Hahn for an ideal, $\theta = 2\pi$, sech pulse travelling in a medium with infinite T_2'. Since $\alpha L \simeq 10$, $t_p \simeq 16$ nsec, one would expect $t_d(2\pi) \simeq 80$ nsec; but since $T_2' < 80$ nsec, relaxation processes cannot be ignored. The latter are complicated in ruby, but preliminary computer calculations using $T_2' < 80$ nsec yield $t_d(2\pi) \sim 20$ nsec as observed [F. A. Hopf (private communication)]. Similar but lesser divergences are observed in SF_6 (Ref. 6, $\alpha L = 4.8$, $t_p = 180$ nsec) where $t_d(2\pi) \simeq 230$, not 432 nsec. Although the pioneering experiments of McCall and Hahn were qualitative, they obtained $t_d(2\pi) \sim 25$ nsec, not >60 nsec for $\alpha L > 24$, $t_p = 5-10$ nsec.

[13]Since the input beam has a roughly Gaussian cross section, the center corresponds to a higher intensity and θ than the periphery. Thus we detected an *average* over the beam cross section; this corresponds to the methods of Ref. 5 with which we compare our results. This is not too serious for low values of θ where only the intense center portion will escape severe absorption. However, unless an aperture is placed (or imaged) at the output face of the attenuator, it is difficult to observe pulse breakup and transmission oscillations in the full transparency region (see Ref. 6). This procedure enhances the plane-wave nature of the effective beam; it should not be confused with the use of apertures in front of the photomultiplier to increase the reproducibility of the detection system. By the time the output beam travels 1.5 m to the photomultiplier, its light is thoroughly mixed.

[14]E. Courtens, in *Proceedings of the Chania International Conference on the Optical Properties of Solids*, 1969, edited by E. D. Haidemenakis (Gordon and Breach, New York, to be published).

[15]This is equivalent to measuring the vertical scale in units such that the maximum transmission corresponds to unity. Since the curves are plotted on a log/log scale, their shape is unaffected by the multiplicative factors involved.

[16]I. M. Asher and M. O. Scully, Opt. Commun. 3, 395 (1971).

[17]D. E. McCumber and M. D. Sturge, J. Appl. Phys. 34, 634 (1963).

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Parametric Excitation of Transverse and Longitudinal Waves near the Plasma Frequency*

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Intense laser light with frequency ω_0 near the plasma frequency excites transverse waves with real frequency components 0 and $\pm 2\omega_0$ and longitudinal plasma oscillations at $\pm\omega_0$. The resonant coupling is due to the $\vec{v} \times \vec{B}$ forces and relativistic mass corrections. The maximum growth rate is $\omega_0 v_0^2/16c^2$, where v_0 is the maximum electron velocity in the incident laser light.

For a hot plasma interacting with intense laser light, plasma instabilities can become the dominant mechanism that couples the laser energy to the plasma. We will derive a new instability with the incident light frequency ω_0 near the plasma frequency ω_p. For laser intensities not much beyond the present state of the art, this new instability is competitive with the so-called modified two-stream electrostatic instability.[1] The incident wave couples the perturbed-electron-plasma oscillation and the perturbed transverse electromagnetic wave. The transverse wave has two real frequency components at zero and $\pm 2\omega_0$. These two waves can be thought of as the analogs of the Stokes and anti-Stokes lines of stimulated Raman scattering.

We will assume that the wavelength of the excited wave is much less than that of the incident wave, $k \gg k_0$. We justify this assumption later in the article. The coordinate system has the incident electric field and the perturbed transverse electric field in the x direction, the perturbed magnetic field in the y direction, and the perturbed longitudinal electric field and wave vector in the z direction. Only electron density perturbations are considered; the high-frequency density fluctuations have negligible effect on the heavy ions.

Using the transverse gauge $\nabla \cdot \vec{A} = 0$, with A_1 the perturbed transverse vector potential, we easily obtain

$$\frac{\partial}{\partial t}\left(\gamma \frac{\partial n_1}{\partial t}\right) + \omega_{pk}^2 n_1 = -\frac{e^2 k^2 E_0 n_0}{m^2 \omega_0 c \gamma} A_1 \sin\omega_0 t, \quad (1)$$

$$\frac{\partial^2 A_1}{\partial t^2} + k^2 c^2 A_1 + \frac{\omega_p^2}{\gamma} A_1 = -\frac{4\pi e^2 c E_0}{m \omega_0 \gamma} n_1 \sin\omega_0 t \ . \quad (2)$$

Because the growth rate will be of the order of v_0^2/c^2, where v_0 is the maximum velocity in the incident electric field $E_0 \sin\omega_0 t$, we have included the relativistic correction factor γ. With $k^2 K T_e/m$ of order $\omega_0^2 v_0^2/c^2$, we only need the rest-frame pressure, so that $\omega_{pk}^2 = \omega_p^2 + 3k^2 K T_e/m$. n_1 is the electron density in the lab frame of reference.

We now expand A_1 and n_1 in harmonics[2] of ω_0:

$$n_1 = (n_1^{(+1)} e^{i\omega_0 t} + n_1^{(-1)} e^{-i\omega_0 t}) e^{i\omega t} + \cdots ,$$

$$A_1 = (A_1^{(0)} + A_1^{(+2)} e^{2i\omega_0 t} + A_1^{(-2)} e^{-2\omega_0 t}) e^{i\omega t} + \cdots .$$

We have kept only the dominant terms driven by the resonance $\omega_0 \sim \omega_p$. On the left-hand side of Eq. (2) we keep only the $k^2c^2A_1$ term, with the approximation $kc \gg \omega_p$. The solution for A_1 is then substituted into Eq. (1). The γ factors on the right-hand side of Eq. (1) can be neglected to the order of interest, and expanded to order v_0^2/c^2 on the left-hand side. Thus the electron density fluctuations at $\pm\omega_0$ are coupled by the incident E_0 through relativistic mass corrections and through the $\vec{v}\times\vec{B}$ force. The dispersion relation becomes

$$\omega^2 = (\omega_p^2 - \omega_0^2 + 3k^2KT/m - \tfrac{3}{4}v_0^2\omega_0^2/c^2)^2/4\omega_0^2 - \omega_0^2 v_0^4/256c^4 . \quad (3)$$

The maximum growth rate is[3,4]

$$\omega_I = \tfrac{1}{16}\ (v_0^2/c^2)\ \omega_0 ,$$

when the bracket in Eq. (3) equals zero. As ω_{pe} approaches ω_0, the most unstable wave number satisfies $k^2c^2 = \omega_0^2 v_0^2/4v_t^2 \gg \omega_0^2$. Since the incident light satisfies the approximate equation $k_0^2c^2 = \omega_0^2 - \omega_p^2$, we were justified in ignoring the wavelength dependence of the incident light. [The modified two-stream instability has a maximum growth rate given by $\omega_I^3 = 0.17\ \omega_{pi}^3 (M/m)^{1/2}$.]

Energy is being deposited in both the transverse electromagnetic wave E_t and in the longitudinal electrostatic wave E_l. The ratio of the two can also be found from Eqs. (1) and (2):

$$E_l/E_t \cong (c/v_0)(kc/\omega_0) \gg 1.$$

Most of the energy will be deposited in the longitudinal plasma oscillations. The wave vector $\vec{k}_0$ of the laser light can be paralled to $\vec{k}$ or parallel to $\vec{B}_t$. If $\vec{k}$ is perpendicular to the plasma density gradient, then the effects of the plasma density gradient should be minimized.

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[1]K. Nishikawa, J. Phys. Soc. Japan 24, 4 (1968).
[2]V. P. Silin, Zh. Eksperim. i Teor. Fiz. 48, 1679 (1965) [Sov. Phys. JETP 21, 1127 (1965)].
[3]J. Sanmartin, Phys. Fluids 13, 1533 (1970).
[4]N. L. Tsintsadze, Zh. Eksperim. i Teor. Fiz. 59, 1251 (1971) [Sov. Phys. JETP 32, 684 (1971)].

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Motion of Ions in Helium II

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The general motion of ions in He II is investigated assuming that nucleation of and escape from vortex rings is a random thermally activated process. The mean drift velocity v_d is calculated as a function of temperature, applied field, and ionic species. It is shown that low-field data up to and just beyond the giant discontinuity can be explained, provided careful attention is paid to the friction forces on small rings, by assuming that v_d is the equilibrium drift velocity. The transition between bare-ion and vortex-ring behavior is discussed in some detail. At higher fields one must take into account vortex-ring dynamics and the possibility of escapes. In general v_d is larger than the equilibrium velocity and, for very large fields, increases with field. Predictions of the theory are compared with experimental drift-velocity data. Also considered are the characteristics of ion currents in nonuniform fields. In particular, predictions are made for the "persistence current" observed when ions propagate first through a region of constant field, then through a region of zero or retarding field.

I. INTRODUCTION

The diversity and novelty of problems associated with the motion of ions in superfluid helium has made this subject a popular area of investigation for both theorists and experimenters. The diversity arises because the ion has separate and quite different interactions with the normal and the superfluid components. The relevant physical processes have been mostly identified, however, and the main features of ion motion can now be understood. It is the purpose of this paper to discuss the interplay of these processes and to present quantitative calculations of the average ion drift velocity and persistence currents.

The main physical effects governing the motion of an ion under the influence of an electric field are the viscous force on a bare ion, the possibility the ion will nucleate and be captured by a quantized vortex ring, the electric and viscous forces on the