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Influence of oscillations of plasma particles in a radial potential well on the beam-plasma interaction

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We investigate the influence of the limited radial dimensions of a plasma on beam-plasma interaction in the absence of a magnetic field. It is shown theoretically and experimentally that in the presence of a static transverse electric field, which causes transverse oscillations of the plasma, the limited radial dimensions of the system greatly influence the dispersion of the excited electronic oscillations if two conditions are satisfied: $k_z r_0 < 1$ and $\omega < v_T/r_0$, where k_z is the longitudinal wave number and r_0 is the inhomogeneity dimension (the radius of the plasma), v_T is the thermal velocity of the plasma electrons, and ω is the oscillation frequency. The observed effect, that the frequency is lower than the electron Langmuir frequency, is due to the decrease of the transverse conductivity of the plasma as a result of the indicated oscillations.

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One of the fundamental problems in beam-plasma interaction is the influence of the limited radial dimensions of the system on the dispersion of the excited oscillations. In the one-dimensional case, which is realized in a strong magnetic field ($\Omega_e \ll \Omega_c$, where Ω_e and Ω_c are the plasma and cyclotron frequencies of the electrons), it was shown theoretically and experimentally that the radial restriction leads to a substantial decrease of the frequencies and increments of the oscillations, and critical values of the parameters it stops the instability.[1–4] As to systems without a magnetic field, the situation remained unclear until the very latest time. On the one hand, a rigorous theoretical analysis with account of both the fact that the oscillations are not potential[5] and of the inhomogeneity of the system[3] has led inevitably to the conclusion that at all $k_z r_0$ (k_z is the longitudinal wave number and r_0 is the characteristic dimension of the inhomogeneity or the radius of the plasma) the increments and frequencies of the volume oscillations are equal to the corresponding values in an unbounded homogeneous system with particle concentrations close to the corresponding concentrations on the axis of the considered system.

On the other hand there exist experimental indications that in some conditions and in the absence of an external magnetic field the decrease of r_0 can lead to a decrease of the frequencies and increments of the excited volume oscillations.[6,7] This contradiction is resolved in the present paper, in which we construct a theory of beam-plasma interaction with account taken of the transverse static electric fields which confine the plasma particles within the beam, and show that under certain conditions these fields, in analogy with longitudinal magnetic field, lead to a substantial dependence of the dispersion of the excited oscillations on the characteristic dimension of the inhomogeneity (radius) of the system. The experimental data obtained in the present paper in an investigation of the collective interaction of a beam of positive ions with plasma electrons have confirmed the main conclusions of theory.

THEORY OF BEAM-PLASMA INTERACTION WITH ALLOWANCE FOR TRANSVERSE OSCILLATIONS OF THE PLASMA ELECTRONS IN AN ELECTROSTATIC WELL

We consider an axially symmetrical system in which a cold beam of charged particles passes with velocity V_0 along the Z axis through a plasma with electron temperature T_e. The potential of the plasma varies along the radius like

$$\varphi_0(r) = -T_e r^2/e r_0^2$$

and is the potential well for the electrons.

An investigation of the wave processes in the plasma will be carried out in a quasistatic approximation ($\mathbf{E} = -\nabla\varphi$) on the basis of linearized equations—the Vlasov equation for the electrons captured in the radial potential well, the hydrodynamic equations for the beam particles, and the Poisson equation. For simplicity we consider only axially symmetrical perturbations, and assume the upperturbed beam density n_{0b} to be homogeneous. Then the equation is describing the system can be written in the form

$$-i(\omega-k_zv_z)f+v_r\frac{\partial f}{\partial r}+\left(\frac{v_\varphi^2}{r}+\frac{e}{m}\frac{\partial\varphi_0}{\partial r}\right)\frac{\partial f}{\partial v_r}-\frac{v_rv_\varphi}{r}\frac{\partial f}{\partial v_\varphi}=-\frac{e}{m}\frac{\partial\varphi}{\partial r}\frac{\partial f_0}{\partial v_r}-ik_z\frac{e}{m}\varphi\frac{\partial f_0}{\partial v_z},\tag{1}$$

$$\left[1-\frac{\Omega_b^2}{(\omega-k_zV_0)^2}\right]\left(\frac{1}{r}\frac{\partial}{\partial r}r\frac{\partial\varphi}{\partial r}-k_z^2\varphi\right)=4\pi en_{0e}\int_{-\infty}^{\infty}dvf.\tag{2}$$

In Eqs. (1) and (2), we have carried out Fourier transformations in t and z, while f and φ are the Fourier components of the distribution function of the electrons and of the self-consistent potential, and Ω_b is the plasma frequency of the beam.

The stationary distribution function of the plasma electrons is chosen in the form of a Maxwell-Boltzmann function

$$f_0=\frac{1}{\pi^{3/2}v_T^3}\exp\left[-\frac{mv^2}{2T_e}+\frac{e\varphi_0(r)}{T_e}\right],\quad v_T=\left(\frac{2T}{m}\right)^{1/2}.\tag{3}$$

In this case the profile of the electron density takes the form

$$n_e(r)=n_{0e}\exp\left(-\frac{r^2}{r_0^2}\right).$$

(similar equations without the beam were used in Ref. 8 in an investigation of the anomalous skin effect in a plasma pinch.)

We use henceforth dimensionless variables and parameters:

$$\Phi=\frac{2}{v_T^2}\frac{e}{m}\varphi,\quad \mathbf{v}\to\frac{\mathbf{v}}{v_T},\quad r\to\frac{r}{r_0},\quad \omega\to\frac{\omega}{v_T}r_0,$$

$$k_z\to k_zr_0,\quad V_0\to\frac{V_0}{v_T}\quad \Omega_b\to\frac{\Omega_b}{v_T}r_0,\quad \Omega_e\to\frac{\Omega_e}{v_T}r_0.$$

We change over in (1) and (2) to new variables—the integrals of the unperturbed motion

$$rv_\varphi=\theta,\quad v_\varphi=\theta/r,$$
$$v_r^2+\theta^2/r^2-\Phi_0(r)=\varepsilon,\quad v_r=\pm(\varepsilon+\Phi_0(r)-\theta^2/r^2)^{1/2}.\tag{4}$$

Introducing the functions F^α defined by the formula $f=f_0(F^\alpha+\Phi)$ and integrating the Vlasov equation with respect to r, we obtain ($\alpha=\pm$ corresponds to $v_r\gtrless 0$)

$$F^\alpha(r)=F^\alpha(r_1)\exp(i\alpha\psi(r,r_1))+i\alpha\omega\int_{r_1}^{r}dr'\frac{\Phi(r')}{(\varepsilon+\Phi_0(r')-\theta^2/r'^2)^{1/2}}\exp(i\alpha\psi(r,r')),\tag{5}$$

$$\left[1-\frac{\Omega_b^2}{(\omega-k_zV_0)^2}\right]\left(\frac{1}{r}\frac{\partial}{\partial r}r\frac{\partial\Phi}{\partial r}-k_z^2\Phi\right)-2\Omega_e^2\exp(\Phi_0)\Phi$$
$$=\frac{\Omega_e^2}{\pi^{3/2}r}\int_{-\infty}^{\infty}dv_z\int_{-\infty}^{\infty}d\theta\int_{\theta^2/r^2-\Phi_0(r)}^{\infty}\frac{d\varepsilon}{(\varepsilon+\Phi_0(r)-\theta^2/r^2)^{1/2}}(F^++F^-)\exp(-v_z^2-\varepsilon),\tag{6}$$

where

$$\psi(r,r')=\pi(\omega-k_zv_z)T(r,r')=(\omega-k_zv_z)\int_{r'}^{r}\frac{dr''}{(\varepsilon-\theta^2/r''^2+\Phi_0(r''))^{1/2}}$$
$$=-\frac{1}{2}(\omega-k_zv_z)\left[\arcsin\frac{\varepsilon-2r^2}{(\varepsilon^2-4\theta^2)^{1/2}}-\arcsin\frac{\varepsilon-2r'^2}{(\varepsilon^2-4\theta^2)^{1/2}}\right].\tag{7}$$

The integration constants $F^\alpha(r_1)$ are obtained from the boundary conditions for the distribution function. As follows from (4), the velocity v_r vanishes at the points

$$r_{0,1}=(\varepsilon/2\pm(\varepsilon^2/4-\theta^2)^{1/2})^{1/2},\quad r_0<r_1.\tag{8}$$

At these points the particle velocity v_r reverses sign and the boundary conditions can be written in the form

$$F^+(r_0)=F^-(r_0),\quad F^+(r_1)=F^-(r_1).\tag{9}$$

From (5) we get with the aid of conditions (9)

$$F^\alpha(r_1)=-\frac{\omega}{\sin\psi(r_0,r_1)}\int_{r_1}^{r_0}dr'\frac{\Phi(r')}{(\varepsilon+\Phi_0(r')-\theta^2/r'^2)^{1/2}}\cos\psi(r_0,r').\tag{10}$$

It follows then from (5) and (10) that

$$F^++F^-=-\frac{\omega}{\sin\psi(r_0,r_1)}\int_{r_1}^{r_0}dr'\frac{\Phi(r')}{(\varepsilon+\Phi_0(r')-\theta^2/r'^2)^{1/2}}$$
$$\times[\cos\psi(r,r')\cos\psi(r_0,r_1)+\cos(\psi(r_0,r')-\psi(r,r_1))-\sin\pi(\omega-k_zv_z)|T(r,r')|\sin\psi(r_0,r_1)].\tag{11}$$

We expand the integrand in (11) in series, using the relations

$$\cos\psi(r,r')=\sum_{n=0}^{\infty}\frac{1}{\pi}\frac{2N}{N^2-n^2}\sin N\pi\cos n\pi\cos n\frac{\psi(r,r')}{\psi(r_1,r_0)}\pi,$$

$$\sin\pi(\omega-k_zv_z)|T(r,r')|=\sum_{n=0}^{\infty}\frac{1}{\pi}\frac{2N}{N^2-n^2}(1-\cos N\pi\cos n\pi)\cos n\frac{\psi(r,r')}{\psi(r_1,r_0)}\pi,$$

$$N=\pi^{-1}\psi(r_1,r_0).\tag{12}$$

As a result, after substituting (11) in (6) and integrating with respect to v_z, we obtain for the potential an equation that is exact within the framework of our model:

$$\left[1-\frac{\Omega_b^2}{(\omega-k_zV_0)^2}\right]\left(\frac{1}{r}\frac{\partial}{\partial r}r\frac{\partial\Phi}{\partial r}-k_z^2\Phi\right)-2\Omega_e^2\exp(\Phi_0(r))\Phi$$
$$=2\Omega_e^2\frac{\omega}{k_z}\frac{1}{\pi^2r}\int_0^\infty d\theta\int_{\theta^2/r^2-\Phi_0(r)}^{\infty}\frac{d\varepsilon\, e^{-\varepsilon}}{(\varepsilon+\Phi_0(r)-\theta^2/r^2)^{1/2}}\frac{1}{T(r_1,r_0)}$$
$$\times\sum_{n=-\infty}^{\infty}W\left(\frac{\omega}{k_z}-\frac{n}{k_zT(r_1,r_0)}\right)\int_{r_1}^{r_0}dr'\frac{\Phi(r')}{(\varepsilon+\Phi_0(r')-\theta^2/r'^2)^{1/2}}$$
$$\times\left[\cos n\pi\cos n\frac{T(r_0,r')-T(r,r_1)}{T(r_1,r_0)}\pi+\cos n\frac{T(r,r')}{T(r_1,r_0)}\pi\right],\tag{13}$$

where

$$W(z)=\frac{1}{\pi^{1/2}}\int_{-\infty}^{\infty}\frac{e^{-t^2}}{z-t}dt.$$

The analysis that follows will be made under the assumptions 1) $\omega/k_z\gg 1$ and 2) $\omega T(r_1,r_0)=\omega/2\ll 1$. The first condition means that we confine ourselves to the case of greatest practical interest, when the Landau damping on the electrons is negligible, and consequently the instability growth rate is largest. The second condition means that the frequency of the excited oscillations is lower than the frequency of the oscillations of the electrons in the static well. It is precisely under these conditions that we can expect a substantial decrease of the transverse conductivity of the plasma. It follows also from the foregoing assumptions that $k_z\ll 1$. Under the indicated conditions, the main contribution to (13) is made by the term of the sum with $n=0$.

We apply to (13) the Hankel transformation

$$\Phi(k_\perp)=\int_0^\infty rJ_0(k_\perp r)\Phi(r)dr,\quad \Phi(r)=\int_0^\infty k_\perp J_0(k_\perp r)\Phi(k_\perp)dk_\perp.$$

After simple transformations we obtain the equation

$$\left[1-\frac{\Omega_b^2}{(\omega-k_zV_0)^2}\right](k_\perp^2+k_z^2)\Phi(k_\perp)+\Omega_e^2\int_0^\infty k_\perp'\Phi(k_\perp')\times\exp\left(-\frac{k_\perp^2+k_\perp'^2}{4}\right)I_0\left(\frac{k_\perp k_\perp'}{2}\right)dk_\perp'=\Omega_e^2\frac{\omega}{k_z}W\left(\frac{\omega}{k_z}\right)\times\int_0^\infty k_\perp'\Phi(k_\perp')\exp\left(-\frac{k_\perp^2+k_\perp'^2}{4}\right)I_0^2\left(\frac{k_\perp k_\perp'}{4}\right)dk_\perp'. \qquad (14)$$

We find the condition for the solvability of this homogeneous equation. Integrating (11) with respect to $k_\perp$ with weight

$$\frac{k_\perp}{k_\perp^2+k_z^2}\exp\left(-\frac{k_\perp^2}{4}\right)$$

and expanding the modified Bessel function in power series, we arrive at the relation

$$\left[1-\frac{\Omega_b^2}{(\omega-k_zV_0)^2}\right]\int_0^\infty dk_\perp k_\perp\Phi(k_\perp)\exp\left(-\frac{k_\perp^2}{4}\right)=\int_0^\infty dk_\perp' k_\perp'\Phi(k_\perp')\exp\left(-\frac{k_\perp'^2}{4}\right)\Omega_e^2\sum_{m=0}^\infty\frac{1}{(m!)^2}\left[\frac{\omega}{k_z}W\left(\frac{\omega}{k_z}\right)\frac{(2m)!}{2^{2m}m!}-1\right]\times\int_0^\infty dk_\perp\frac{k_\perp\exp(-k_\perp^2/2)}{k_\perp^2+k_z^2}\left(\frac{k_\perp k_\perp'}{4}\right)^{2m}. \qquad (15)$$

Since $k_z\ll 1$, the main contribution to the sum over m will be given by the term with $m=0$. As a result we obtain the dispersion equation

$$1-\frac{\Omega_b^2}{(\omega-k_zV_0)^2}-\frac{\Omega_e^2}{\omega^2}U^2(k_z)=0, \qquad (16)$$

where

$$U^2(k_z)=k_z^2\int_0^\infty dk_\perp\frac{k_\perp\exp(-k_\perp^2/2)}{k_\perp^2+k_z^2}=-\frac{k_z^2}{2}\exp\left(\frac{k_z^2}{2}\right)\mathrm{Ei}\left(-\frac{k_z^2}{2}\right)\approx k_z^2|\ln k_z|.$$

$\mathrm{Ei}(z)$ is the integral exponential function. The dispersion equation (16) corresponds to density perturbations $n_e(r)\sim\exp(-r^2)$, which are the largest-scale perturbations possible in this system. Changing to dimensional quantities, we get

$$U^2(k_zr_0)\approx k_z^2r_0^2|\ln k_zr_0|.$$

As seen from the obtained dispersion equation, under our assumptions the frequency of the plasma electron oscillations excited by the beam is lower than the Langmuir frequency by a factor $1/U$. The dependence of the decrease factor U on k_zr_0, shown in Fig. 5 below, is of the same form as in the presence of a strong magnetic field.[1] In either case, the decrease of the oscillation frequencies due to the small transverse conductivity of the plasma. However, whereas in the latter case this is due to the magnetization of the electrons, in the system considered here it is due to rapid oscillations of the electrons in the static potential well.

It should be noted that a similar dispersion in a bounded plasma without a magnetic field at $k_zr_0\ll 1$ is possessed also by a surface wave. This wave, however, differs in principle from a volume wave and can be easily identified experimentally. As shown in Ref. 9, it can be excited in the investigated system described below only in the direction opposite to the direction of motion of the beam, has an enitrely different radial structure, and at $k_zr_0\gg 1$ its frequency is smaller by an approximate factor $2^{1/2}$ than the frequency of the volume oscillations excited by the beam. The decrease of its frequency at $k_zr_0<1$ is due to another physical mechanism—the emergence of the electric-field force lines to the outside of the plasma.

EXPERIMENTAL INVESTIGATION OF THE DISPERSION OF THE ELECTRON OSCILLATIONS EXCITED BY AN ION BEAM. DISCUSSION OF RESULTS

The experiments were performed on the setup whose diagram is shown in Fig. 1. A 20-mA beam of protons of energy 30 keV was extracted from the "duoplasmatron" *1* by the field of the extractor *2*, and was shaped by a magnetic lens *3* into a long weakly diverging beam *4* that passed through an ion guide of length up to 500 cm and diameter 35 cm to collector 5. The concentration of the beam ions was usually 10^7 cm^{-3}. The beam diameter could be varied by an iris diaphragm 6 located at the entrance of the beam into the ion guide. The plasma was produced by impact ionization between the beam ions and the air molecules. When the pressure was varied in the range 1×10^{-5}–10^{-3} Torr the ratio of the electron concentration to the beam-particle concentration varied in the range from 1 to 100. The radius of the plasma was approximately equal to the radius of the beam,[10] a fact that does not agree with the assumption in the theory that the beam is unbounded. However, since this paper deals in fact only with the dispersion of the plasma oscillations, this disparity does not play an essential role.

As shown previously,[10,11] in such a plasma there is always a static potential well for the electrons; the depth of the well is determined by the Coulomb conditions of the beam ions with these electrons. The experimentally obtained[10] radial distributions of the potential within the limits of the beam were close to $\varphi=\Delta\varphi(1-r^2/r_0^2)$, as is assumed in the theory.

Under the conditions of excitation of electron oscillations ($P=6\times10^{-5}$–6×10^{-4} Torr) the depth $e\Delta\varphi$ of the potential well in the beam does not depend strongly on the gas pressure or on the radial dimension of the beam,[10] and does not exceed 5 eV. When the plasma is produced only electrons with energy less than $e\Delta\varphi$ are accumulated in the system, while the faster electrons leave immediately and make no substantial contribution to the space charge. If the accumulated electrons are in thermal equilibrium, then in accordance with the foregoing their average thermal velocity remains

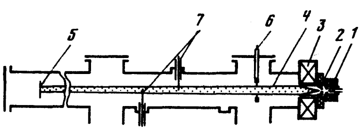

FIG. 1. Experimental setup: 1—ion source, 2—extractor, 3—magnetic lens, 4—beam, 5—target, 6—iris diaphragm, 7—probes.

practically unchanged under the experimental conditions, and does not exceed $v_T \sim (2e\Delta\varphi/m)^{1/2} \sim 1.4 \cdot 10^8$ cm/sec, so that $V_0 > v_T$.

The excited electron oscillations were registered with capacitive probes.[4] The same probes were used to measure the frequencies of the short-wave ion oscillations transverse to the beam ($k^2 d_e^2 > 1$, d_e is the electron Debye radius), whose maximum amplitudes corresponded to the plasma ion Langmuir frequency at the given distance along the radius.[12,13] This localization of the ion oscillations was used to calculate, from the measured ion Langmuir frequency, the concentration of n_{i0} the ions on the system axis. Using next the quasineutrality relation

$$n_{e0} \approx n_{0b} + n_{i0}, \tag{17}$$

we could calculate the axial concentration of the electrons. In those regimes when the ion oscillations were not excited, the density of the ions on the axis was calculated from a formula introduced from the condition for the balance of the production and free spreading of the ions under the influence of the electric field[10]:

$$n_{i0} = n_{0b} V_0 n_a \sigma_i r_0 / \bar{v}_i, \tag{18}$$

where σ_i is the cross section for the production of plasma waves, n_a is the concentration of air molecules and $v_i = (e\Delta\varphi/m_i)^{1/2}$ is the average velocity of the spreading ions.

As a confirmation of the applicability of Eq. (18), Fig. 2 shows plots of n_{i0} measured with the procedure described above against the radius of the system for two values of the pressure. It is seen that in accordance with expression (18) n_{i0} increases in proportion to r_0 and to the gas pressure, and the measured and calculated values of n_{i0} are equal if $\sigma_i \approx 1 \times 10^{-5}$ cm^2.

The measurements have shown that the amplitude of the oscillations of the potential is maximal on the beam axis, and the phase remains unchanged along the entire diameter of the system. Thus, the largest-scale mode of volume oscillations is excited in the plasma, and it was this mode which was investigated theoretically.

The dispersion characteristics of the oscillations were obtained by directly measuring the dependences of the frequencies of the excited oscillations on the radial dimension of the system. Figure 3 shows the indicated dependences, obtained for different gas pressures. The measured or calculated values of the plasma concentration were used to plot ω/Ω_e and $k_z r_0$ for all the curves of Fig. 3. In fact all the points fitted a single curve (Fig. 4), which is indeed the experimentally measured dependence of the coefficient U of the frequency decrease on $k_z r_0$.

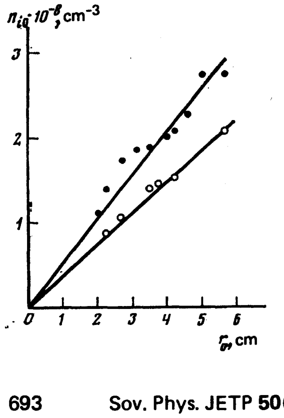

FIG. 2. Dependence of the plasma ion concentration on the system radius. ○ – $P = 1.2 \times 10^{-4}$ Torr, ● – $P = 1.7 \times 10^{-4}$ Torr.

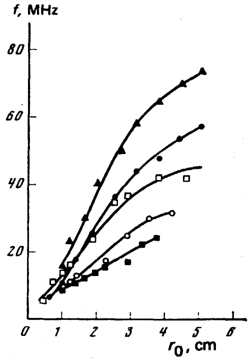

FIG. 3. Dependence of the frequency of the electron oscillations on the system radius: ■ – $P = 6 \cdot 10^{-5}$ Torr, ○ – $P = 8.10^{-5}$ Torr, □ – $P = 1.8 \cdot 10^{-4}$ Torr, ● – $P = 3.2 \cdot 10^{-4}$ Torr, ▲ – $P = 5.4 \cdot 10^{-4}$ Torr.

It follows from the theory that a substantial dependence of the dispersion of the oscillations on $k_z r_0$ should be observed when the following two conditions are satisfied:

a) $k_z r_0 < 1$, b) $\omega < v_T/r_0$ or $k_z r_0 < v_T/V_0$.

Inasmuch as in these experiments the beam velocity did not greatly exceed the thermal velocity of the electrons, the two conditions are actually equivalent. In accordance with the conclusion of the theory, a noticeable decrease of the ratio ω/Ω_e with decreasing $k_z r_0$ becomes observable at $k_z r_0 \sim 1$ (Fig. 4). At $k_z r_0 \ll 1$ the experimentally measured values of ω/Ω_e are close in order of magnitude to the theoretical values of U (Fig. 5). At $k_z r_0 \gg 1$, as expected, the experimentally measured values of ω/Ω_e do not depend in this case on the plasma radius. The very fact itself that ω/Ω_e does not reach unity is due primarily to the inhomogeneity of the plasma. The frequency of the electron plasma at oscillations corresponds in this case not through the maximum electron concentration, but to a certain value averaged over the radius.[3]

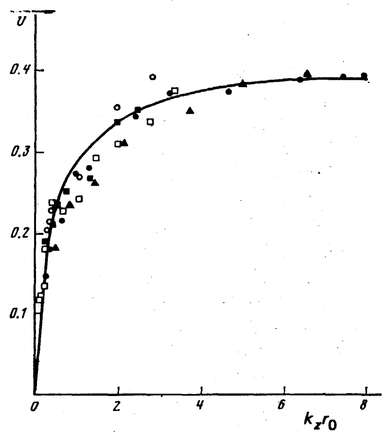

FIG. 4. Experimental dependence of the coefficient of frequency decrease on $k_z r_0$: ● – $P = 3.2 \cdot 10^{-4}$ Torr, ○ – $P = 8 \cdot 10^{-5}$ Torr, □ – $P = 1.8 \cdot 10^{-4}$ Torr, ▲ – $P = 5.4 \cdot 10^{-4}$ Torr, ■ – $P = 6.0 \cdot 10^{-5}$ Torr.

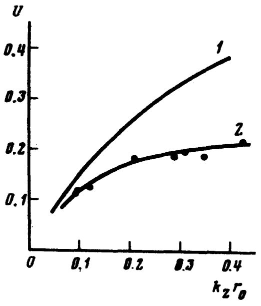

FIG. 5. Dependence of the frequency-decrease coefficient $U=\omega/\Omega_e$ on the system radius at small $k_z r_0$. 1—theory, 2—experiment.

The agreement between the experimental data and the theoretical calculation allows us to conclude that when plasma electron oscillations are excited by a beam of positive ions the electron oscillations in the potential well lead to a dependence of the dispersion of the oscillations on the radial dimension of the system.

The observed effect can play a substantial role also in other beam systems. Thus, in a plasma produced by a beam of negative ions or electrons, at sufficiently high pressure

$$n_a > n_{a\,cr} = \bar{v}_i / V_0 \sigma_i r_0, \qquad (19)$$

there is likewise produced a static potential well for the electrons. In the case of a negative-ion beam, the conditions for the onset of the frequency decrease will in fact be the same as in the case of the corresponding beam of positive ions.[7] On the other hand, in a plasma produced by an electron beam, we practically always have $v_T/V_0 \ll 1$, and therefore the fact that the radial dimensions are limited will come into play at sufficiently small $k_z r_0$ ($k_z r_0 < v_T/V_0 \ll 1$). For this reason the effect can be observed only at very low densities and small radii of the electron beam.

It is of interest to note that the radial limitation of the system can affect also the excitation of ion oscillations. This is possible in a plasma produced by a beam of negatively charged particles at $n_a < n_{a\,cr} = \bar{v}_i / v_0 \sigma_i r_0$, when there are practically no electrons in the system, and the ions oscillate in a potential well made up by the undercompensated space charge of the beam. This effect was possibly observed experimentally in a plasma produced by a beam of negative ions.[14] Actually, in this case the oscillation frequency became lower than the ion Langmuir frequency at $n_a < n_{a\,cr}$, when a potential well existed in the system for the plasma ions, and the condition $\omega < v_{Ti}/r_0$ was satisfied.

We have thus observed in the present study a new effect—the influence of the plasma-particle oscillations in radial static electric field on the dispersion of the oscillations in the absence of a magnetic field. We have shown that when the conditions $k_z r_0 < 1$ and $\omega < v_T/r_0$ are satisfied, these oscillations lead to a decrease of the transverse conductivity of the plasma, and this in turn causes a decrease of the frequency of the excited oscillations. The observed effect can play an important role not only in a beam plasma, but also in a plasma produced by an independent method, if the corresponding electric fields are present in the latter. Lowering the frequencies of the plasma oscillations causes a decrease in the growth rates, and this can lead, in systems of finite length, to stabilization of beam instability.

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