/3)} \frac{e^2E^2 n_0 V_T}{m} \omega^{-2/3} \left(\frac{L}{V_S}\right)^{4/3} \exp(-V_S^2 V_T^{-2}). \tag{11}$$

Here $V_S = \Delta\omega/k$. The oscillation absorption coefficient η is equal to the ratio of the energy absorbed per unit time to the energy flux in the oscillations. Using the expression for the flux (see, e.g., [8]), we obtain

$$\eta \approx \frac{\pi^{5/2}}{2^{5/3}3^{8/3}\Gamma^2(2/3)} \frac{\omega_p^2 V_T}{c^2} (k\omega)^{1/3} \left(\frac{L}{\Delta\omega}\right)^{4/3} \exp\left(-\left(\frac{\Delta\omega}{kV_T}\right)^2\right)$$
$$k \cong \max\left(\frac{\omega}{c}, \frac{\omega_p}{c}\left(\frac{\omega}{\Delta\omega}\right)^{1/2}\right). \tag{12}$$

If the inverse condition $\omega k^{-1}(r_e/L)^{2/3} \gg V_T$ is satisfied, then the effect of the finite time of the resonance interaction prevails over the Doppler effect. The characteristic scale of the variation of $|\delta V_-|^2$ as a function of V greatly exceeds V_T everywhere in the region $V \gg V_T$. Calculating (9) by the saddle-point method, we arrive at the expression

$$\eta \approx \frac{\pi^2}{6^{11/6}} \frac{\omega_p^2 L}{kc^2} \left(\frac{\omega}{\Delta\omega}\right)^{1/2} \exp\left(-3^{1/2}2^{1/3}\left(\frac{L\omega}{V_T}\right)^{2/3} \frac{\Delta\omega}{\omega}\right). \tag{13}$$

C. Reflection of oscillations

It can be shown that the reflection connected with the resonance interaction is negligibly small, and therefore the reflection coefficient can be obtained with the aid of the conservative equation (1). If the magnetic field has a parabolic dependence on the coordinate, this equation does not reduce to the standard one, and we must use approximate methods to calculate the reflection coefficient ρ.

Assume that the condition $C \ll 1$ is satisfied. In this case the "plasma" term in (1) can be taken into account as a small correction within the framework of the so-called Born approximation (see, e.g., [9]). The standard procedure yields:

$$\rho = \frac{\pi^2}{4} \frac{\omega_p^4 L^2}{\omega\Delta\omega c^2} \exp\left(-4\frac{\omega L}{c}\left(\frac{\Delta\omega}{\omega}\right)^{1/2}\right). \tag{14}$$

At $C \gg 1$ the quasiclassical approximation can be used. To calculate the reflection coefficient in this case we shall use a method proposed in[10]. Following [10], we consider the analytic continuation of (1) to the complex z plane. The quasiclassical approximation ceases to hold at the points z_0 and z_S at which the wave number $k(z)$ becomes respectively equal to zero and infinity. We draw from these points lines on which only the real part of the complex "phase" of the quasiclassical solution changes. In Fig. 2 these are the line Az_0B and the segment z_0z_S. Since z_S is a branch point of the solution, a cut should be drawn from it. In this case it is convenient to draw it vertically upward (wavy line). The solution of interest to us, at large positive values of z, takes the form of a wave traveling to the right. We continue it to the left along the line 1–10. In the vicinity of the points z_0 and z_S, where the quasiclassical approximation is violated, Eq. (1) reduces to Bessel equations.

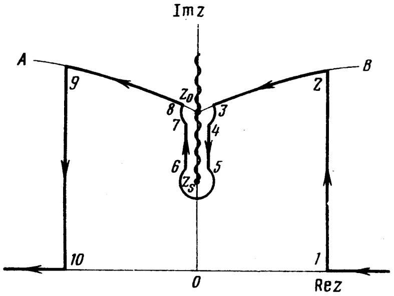

FIG. 2. Complex z plane for Eq. (1). The solution is continued from the region Re$z > 0$ into the region Re$z < 0$ along the path 1–10.

It can be found with their aid that circuiting around the points z_0 and z_S changes the quasiclassical representation of the solution—it acquires terms corresponding to the reflected waves. On the real-phase lines their amplitude is comparable with that of the "incident" wave. On segment 9–10, however, the "incident" wave increases exponentially, while the "reflected" one decreases. As a result, the reflection coefficient is exponentially small:

$$\rho=\left(\frac{\pi}{2}\right)^4\sin^4\Phi\exp(-4A),\qquad \Phi=\frac{\omega L_1}{c}\left(\bar{K}'\left(\frac{L_0}{L_1}\right)-\bar{E}'\left(\frac{L_0}{L_1}\right)\right),$$
$$A=\frac{\omega L_1}{c}\bar{E}\left(\frac{L_0}{L_1}\right),\qquad L_0=L\left(\frac{\Delta\omega}{\omega}\right)^{1/2},\qquad L_1=L\left(\frac{\Delta\omega}{\omega}+\frac{\omega_p^2}{\omega^2}\right)^{1/2}. \tag{15}$$

Here Φ is the phase shift of the quasiclassical solution on the segment z_0z_S, A is the change of the imaginary part of the phase on going from the real axis to the line AB, while $\bar{K}'$, $\bar{E}'$, and $\bar{E}$ are complete elliptic integrals (see, e.g., [11]).

The reflection coefficient vanishes at $\Phi = n\pi$. In this case the waves "reflected" by the points z_0 and z_S cancel each other. A similar phenomenon is known in optics as bleaching, and in the theory of electron scattering by atoms it is called the Ramsauer effect (see, e.g., [12]).

We have approximated the dependence of the magnetic field on the coordinate by a parabola. This approximation is suitable at low values of z, i.e., at the center of the system, but is obviously violated in the region of the mirrors. A more realistic relation is

$$B(z)=B_0\left[1+\frac{z^2}{L^2}\left(1+\alpha\frac{z^2}{L^2}\right)^{-1}\right]. \tag{16}$$

The reflection coefficient can be obtained in this case from (15) by the redefinition

$$L_0\to L_0\left(1+\alpha\frac{\Delta\omega}{\omega}\right)^{-1/2},\quad L_1\to L_1\left[1+\alpha\left(\frac{\Delta\omega}{\omega}+\frac{\omega_p^2}{\omega^2}\right)\right]^{-1/2},$$
$$c\to c\left(1+\alpha\frac{\Delta\omega}{\omega}\right)^{1/2}\left[1+\alpha\left(\frac{\Delta\omega}{\omega}+\frac{\omega_p^2}{\omega^2}\right)\right]^{-1/2}. \tag{17}$$

To find the new value of the reflection coefficient (14), it must be represented in the form

$$\rho=(\pi/2)^2(1-L_0^2/L_1^2)\exp(-4\omega L_1/c) \tag{18}$$

and the same substitutions must then be made. If α is of the order of unity, as is typical of real traps, then (15) can be significantly changed, whereas (14) remains practically unchanged. The last fact follows from the condition $\omega_p^2/\omega^2 \ll 1$, which ensures satisfaction of the inequality $C \ll 1$. The reason for this difference between (14) and (15) is the following: At $C \ll 1$ the oscillations are reflected by a region in which the coefficient of Eq. (1) changes most strongly, i.e., by the vicinity of the magnetic-field minimum. But if $C \gg 1$, then the reflection takes place in a region where the quasiclassical theory parameter $k^{-2}dk/dz$ is maximal. In view of the increase of k as $z \to 0$, this region need not necessarily coincide with the vicinity of the origin. Thus, to calculate the reflection coefficient at $C \gg 1$ it is necessary to take into account the real profile of the magnetic field. In addition, it is necessary also to take into account the reflection from the plasma density gradient. All this makes the calculations quite difficult.

2. EXPERIMENT

A. Description of setup

The experiments were shown with the setup shown in Fig. 3. The plasma was produced in a cylindrical glass tube (P) by a high-frequency generator with a symmetrical output with respect to the ground. The generator voltage was applied to two thin copper electrodes (E) wrapped around the tube near its ends. The electrodes were 1.5 cm wide and 0.02 cm thick. The generator power was 100 W and its frequency 40 MHz. The ends of the discharge tubes were circular cones 3 cm high. The length of the uniform part of the tube was 50 cm and its diameter 2.8 cm. The tube was filled with helium, argon, krypton, and mercury vapor. The operating pressures ranged from 5×10^{-2} to 1×10^{-4} mm Hg. The plasma parameters were determined with an electric double probe perpendicular to the tube axis. The length of the active part of the probes was 3 mm and the probe diameter was 0.1 mm. The probe measurements have shown that the electron temperature varies approximately from 1.5 to 20 eV, and the plasma concentration does not exceed 5×10^9 cm^{-3}. The degree of ionization of the gas was $\sim10^{-4}$.

The discharge tube was placed inside a rectangular

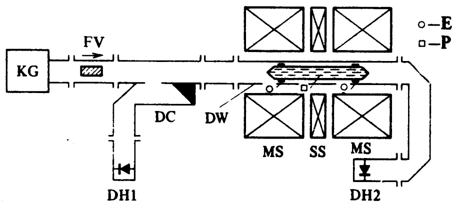

FIG. 3. Diagram of experimental setup: KG—klystron generator, FV—ferrite valve, DC—directional coupler, DH—detector head, DW—dismountable waveguide, E—RF electrodes, P—plasma in dielectric vessel, MS—multisection solenoid, SS—supplementary solenoid.

waveguide (DW) of 72×34 mm cross section, in which an electromagnetic wave of type H_{10} was excited at a fixed frequency 2.92 GHz by a 0.5 mW klystron generator (KG). The parasitic deviation of the generator frequency was 0.01%. The plasma-sounding microwave signal was amplitude-modulated at a frequency 1 kHz.

A magnetic field of mirror geometry was produced with a multisection solenoid (MS) of total length 100 cm. A typical distribution of the magnetic field along the axis of the plasma column is shown in Fig. 1. The plasma occupied in this case the region $-l<z<l$. The maximum magnetic field in the experiments was 2 kG. The field inhomogeneity over the cross section of the tube did not exceed 0.2%. The magnetic-field stability was 0.1%. Largest value of the mirror ratio $R=B_1/B_0$ reached in the experiments was 1.7. The dependence of the magnetic field on the coordinate z in the region of the minimum of $B(z=0)$ could be approximated with sufficient accuracy by the expression $B=B_0(1+z^2/L^2)$. Under the experimental conditions the value of L changed from 10 to 20 cm. In the experiments, provision was made for slight (on the order of 4%) local variation of the magnetic field at a frequency $\Omega=50$ Hz in the region of the minimum distribution $B=f(z)$. This was done by adding a supplementary solenoid section (SS), shown in Fig. 3. The minimal electron cyclon density then varied in accordance with the law

$$\tilde{\omega}_{e0}=\omega_{e0}+\frac{eB_g}{mc}\cos\Omega t.$$

The alternating magnetic field B_g was measured with a magnetic probe. Signals from the same probe synchronized the horizontal sweep of an S8-2 two-beam oscilloscope. The vertical inputs of the oscilloscope were generated by the detector sections (DH) and fed through V8-6 signal-ratio meters.

B. Results of measurements

Figure 4 shows an oscillogram of the output signal from detector head DH-2. It represents the dependence of the power of the transmitted electromagnetic field, at a fixed frequency, on the alternating magnetic field $B_g\cos\Omega t$ produced by the supplementary solenoid. During the sweep time, marked on the oscillogram by the numbers 1 to 4, the auxiliary magnetic field changed from $-B_g$ to $+B_g$. The value of the minimal stationary magnetic field B_0 was chosen to produce electron cyclotron resonance at the minimum of the magnetic field and was equal to $B_0=mc\omega/e=1.04$ kG. The mirror ratio was $R=1.3(L=14$ cm). The envelope sections marked on the oscillograms by the numbers 1–2, 2–3, and 3–4 correspond to the conditions $\omega>\tilde{\omega}_{e0}$, $\omega\approx\tilde{\omega}_{e0}$, and $\omega<\tilde{\omega}_{e0}$ respectively. The segment 1–2 of the curve is characterized by the presence of two resonance points separated in space. Segments 2–3 is characterized by the absence of points of exact cyclotron resonance ($\omega=\tilde{\omega}_{e0}$). This segment of the oscillogram makes it possible to obtain the profile of the cyclotron-absorption line in the plasma in the region where the minimum of the magnetic field is located. Segment 3–4 of the oscillogram corresponds to the conditions at which the cyclotron resonance in the plasma no longer takes place and the electromagnetic wave propagates through it without absorption.

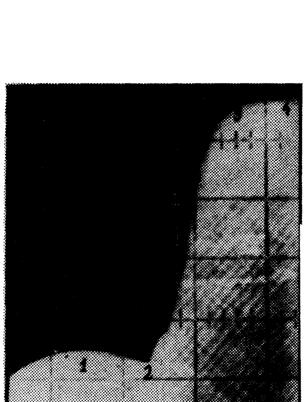

FIG. 4. Oscillogram describing the variation of the transmission coefficient τ with changing magnetic field $B_g\cos\Omega t$.

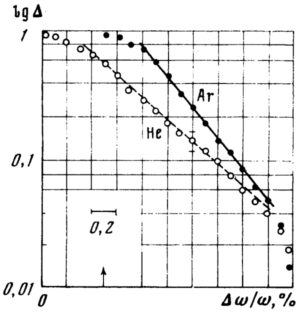

FIG. 5. Plot of log Δ against $\Delta\omega/\omega$.

Simultaneously with the transmitted wave power, we registered also the reflected power, using a directional coupler (DC). The experiments have shown that the reflected power depends on the electron density and on the value of L. The reflection of the wave becomes noticeable only at electron densities higher than 1×10^9 cm^{-3} and at values of L smaller than 10 cm. The reflected power, however, did not exceed 10% of the incident-wave power. Under the conditions when oscillograms similar to that shown in Fig. 4 were obtained, the reflected signal was negligibly small. This circumstance makes it possible to use the oscillogram of Fig. 4 to obtain information on wave absorption in the plasma. Inasmuch as under the conditions of the experiment it can be assumed that $\eta=1-\tau$, the wave-power absorption coefficient in the plasma is directly proportional to the ordinate Δ measured from the continuation of the horizontal part of segment 3–4 of the curve corresponding to the condition $\tau=1$. To measure the electron temperature it is necessary to know the slope of log η as a function of $\Delta\omega/\omega$. Since $\eta\sim\Delta$, this slope is exactly equal to the slope of the plot of log Δ against $\Delta\omega/\omega$.

Figure 5 shows the plots of log Δ obtained in this manner as a function of the quantity $\Delta\omega/\omega$, which characterizes the detuning from electron-cyclotron resonance ($\Delta\omega=\tilde{\omega}_{e0}-\omega$). The results are presented for a plasma produced in helium and argon at a pressure 1×10^{-3} mm Hg. The resonance detuning is reckoned from point 2 of Fig. 4. For the case of the argon plas-

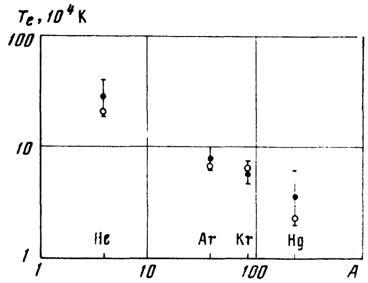

FIG. 6. Electron temperatures T_e in inert-gas and mercury-vapor plasma, measured by means of cyclotron resonance (●) and with probes (○).

ma, $\Delta\omega/\omega$ in Fig. 5 is reckoned from the arrow. It is seen from Fig. 5 that the plots of log Δ against $\Delta\omega/\omega$ are linear in a certain interval of $\Delta\omega/\omega$. In accord with formula (13), this dependence can take place at a fixed frequency ω if the cyclotron-absorption line profile is determined at the minimum of the magnetic field by the finite character of the time of the resonance interaction of the particles with the wave. To this end, as shown in Sec. 1, it is necessary to satisfy the inequality

$$D=\frac{\omega}{kV_T}\left(\frac{r_e}{L}\right)^{2/3}>1.$$

Under the experimental conditions we have

$$k=\frac{\omega}{c}>\frac{\omega_p}{c}\left(\frac{\omega}{\Delta\omega}\right)^{1/2}$$

and $D=(c^3/\omega^2 V_T L^2)^{1/3}$. The value of D calculated with allowance for the variation of the electromagnetic-wave velocity in the waveguide ranged from 1.5 to 3 in the experiments.

For the linear segment of the plot of log Δ against $\Delta\omega/\omega$, shown in Fig. 5, the condition $C<1$ is satisfied. As follows from the theoretical part of the paper, the wave reflection can then be described by formula (14). Calculation by this formula shows that, in accord with experiment, the reflected wave power is negligibly small at electron densities lower than 1×10^9 cm^{-3}.

In the theoretical analysis of the cyclotron absorption of the oscillations the plasma was assumed to be collisionless. In the gas-discharge plasma used in our experiments, however, it is necessary in general to take into account the collisions of the electrons with the neutral atoms. Under the influence of these collisions, just as under the influence of the effects described in the preceding section, the cyclotron-absorption line becomes broadened. This broadening by electron-neutral collisons was observed in our study at pressures higher than a certain limit that depended on the type of gas. No line broadening was observed, however, at pressures lower than 2×10^{-3} mm Hg in any of the investigated cases

Knowledge of the slope of the linear part of the plot of log Δ against $\Delta\omega/\omega$, shown in Fig. 5, makes it possible to estimate the electron temperature in the plasma by formula (13). The filled circles in Fig. 6 represent the values of T_e obtained in this manner for inert-gas and mercury plasmas. The results shown in the figure were obtained by averaging a large number of measurements at a pressure 1×10^{-3} mm Hg at different values of the parameter L. The light circles in the same figure show for comparison the electron temperatures obtained by probe measurements. It is seen that the results obtained by the different methods are close.

The data reported here were obtained for a weakly ionized low-temperature gas-discharge plasma. The method proposed by us, however, can be used also to determine the temperature in a hot fully ionized plasma.

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