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Observation of Transverse Zero Sound in Normal ^{3}He†

Pat R. Roach
Argonne National Laboratory, Argonne, Illinois 60439

and

J. B. Ketterson
Department of Physics, Northwestern University, Evanston, Illinois 60201, and Argonne National Laboratory, Argonne, Illinois 60439

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The direct transmission of transverse zero sound between two AC-cut quartz transducers is observed in the vicinity of 3 mK; at higher temperatures the transition to classical viscous shear-wave behavior is observed through the effect on the damping of a single transducer.

Following the development of his theory of a Fermi liquid in 1957,[1] Landau predicted the existence of two collective excitations in normal ^{3}He at very low temperatures[2]; these excitations are referred to as longitudinal (symmetric) and transverse (asymmetric) zero sound. The onset of the longitudinal zero-sound mode was first observed by Keen and co-workers[3] while the complete temperature dependence of the velocity and attenuation was first studied by Abel, Anderson, and Wheatley[4]; the transverse mode has not been observed previously.

A requirement for the propagation of zero sound is that its velocity, V, be greater than the Fermi velocity, V_F. The strength of the Landau-quasiparticle interaction parameters, F_i, determines V and for the transverse mode the condition $V > V_F$ requires $F_1 > 6$ for the case $F_i = 0$, $i \geq 2$. Experimental data indicate $F_1 > 6$ for all pressures, but reliable data on higher F_i are not available.

The zero (or collisionless) sound regime is characterized by the condition $\omega\tau \gg 1$, where ω is the angular sound frequency and τ is a quasiparticle collision time; in the degenerate regime $\tau \propto T^{-2}$. As the temperature is increased, a hydrodynamic regime is entered where $\omega\tau \ll 1$. The modification of the transverse mode at high temperatures is radically different from the longi-

tudinal case: The transverse wave goes over into the classical damped viscous shear mode with a complex propagation constant, $K \equiv k + i\alpha$, given by

$$K = (1+i)(\omega\rho/2\eta)^{1/2}; \quad (1)$$

here ρ is the density and η is the viscosity. This mode is diffusive and does not propagate in the usual sense.

The attenuation of transverse zero sound was first calculated by Corruccini, Clarke, Mermin, and Wilkins.[7] The full dispersion relation for all values of $\omega\tau$ was later calculated by Lea, Birks, Lee, and Dobbs.[8] Both sets of investigators considered the case $F_2 = 0$. We have extended the calculations to the case $F_2 \neq 0$ and find for the dispersion relation

$$(\xi^2 - 1)w = \frac{F_1 - 6 - 9\beta + 12\xi^2 a(1+\beta)}{3F_1 - 9\beta + 36\xi^2 a(1+\beta)}, \quad (2)$$

where $a = \frac{1}{4}F_2/(1+\frac{1}{5}F_2)$, $\xi = (i\omega\tau - 1)/i\tau K V_F$, $\beta = (i\omega\tau - 1)^{-1}$, and $w = \frac{1}{2}\xi \ln[(\xi+1)/(\xi-1)] - 1$. Recently Bolton[9] has shown that, for the case where only F_1 is retained, the results of Lea *et al.* may be closely approximated by use of a simplified transport theory. He finds for the dispersion relation

$$K^2 l^2 = \frac{-\omega\tau(1 - i\omega\tau)^2}{\frac{1}{5}i m^*/m + \omega\tau}, \quad (3)$$

where m^* is the effective mass, $l = V_F\tau$, $m^*/m = 1 + \frac{1}{3}F_1$, $V_F = p_F/m^*$, and $p_F = \hbar(3\pi^2 n)^{1/3}$, with n the number density of ^{3}He atoms. This expression yields the classical expression for K at high temperatures [$\eta = \frac{1}{5}(m^*/m)\rho V_F{}^2\tau$], and $\omega/K = V_F$ and $\alpha = (V_F\tau)^{-1}(1 - m^*/10m)$ at low temperatures; note that the zero-sound attenuation is proportional to T^2.

Bolton has also taken into account the effect of slip between the liquid ^{3}He and a transversely oscillating boundary with maximum velocity u_0 and has calculated the work done per unit area per cycle, W, and finds

$$W = \frac{\pi u_0{}^2}{\omega}\frac{\eta}{l}\frac{\mathrm{Im}Kl + (2/\theta)|Kl + \omega\tau|^2}{1 + (4/\theta)\,\mathrm{Im}Kl + (2/\theta)^2|Kl + \omega\tau|^2}; \quad (4)$$

for diffuse (specular) reflection of the ^{3}He quasiparticles striking the surface, θ has the value 1 (0); no damping occurs in the limit $\theta = 0$.

A pair of AC-cut quartz transducers was used to generate and detect the transverse waves. Calculations using Eqs. (2) or (3) show that the attenuation is quite high and thus a very short path length is required. The path length employed in our experiments was 0.0025 cm, which was provided by a fine platinum wire separating the two transducers.[10] Some bowing of the transducer is possible as well as a compression of the wire spacer; both of these effects can lead to some uncertainty in the path length. Because of the extremely short wavelength of the transverse mode (5×10^{-4} cm at 12 MHz) nonparallelism between the transducers can reduce the amplitude of the signal. By viewing with coherent light two optical flats separated by a wire spacer, it was determined that parallelism of order a fringe could be achieved. With the path length employed there is no advantage in using a pulse technique since a time separation between the feedthrough and the signal is not possible; thus a continuous, amplitude-modulated, rf signal was applied to the first transducer. The output of the second transducer was applied to an rf amplifier detector and then to a lock-in amplifier which was, in turn, driven by the modulation frequency, f_m. The modulation frequency of 20 Hz was such that $f_m \ll 1/\tau_R$ (where τ_R is the ringing time of the transducers) and thus the transducer can follow the rf envelope. A certain amount of electromagnetic coupling or feedthrough exists between the receiving and transmitting transducers; for 12 MHz the feedthrough was the same order of magnitude as the zero-sound signal at 2.5 mK. The feedthrough was nulled by passing amplitude-modulated rf from the oscillator through a phase shifter and attenuator and adding it to the receiver input. Some temperature dependence of the feedthrough is expected because of the temperature dependence of the electrical impedance of the transducers (which in turn arises from the ^{3}He acoustic impedance); this affects the magnitude of the rf signal in the cell and thus the feedthrough. However, this small effect on the feedthrough should be completely temperature independent below ~5 mK according to our acoustic impedance measurements. Some spurious signal from a small amount of longitudinal zero sound could also exist. For longitudinal sound to exhibit an appreciable attentuation change in the range of 2.5–4 mK, however, it must exist as a standing wave of very high standing-wave ratio. Otherwise, in transiting the cell only once, for example, the wave would be attenuated by an amount that would change by less than 0.1 dB in going from 2.5 to 4 mK. A high standing-wave ratio is very unlikely in our cell since it would require a much higher degree of parallelism of our transducers than

would appear reasonable. Even if the parallelism were adequate, it would not be possible for the transducer spacing to satisfy the very sharp standing-wave resonance condition at all the different pressures (and sound velocities) at which measurements were made. Thus the signals arising from the temperature-dependent loading effect and a possible longitudinal component are expected to have little temperature dependence below about 5 mK; therefore, we nulled the detected signal near this temperature and the signal which grows rapidly out of the noise below about 4 mK can only arise from transverse zero-sound transmission. In order to measure the attenuation a second attenuator–phase-shifter combination (in parallel with the previously mentioned one) was used to null the transverse zero-sound signal. In order to increase the signal-to-noise ratio, an unmodulated rf bias of appropriate phase and amplitude was also added to the receiver input.

We now discuss the experiments involving the damping of a transducer following the removal of its rf excitation energy. If one assumes that the only mechanism for energy loss from the transducer is the coupling of energy into the liquid which loads the two faces, the time dependence of the detected signal, ψ, is given by

$$\psi = \mathrm{Re}\psi_0 \exp[(-t/\tau_R) - i(\omega + \Delta\omega)t], \tag{5}$$

where $1/\tau_R = (2R_H/\pi R_Q)\omega_0$ and $\Delta\omega = (2X_H/\pi R_Q)\omega_0$; here $f_0 = \omega_0/2\pi$ is the fundamental frequency of the transducer, $\omega = (2n+1)\omega_0$, R_H and X_H are the real and imaginary parts of the acoustic impedance, Z_H, of the ^{3}He, and R_Q is the acoustic impedance of the quartz (assumed real). Equation (5) neglects the effects of nonparallelism of the transducers and the extraction of energy for detection; it further assumes $Z_H/Z_Q \ll 1$.

Classically (and neglecting slip), the acoustic impedance of a material is defined by ρV, where ρ is the density and V is the complex acoustic velocity, $V = \omega/K$. Thus

$$\frac{R}{\rho} = \frac{\omega k}{k^2+\alpha^2}, \quad \frac{X}{\rho} = \frac{\omega\alpha}{k^2+\alpha^2}, \tag{6}$$

where k and α are the real and imaginary parts of K. Based on simple arguments relating the work done on the liquid to the decay time, τ_R, used to determine R_H, we obtain $W\omega/\pi u_0^2 = R_H$. Using Eq. (4) we find $R_H = \rho\,\mathrm{Re}V$ in the limit $\omega\tau \ll 1$, $\theta = 1$, or $R_H = (\omega\rho\eta/2)^{1/2}$; this is consistent with the classical interpretation of R_H. However at low temperatures ($\omega\tau \gg 1$) we find $R_H = \rho V_F m^*/10m$; thus an additional factor $m^*/10m$ multiplies the classical result.

The ringing experiments were performed as follows: An rf pulse whose width was long relative to τ_R was applied to the transducer. The time dependence of the amplitude and phase of the ringing signal was then studied by adding a second comparison or reference signal of variable phase and amplitude to the receiver input; by setting the reference signal such that its amplitude was identical but its phase opposite to that of the signal at some particular instant in the decay, the time dependence of the amplitude and phase of the ringing signal could be determined. The existence of spurious effects on the amplitude and phase of the ringing signal made it necessary to normalize the data. The decay was measured at high temperatures where the viscous damping is negligible and the spurious decay time, τ_S, determined. At lower temperatures, where ^{3}He loading effects appear, the total decay time, τ_T, was measured and τ_R was deduced from $1/\tau_T = 1/\tau_R + 1/\tau_S$.

The cryogenic techniques and the sonic cell employed are essentially identical to those used in our earlier studies of sound propagation in the superfluid phases of ^{3}He.[11]

Figure 1 shows the temperature dependence of the attenuation of zero sound as determined in the transmission experiments for pressures of 2.0, 8.0, 15.0, 23.0, and 28.9 bar. As in all

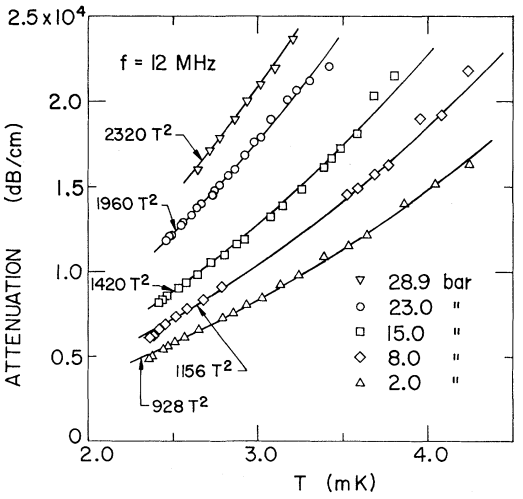

FIG. 1. The temperature dependence of the attenuation of transverse zero sound for pressures of 2.0, 8.0, 15.0, 23.0, and 28.9 bar. The fitted coefficient of the T^2 temperature dependence is shown next to the data at each pressure.

acoustic experiments where only the relative attenuation can be determined, a normalization was required. The normalization was acomplished by requiring a T^2 dependence for the attenuation. A smoothly varying behavior of the extrapolated zero-temperature normalization constant with pressure provided an overall consistency check on the data.

The results of Fig. 1 may be compared with the attenuation coefficients resulting from Eq. (2) when various values of F_2 are put in along with known values[12] for F_1, V_F, and τ. In the limit $\omega\tau \gg 1$, Eq. (2) yields attenuation coefficients corresponding to our measured results when values of F_2 from 1.5 (at 2 bar) to -1.0 (at 28.9 bar) are used. (The opposite pressure dependence of F_2 resulted from calculations of Ostgaard.[13]) Considerable caution is required before accepting these values of F_2 at face value. By assuming, when we determined the attenuation coefficients, that the sound path length is given by the diameter of our wire spacer, we could be introducing an error of 1 to 2 units in the magnitude of F_2. This path-length uncertainty would have little effect on the pressure dependence of F_2, however. Of the remaining parameters used in arriving at F_2, both the values of τ and the values of F_1 are somewhat uncertain. At 29 bar it would only require a 10% decrease in the value of τ used along with a 15% decrease in the value of F_1 used in order to account for our 29-bar attenuation data using a value $F_2 = 0$. Such an adjustment to τ is probably within the uncertainty limits of the viscosity data from which it is obtained[14,15] and the adjustment to F_1 is suggested by recent heat capacity measurements of Halperin *et al.*[16]

Figure 2(a) shows, for a pressure of 23 atm, the values of R_H/ρ as a function of temperature for frequencies of 36, 60, and 108 MHz determined from the measurements of $1/\tau_R$; measurements were also made at 12 MHz. Figure 2(b) shows the corresponding values of X_H/ρ determined from the frequency shifts, $\Delta\omega$. The values of X_H/ρ have been normalized such that the results for the three frequencies all extrapolate to zero at high temperature since this is the theoretically predicted behavior. In Fig. 2(a) we have plotted the theoretical values of R_H/ρ using Eq. (4); here we used an effective mass $m^* = 5.31m$ (corresponding to $F_1 = 12.94$) and $\tau(T)$ was deduced by interpolating between the viscosity data of Bertinat *et al.*[14] at low pressure and of Alvesalo *et al.*[15] at the melting pressure. The behavior of X_H was not reported by Bolton. The

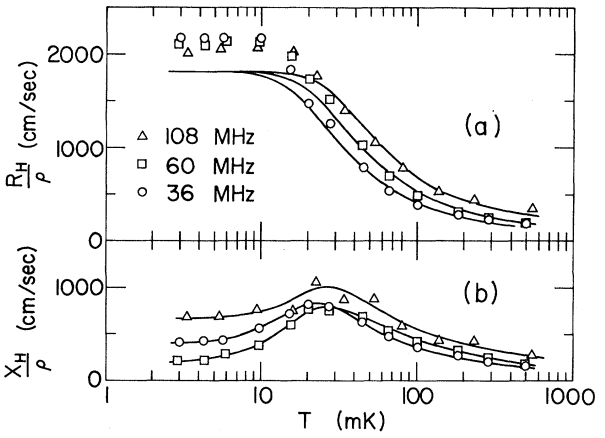

FIG. 2. The temperature dependence of the measured values of (a) R_H/ρ and (b) X_H/ρ for a pressure of 23.0 bar. The solid lines in (a) show the behavior predicted by the Bolton theory as contained in Eqs. (3) and (4). The lines in (b) are smooth curves through the points.

Bolton theory is in good agreement with the data for R_H for the frequencies shown in Fig. 2(a). The 12-MHz data (not shown) were only in qualitative agreement, the reason for which is not clear; it may be because of the more severe coupling to spurious modes of the transducer at lower frequencies. The data for both R_H and X_H agree qualitatively with Eq. (6) where k and α are derived from Eq. (3); note that we would expect precise agreement with classical acoustic impedance theory at high temperatures but only qualitative agreement in the zero-sound regime.

Taken together, the transmission and ringing experiments provide convincing proof for the existence of transverse zero sound in normal ^{3}He.

We have extended the ringing experiments into the A and B superfluid phases and observe a peak in both components of the acoustic impedance just below the superfluid transition. Unfortunately, reliable transmission experiments throughout the superfluid region are not possible with our technique because the change in the acoustic impedance upon entering the superfluid can upset our cancelation of the spurious signal. Details of these measurements will be published shortly.

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Pressure Dependence of Reflectivity of Se: Experimental Evidence for Large Local-Field Corrections*

Marc Kastner and R. R. Forberg
Department of Physics and Center for Materials Science and Engineering, Massachusetts Institute of Technology, Cambridge, Massachusetts 02139
(Received 8 December 1975)

The pressure dependence of the reflectivity of trigonal and amorphous Se has been measured between 1.1 and 4.5 eV using a new technique. The results indicate that local-field corrections are large. Structure near 2 eV suggests a localized excitation in both the amorphous and the crystalline material.

Spectroscopic theories of chemical bonding in solids have proved to be powerful for predicting a wide range of properties of the A^NB^{8-N} crystals (Ge, GaAs, etc.). The starting point of these theories is the identification of the bond strength with a spectroscopically determined energy. In Phillips's theory[1] this energy is determined by the dielectric constant, in Harrison's[2] theory it is the energy of the E_2 peak in the imaginary part of the dielectric function. In the case of the covalent A^NB^{8-N} crystals, it may be reasonable to assume that microscopic fields and the modification of matrix elements resulting from exchange are small. (All of these effects will be called "local-field corrections" in this Letter.[3]) In such solids the charge density is fairly uniformly distributed around the atoms. On the other hand, in very ionic or molecular solids the charge density may be quite localized. In this case the determination of the bond strength from the spectroscopic energy is strongly modified by local-field corrections.[4] Therefore, in order that spectroscopic theories of bonding can be extended to ionic and molecular materials, local-field corrections must be better understood. This Letter presents the results of a new technique for studying local-field corrections in solids.

Since local-field corrections are density dependent, the pressure dependence of the optical properties of solids can indicate the size of these effects if the electrons' energy states are not strongly density dependent. For example, measurements of the pressure coefficient of the infrared refractive index for amorphous alloys containing Ge and Se show that Se-rich alloys obey the Lorenz-Lorentz relation,[5] indicating large local-field effects in these materials. The pressure dependence of the reflectivity will reveal more information about local-field corrections, and since for Se these corrections are expected to be large, Se is an ideal material for study.

Early measurements of the reflectivity as a function of pressure were limited to observing energy shifts of the reflectivity peaks.[6] Measuring the variation of the magnitude of the reflectivity is difficult because the refractive indices