with a damping decrement

$$\gamma^{\pm}(\omega', k) = \gamma_e^{\pm} + \gamma_\mu^{\pm}, \qquad (5)$$

$$\gamma_e^{\pm} = \gamma_o^{\pm} \sum_{n,\sigma} (n \pm 1/2 + 1/2)\,[f^{\sigma}(E_{n,p^{\pm}(k,\Omega)} - \hbar\omega) - f^{\sigma}(E_{n,p^{\pm}(k,\Omega)})]^3,$$

$$\gamma_\mu^{\pm} = \gamma_o^{\pm} \sum_n (ak)^2 \left(\frac{\Omega_o}{\Omega}\right)^2 [f^{\pm}(E_{n,p^{\pm}(k,\Omega_o)} - \hbar\omega) - f^{\pm}(E_{n,p^{\pm}(k,\Omega)})],$$

$$\gamma_o^{\pm} = -\left(\frac{\partial G}{\partial\omega}\right)^{-1} e^2 (mc^2 a^4 k^3)^{-1}, \qquad a^2 = c\hbar/|e|H.$$

If we neglect the influence of the self-consistent fields on the US spectrum then we can put $\omega \approx v_0 k$ in (5). Equation (6) then describes the US damping decrement. It follows from (5) that when $\omega \ll \Omega$ or Ω_0, the absorption of sound exists only for waves with $k > k_{min}(\Omega_0) \sim \Omega_0\sqrt{m/2\zeta}$ (the first term in (5)) and $k > k_{min}(\Omega_0) \sim \Omega_0\sqrt{m/2\zeta}$ (second term in (5)). $\gamma_\mu^{\pm}$ is due to the interaction of the US with the electron spins responsible for the Pauli paramagnetism. The absorption maxima $\gamma_e^{\pm}$ and $\gamma_\mu^{\pm}$ are shifted because of the difference between Ω_0 and Ω.

If $\hbar\omega' \ll T$, the difference between the Fermi functions in (5) can be replaced by the derivative, and it is then obvious that the positions of the maxima of the US absorption oscillations are determined by Eqs. (1) and (2). The conditions necessary for the experimental observation of the oscillations of the paramagnetic absorption of ultrasound is given by the equations

$$\hbar\Omega,\ \hbar\Omega_o \gg T,\ \hbar/\tau,$$

$$|\Omega - \Omega_o| > \omega \gg (mv_o^2/\hbar\tau)^{1/2}.$$

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PLASMA DIFFUSION IN STELLARATORS

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As is well known, in many stellarators anomalous leakage of plasma is observed in a wide range of plasma-parameter variation [1 - 4]. The cause of the anomalous diffusion is not clear at present. On the one hand, plasma in toroidal systems is subject to a wide class of dirft instabilities (cf., e.g., [5] and the literature cited there), and some experiments [4, 7] the observed leakage can be related to oscillations in the drift-frequency range. On the other hand, even classical diffusion, with allowance for the trapped and localized particles, can reach appreciable values in asymmetrical toroidal systems [8 - 10]. However, in the most typical region of plasma parameters, where the electron collision frequency is high, and the ion frequency is low compared with the drift frequency, classical diffusion cannot explain the observed leakage, inasmuch as in this case the electron diffusion coefficient is small (the latter circumstance was noted in [10]).

But it is precisely in this region, as will be shown below, that a sufficiently dangerous instability of the drift type can develop, when account is taken of the finite orbits [5, 6]; this instability suffices for an anomalous plasma leakage on the order of the Bohm leakage, and naturally transports both ions and electrons to an equal degree.

Since the curvature of the field is immaterial in the following, we consider as the stellarator model a straight system with cylindrical magnetic surfaces and with a field

$$B_z = B\,[1 + \epsilon \cos\,(l\nu - k_0 z)\,], \tag{1}$$

where ε is the depth of modulation of the magnetic field by the helical winding, ℓ is the multiplicity of the helical winding, and k_0 is the reciprocal of its period on the z axis.

Let us consider small plasma oscillations which are strongly elongated along the magnetic field, so that the corresponding longitudinal wave number is $k_z \ll k_0$. We assume that the ion collision frequency is small compared with the oscillation frequency ω, and consequently we can use the collision-free kinetic equation for their description. We assume furthat that $\omega \ll \omega_b \sim k_0 v_i = k_0\sqrt{T/m_i}$, where v_i is the thermal velocity of the ions and ω_b is the frequency of ion oscillations in the periodic magnetic field (1). The ions that pass through equalize well the potential of the electric field along z over a length $\sim x_0^{-1}$, so that the potential ϕ can be set equal to

$$\phi = \phi\, e^{-i\omega t + im\nu + ik_z z}\quad,$$

where ϕ does not depend on z (i.e., we can neglect small oscillations of $\tilde{\phi}(k_0 z)$ with the period of the magnetic-field modulation). In addition, we can average in the kinetic equation for the ions over the fast oscillations ω_b, so that the transverse motion of the ions can be described in terms of the conservation of the longitudinal adiabatic invariant

$$I = \sqrt{2/m_i}\,\int (E - e\phi - \mu B)^{1/2}\, d\ell, \tag{2}$$

where E is the total energy of the ion and μ is the transverse adiabatic invariant.

The appropriately linearized kinetic equations for the ions, assuming $k_z v_i \ll 1$, takes the form

$$-i\,\omega f' + i v_0\,\frac{m}{2}\,f' + v_r\,\frac{df_0}{dr} = 0, \tag{3}$$

where $v_r(E, \mu)$ and $v_\nu(E, \mu)$ are the components of the drift velocity averaged over the longitudinal oscillations, and equal

$$v_r = \frac{c}{eB}\frac{dI}{r\,d\nu}\left(\frac{dI}{dE}\right)^{-1}, \quad v_\nu = -\frac{c}{eB}\frac{dI}{dr}\left(\frac{dI}{dE}\right)^{-1}. \tag{4}$$

Assuming that in the equilibrium state the transverse electric field is zero, we obtain from (3), with allowance for (2) and (4), the perturbation of the ion density:

$$n_i' = \frac{en_o}{T}\phi < \frac{\omega_*}{\omega - \omega_m} > \simeq \frac{en_o}{T}\phi\frac{\omega_*}{\omega}\left(1 + \frac{<\omega_m>}{\omega} + \frac{<\omega_m^2>}{\omega^2}\right), \tag{5}$$

where

$$\omega_* = -\frac{m}{r}\frac{cT}{eBn_o}\frac{dn_o}{dr}$$

is the so-called drift frequency, $\omega_m = (m/r)v_\nu$ - is the magnetic-drift frequency, and the angle brackets denote averaging over the Maxwellian distribution function. The averaging over ω_m/ω in (5) is based on the assumption that $\omega >> \omega_m$. The magnitude and sign of $<\omega_m>$ are determined by the magnetic-field characteristics averaged over the length. In order of magnitude we have $<\omega_m> \sim \varepsilon^2\omega_*$, and the sign of $<\omega_m>$ coincides with the sign of ω_*. In the case of magnetic well, max $\oint d\ell/B$, and is opposite to it in the case of min $\oint d\ell/B$. The main contribution to the quantity $<\omega_m^2>$ is made by the trapped ions, the fraction of which is $\sim \sqrt{\varepsilon}$, so that in order of magnitude we have $<\omega_m^2> \sim \varepsilon^{5/2}\omega_*^2$. Assuming that the electron collision frequency is $\nu_e >> \omega_*$ and T = const, we can use for the perturbation of the electron density the well known hydrodynamic expression [11]

$$n_e' = \frac{en_o}{T}\phi\left\{1 + \frac{\omega_* - \omega}{\omega + iDk_z^2}\right\}, \tag{6}$$

where $D = T\nu_e/m_e$ is the electron diffusion coefficient.

Equating (5) and (6) and assuming that $Dk_z^2 << \omega << \omega_*$, we obtain the frequency of the oscillations under consideration

$$\omega = -\frac{<\omega_m^2>}{<\omega_m> + iDk_z^2} \tag{7}$$

We see therefore that the maximum increment $\gamma \sim \sqrt{\varepsilon}\omega_*$ is reached when $Dk_z^2 \sim <\omega_m> \sim \varepsilon^2\omega_*$. It is very important that in this case $\gamma = \mathrm{Im}\,\omega$ is completely insensitive to the sign of $<\omega_m>$, i.e, to the presence or absence of min B.

The instability considered here is analogous to the drift-dissipative instability [12], but differs from in the less stringent requirements with respect to the shear θ, which limits from below the possible values of $k_z > \theta m/r$. Since the instability increment is of the order of the frequency, we can expect development of strong density fluctuations and macroscopic plasma leakage with an effective diffusion coefficient $D_\perp \sim \sqrt{\varepsilon}\, cT/eB$. Incidently, it is not excluded that the instability sets in at a somewhat lower level, since even when $\phi \sim \varepsilon T/e$ the electric field begins to influence the ion drift trajectories I = const. The diffusion coefficient $D_\perp$ is of the order of magnitude of the Bohm coefficient and greatly exceeds the maximum attainable classical coefficient of diffusion of ions on localized particles [10], $D_c = \sim \varepsilon_t^2 \varepsilon^{-1/2} cT/eB$, where $\varepsilon_t = a/R$ is the ratio of the minor radius of the plasma to the radius of curvature (from the condition that there be no disruption of the magnetic surfaces by the curvature it follows that $\varepsilon_t < \varepsilon^2$). In addition, unlike the classical diffusion on localized particles, the turbulent diffusion considered here involves all the ions to an equal degree, and it is therefore insensitive to the formation of a plateau on the distribution function in the region of the localized ions, which leads to a lowering of D_c when the collision frequency is decreased [9].

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DETERMINATION OF THE LEVEL WIDTHS OF GAS MOLECULES BY THE PHOTON-ECHO METHOD

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In analogy with spin echo [1], the photon-echo phenomenon is used successfully to determine the relaxation times in a solid [2] and the collision widths of the levels of gas molecules [3, 4]. The photon-echo method is useful also for the determination of the g-factors of levels [5] and other characteristics of gas molecules [6, 7]. So far, however, the photon-echo method was used only to measure the sum of the widths of the upper and lower resonance levels by exciting the gas with two light pulses.

We demonstrate in this paper the possibility of measuring the width of each resonance level separately by the photo echo method. To this end, the gas medium must be excited by three light pulses and the photon echo produced by the second and third pulses must be investigated. The level widths are determined from the attenuation of the intensity of this photon echo as a function of the time interval between the transmitted pulses. The greatest effect is reached when the first, second, and third excited pulses are respectively 180, 90, and 180 degree pulses, and all move in the same direction.

For concreteness, we consider the photon echo in a gas for an atomic transition with change of total angular momentum $1/2 \to 1/2$. The basic equations are chosen to be the