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Nonlinear ultrasonic excitations and Landau damping in the Fermi liquid ^{3}He

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Pulsed sound-transmission experiments at different pressures (0–16 bars), frequencies (10–50 MHz), and sound-pulse energies (0.01–110 nJ) demonstrated that Landau damping exists in a Fermi liquid even for low-frequency phonons at temperatures as high as 1 mK due to multiple-phonon absorption. This nonlinear damping of zero sound increases with the square of the frequency. A pronounced sound power dependence of the quasiparticle collision time was found as well as an additional damping of first sound at the third harmonic frequency. [S0163-1829(97)51818-2]

In the past few years, nonlinear effects of sound attenuation in ^{3}He have gained growing interest, especially the coupling of zero sound[1] to the order-parameter modes in superfluid ^{3}He-B.[2,3] In addition, pair breaking due to two-phonon absorption has been studied experimentally.[4] The extension of these more recent experiments towards lower temperatures and higher sound intensities led to several unexpected observations which were briefly described in Ref. 5. One of these observations was that the sound attenuation coefficient α increased with increasing sound intensities and decreasing temperatures. It was speculated[6] that this behavior was due to Landau damping.[7] An indication of quantum damping was seen in experiments at very high sound frequencies ($\leqslant$ 590 MHz).[8,9] In this paper we describe experiments at much lower frequencies which clearly demonstrate that quantum-damping of sound can take place in a Fermi liquid at much higher temperatures, as predicted by Landau,[7] and that this new nonlinear damping effect of sound is caused by quasiparticle excitations due to the coherent absorption of many phonons.

According to Landau, the attenuation coefficient of vibrational excitations with frequency ω can be written as $\alpha=\alpha_0[1+(\hbar\omega/2\pi k_B T)^2]$.[7] Here α_0 is the zero-sound attenuation coefficient which depends on the quasiparticle collision time τ. Depending on $\omega\tau$, a crossover from the collisionless regime (zero sound, $\omega\tau>1$) to the hydrodynamic regime (first sound, $\omega\tau<1$) occurs, and α_0 can be written in the viscoelastic theory in the form $\alpha_0=[(c_0-c_1)/c_1^2]\omega^2\tau_\eta/[1+(\omega\tau_\eta)^2]$ (Ref. 10) with c_0 and c_1 the zero sound and first sound velocities, which depend on pressure p but only very slightly on temperature. τ_η depends on the quasiparticle collision time τ (Ref. 11) and hence on pressure and the square of the inverse of temperature according to $\tau_\eta(p)T^2=\mathrm{const}(p)=\tau_0(p)$. We will show in this paper that this relation has to be modified and that τ_0 also depends on sound intensities.

For $T\rightarrow 0$ the damping of sound becomes temperature independent and varies with frequency as $\alpha=[(c_0-c_1)/c_1^2](\hbar/2\pi k_B)^2\tau_0^{-1}\omega^2$. This kind of sound attenuation is called Landau damping or quantum damping. The crossover between zero sound and quantum damping occurs at a frequency $\omega\geqslant(2\pi k_B/\hbar)T$. For ultrasonic excitations in the *linear* regime of sound and a frequency as high as 100 MHz, quantum damping gains importance only at temperatures $T<$ 1 mK. Since ^{3}He becomes superfluid below this temperature, and since most of the experimental investigations had been done at frequencies below 100 MHz, Landau damping was hidden in the experiments. However, in *nonlinear* sound of high intensity many phonons of the same frequency can be absorbed coherently causing quantum damping even at low frequencies and relatively high temperatures.

Here we report on sound transmission experiments in the frequency regime 10 to 50 MHz which were performed as described elsewhere.[4,5,12] The sound pulses were excited by a quartz transducer with an energy deposition for single sound pulses (typical durations of 5 μs) varying from 1 pJ to 100 nJ. This amount of heat was deposited into a free ^{3}He volume of $4\times4\times6$ mm^3. At $T\sim$ 1 mK, this heat was conducted to the cooling stage within a thermal relaxation time of $\sim$30 s. Pulse power calibration was done by a calorimetric method in the temperature range 3 to 5 mK as described elsewhere.[4,12] We checked for overheating effects by various techniques[4,12] and ensured that our results were not influenced by heating even with sound pulses of highest power. In order to improve the signal-to-noise ratio of the received signal we used a cold preamplifier as well as digital filtering techniques and digital lock-in techniques which we will describe in detail elsewhere.[13]

The sound attenuation coefficient was determined from the first signal detected by the receiving quartz after the sound pulse was transmitted (distance $x=4$ mm) through the liquid according to $\alpha(T)=(1/x)\{\ln(S_0)-\ln[S(T)]\}$ which compares to the theoretical expression

$$\alpha(T)=\frac{c_0-c_1}{c_1^2}\,\frac{\omega^2\tau_\eta}{1+(\omega\tau_\eta)^2}\left[1+\left(\frac{\hbar\omega}{2\pi k_B T}\right)^2\right].$$

Here $S(T)$ denotes the detected sound amplitude at temperature T, and S_0 is a constant which depends on the experimental setup, and on how effectively at each frequency electrical energy is converted into sound energy (or vice versa) by the quartz transducers. Therefore, S_0 must be determined independently for each frequency and power used in the experiments. Because of high sound attenuation in the transition regime between zero sound and first sound, echo techniques cannot be used for a calibration of S_0. Therefore, we

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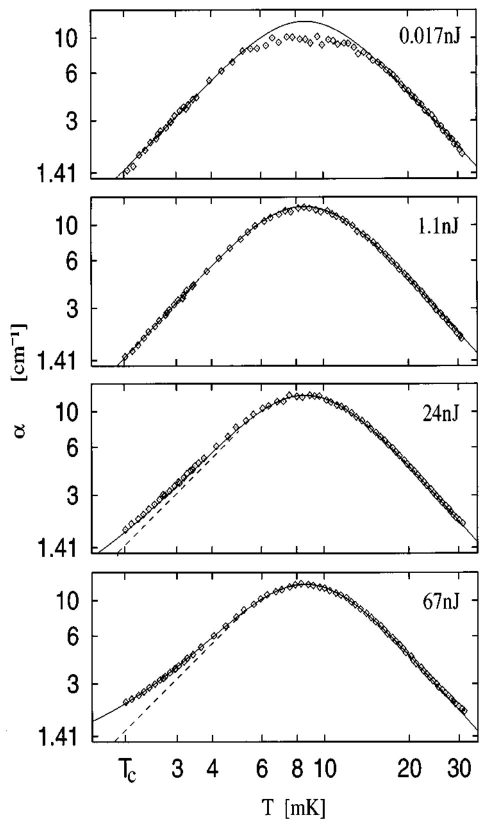

FIG. 1. Dependence of the sound attenuation coefficient α on temperature and sound-pulse energy at frequency 10 MHz and pressure 10.5 bars. The lines are fits to the visco-elastic theory including quasiparticle excitation by multiple-phonon absorption. The discrepancy between fit and data at $E=0.017$ nJ is caused by the spectrometer resolution. The dashed lines display the T^2 law of linear zero sound.

had to follow a different approach to obtain absolute values for α: In the linear regime of sound, all Fermi-liquid parameters entering the attenuation coefficient are known and hence α is known. Thus, $S_0=S(T)\cdot\exp[\alpha(T,p)x]$ is in principle fixed by the data. In the nonlinear regime of sound, S_0 has to be treated as a free parameter and must be obtained from a fit to the data $S(T)$.

For the determination of $S(T)$ from our data, we used three different methods[12] which were sensitive in different time windows after the pulse. All three methods yield identical results for the attenuation coefficient α in the covered temperature range ($T_c<T<200$ mK) and at all investigated pressures ($p=0-16$ bars), frequencies (10–50 MHz) and time windows (0–25 μs, 0–35 μs, and 0–70 μs). This indicates that time-dependent heat currents as well as interference effects due to reflections at the cell walls do not play a significant role. In this paper, we present the data taken in zero magnetic field at 10 MHz and at 10.5 bars. However, the reported effects were also seen at all other pressures and frequencies.

In Fig. 1, the attenuation coefficients determined below 30 mK are shown for sound power levels which vary over more than three orders of magnitude. For the lowest sound energies ($E=0.017$ nJ) it is obvious that the resolution limit of our spectrometer was reached at temperatures $6<T<13$ mK. At sound energies up to about 1 nJ, we found for α the expected temperature dependencies of T^{-2} and T^2 for first sound and zero sound, respectively. At higher sound intensities, however, only first-sound attenuation still varied with T^{-2} whereas zero-sound attenuation deviated more and more from the T^2 law with decreasing temperature and increasing power (compare dashed lines in Fig. 1). We attribute this behavior to quantum damping due to the excitation of quasiparticles by multiple-phonon absorption. In such an excitation, N phonons of frequency ω are coherently absorbed in a time shorter than the quasiparticle lifetime yielding an observable damping which becomes temperature independent for $T\rightarrow 0$, and varies with the square of the phonon frequency. It can be shown that Landau's original equations[7] can be reformulated without restriction of validity in terms of N phonons of energy $\hbar\omega$ yielding $\alpha=\alpha_0[1+(N\hbar\omega/2\pi k_BT)^2]$.

In our experiments, the sound intensity was chosen high enough that multiple-phonon excitation of quasiparticles should not be negligible. The total number n_0 of 10 MHz phonons generated by the strongest sound pulse of energy $E\sim$ 100 nJ can be estimated to be on the order of $n_0=E/\hbar\omega\sim 10^{20}$. These phonons are generated in about 4 μs, the ringing time of the transducer. Thus, within the quasiparticle collision time of about 0.5 μs around 1 mK, roughly $n\sim 10^{19}$ phonons can excite or be scattered by quasiparticles. This number of phonons (n) is about 1000 times the number of thermally excited quasiparticles (n_{qp}). Thus, this nonlinear damping effect can become very effective.

How well the modified Landau model fits our data is demonstrated in Fig. 1. The solid lines in Fig. 1 are fits to the data using $c_1=297.22$ m/s (Ref. 14) as a fixed value. In linearized theory, c_0 would also be fixed, but c_0 might depend on sound intensity. Thus, free fitting parameters were c_0, S_0, τ_0, and N, the average number of coherently absorbed phonons. The results for the dependence of zero-sound velocity, viscous collision time and number N of multiple phonon absorption on sound intensity are compiled in Fig. 2.

For low sound energies ($E<$ 1 nJ) the 10 MHz phonon density is low and the ratio n/n_{qp} is less than one. In this regime, α is insensitive to $N(N=1)$ and c_0 = 300.93 m/s is close to what is expected [300.95 m/s (Ref. 14)]. The viscous collision-time $\tau_0=\tau_\eta T^2$ = 1.16 μs(mK)^2 agrees well with the value 1.02 μs(mK)^2 of Abel *et al.*,[14] if one takes into account that in Ref. 14 the absolute temperature scale has to be shifted by -6% in order to agree with the Greywall scale[15] used in our experiments.

This agreement at small sound powers gives us confidence in the new results at high sound power (see Fig. 2), where the ratio n/n_{qp} reaches values up to 1000 while c_0 drops to 300.56 m/s, τ_0 increases to 1.24 μs(mK)^2, and N rises to 27, respectively. The tiny change in zero-sound velocity at high sound intensities can be explained by the simultaneous increase in the attenuation coefficient. The 8% increase in $\tau_\eta T^2$ is not understood yet. A more systematic investigation for the power dependence of τ_η is required. The lines in Fig. 2 are guides to the eyes and display a fit to the data assuming a power-law dependence on the pulse energy E [J]. We could parametrize our results with c_0 [m/s] =

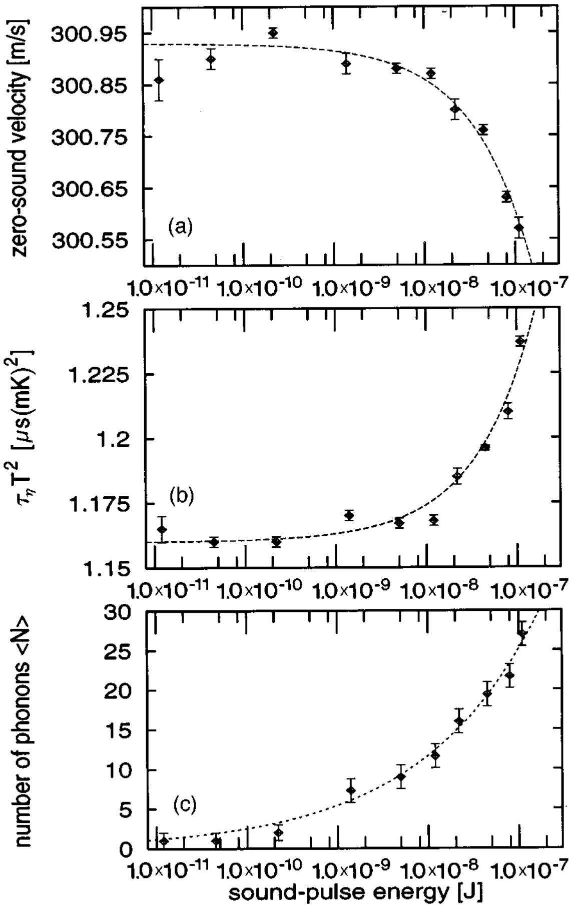

FIG. 2. Plotted versus sound-pulse energies are the (a) zero-sound velocity, (b) viscous collision time, and (c) number of coherently absorbed phonons. The lines are guides to the eyes with sound energy E dependences of $E^{2/3}$, $E^{2/3}$ and $E^{1/3}$ for (a), (b), and (c), respectively.

300.93–15000$E^{2/3}$, $\tau_\eta T^2$ [μs(mK)^2] = 1.16 + 3000$E^{2/3}$, and $N=1+5400E^{1/3}$. An understanding of these dependencies is lacking at present.

If N phonons of a single frequency ω coherently produce a particle-hole excitation, this excitation may decay under emission of several phonons of different frequencies ω_i. The observation of a frequency distribution with $0<\omega_i \leq N\omega$ would be expected, because the density of quasiparticle states is quasicontinuous. Thus, a time-dependent investigation of the frequency distribution in the received nonlinear sound signal would in principle yield information about the quasiparticle dynamics. For such investigations, a broadband detector in connection with a 100 MHz bandwidth detecting electronics is needed. The latter was especially developed,[13] but with our quartzes only odd harmonics of the sound-pulse fundamental frequency could be detected in a frequency band of roughly 200 kHz. Therefore, if any conversion from phonons of the fundamental to higher harmonic frequencies takes place due to nonlinear effects in the Fermi liquid, the amplitude of the higher harmonic Fourier coefficients would be expected to be small.

Unfortunately, nonlinear elements in the electronics also produce harmonic frequencies (even and odd harmonics) against which one has to discriminate in order to determine the contribution generated in the liquid. Figure 3 displays the

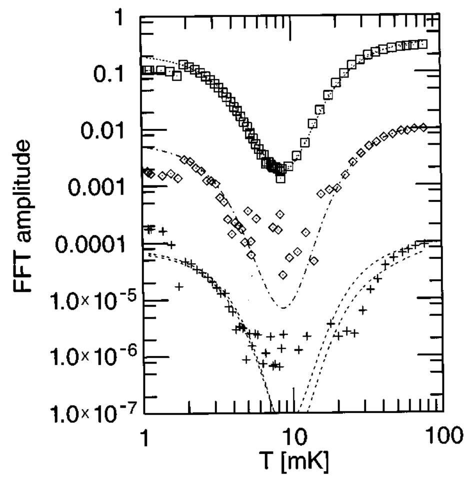

FIG. 3. The Fourier-transform amplitudes of the detected signals at 10 (squares), 20 (diamonds), and 30 MHz (crosses) after a ''pure'' 10 MHz sound-pulse excitation of 110 nJ energy. For clarity, the 30 MHz data are shifted by a factor of 0.01. The larger scatter of the 20 and 30 MHz data around 10 mK is due to the limited resolution of the spectrometer. The curves are fits according to the visco-elastic model when harmonic distortions in the transmitter and receiver electronics are also taken into account (see text).

magnitude of the signals $S(T,j\omega)(j=1,2,3)$ as obtained by Fourier transformation of the signal received after a 10 MHz pulse excitation with a pulse energy of 110 nJ. The second harmonic $(j=2)$ 20 MHz signal is purely due to nonlinearities in the receiving electronics, because the quartzes are almost insensitive to even harmonics. Beside contributions due to distortions of the receiver, the 30 MHz signal $(j=3)$ contains also contributions from sound in the liquid. In order to determine the amount of sound generated by a 30 MHz excitation with the third harmonic of the transmitter, and how much is really caused by nonlinear effects in the liquid (initiated by the strong 10 MHz excitation), the electronically generated effects must be eliminated. If the harmonic distortions $S_{el}(j\omega)$ are generated by the electronics, they must depend on the amplitude of the driving signal $S(\omega)$. These electronically generated contributions can be modelled by $S_{el}(j\omega)=AS^B(\omega)$ where A and B are adjusted to the data (see for instance the 20 MHz data in Fig. 3). Using this simplified model for the electronic contributions and the visco-elastic model (including Landau damping) for the sound signals, we could fit our data as shown in Fig. 3.

The directly excited 10 MHz sound signal (squares in Fig. 3) can be explained very well when Landau damping is added to the visco-elastic theory. The electronically produced second and third harmonics at 20 and 30 MHz have 5 and 1% of the amplitude of the received 10 MHz sound signal, respectively. However, in the third harmonic signal, contributions of 30 MHz sound are also present which has about 4% of the 10 MHz sound amplitude (note, the 30 MHz data are shifted by a factor 0.01 in Fig. 3 for clarity). Obviously, in the zero-sound regime, the third harmonic signal can fully be explained within this model. However, this does not hold in the superfluid[16] as well as in the first-sound regime. The measured third harmonic signal deviates around 30 mK from the visco-elastic model even when phase differences $\Delta\phi$ between the receiver- and transmitter-caused distortions are taken into account (for instance $\Delta\phi = 0°$ and $60°$ for the two dotted lines in Fig. 3, respectively). The origin of this additional attenuation is unknown, but it is only present at the highest sound powers. Below $T = T_c = 1.85$ mK (at $p = 10.5$ bars) ^{3}He becomes superfluid which can be identified in Fig. 3 by the sharp signature in attenuation. In superfluid ^{3}He-B, the situation becomes more complex.[12] But around 1 mK the signals shown in Fig. 3 are also influenced by quantum damping.[16]

Our experiments with 30 and 50 MHz sound-pulse frequencies confirm these observations at 10 MHz. In addition, they also confirm that quantum damping of zero sound in the Fermi liquid ^{3}He varies with ω^2. This is an additional proof that multiple-phonon absorption takes place in the lifetime of quasiparticle excitations. It also demonstrates that Landau's description of vibrations in a Fermi liquid holds down to lowest temperatures.

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