in progress.

In conclusion, we have demonstrated numerically the existence of shear Alfvén waves in tokamak geometry. Both the drift wave and the shear Alfvén wave are stable in a collisionless plasma. When trapped electrons are included, both waves can be unstable. However, their behaviors with beta are different. While the drift mode is stabilized by effects of beta, the shear Alfvén mode is destabilized. Of course, when the poloidal beta approaches $1/\epsilon$, the simple concentric, circular flux-surface model used to calculate the trapped-electron response and the slab model underlying Eqs. (1) to (4) are no longer valid. For these reasons, the high-beta ends of the curves in Figs. 1, 2, and 4 are not very accurate but should still represent the beta dependence of the usual drift and the shear Alfvén branch in a tokamak with $\beta \lesssim 5\%$.

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Incipient Superfluidity in Liquid ^{3}He above the Superfluid Transition Temperature

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The large zero-sound attenuation just below the transition temperature T_c to a superfluid state in liquid ^{3}He is anticipated above the transition by an excess attenuation over that for a normal Fermi liquid. The excess is greatest at T_c but can be readily discerned at temperatures several percent above T_c. The excess is largest near melting pressure and exists adjacent to both ^{3}He-*A* and ^{3}He-*B*.

The attenuation of zero sound[1-3] in a normal Fermi liquid is frequency independent and proportional to T^2. We have discovered that in liquid ^{3}He at temperatures above the transition temperature T_c the zero-sound attenuation rises above the extrapolated value for normal zero-sound attenuation. The excess attenuation is very significant, readily measurable, and can be discerned at temperatures as much as 10% above T_c. That such an effect might exist for zero sound was pointed out some time ago by Paulson, Johnson, and Wheatley.[4] A small effect above T_c was observed in static magnetization by Paulson, Kojima, and Wheatley[5] and in viscosity coefficient by Parpia, Sandiford, Berthold, and Reppy,[6] the measurable effects extending above T_c by only a tenth of a percent of T_c. It is possible that the excess zero-sound attenuation reflects the fluctuation superfluidity in fermion systems first discussed by Thouless[7] (equilibrium properties) and by Emery[8] (transport properties). In particular, Emery[8,9] has stressed the possible importance of fluctuation superfluidity to an understanding of the properties of liquid ^{3}He, while fluctuations in *superconductors* have proven to be of great interest.[10] Although there does not yet exist a theory of the excess attenuation which we describe here, we hope that this work will stimulate a theoretical interpretation so that, in analogy to superconductivity, a deeper understanding of the superfluid state will result.

Our measurements were made in a nuclear

cooling cryostat with very precise temperature control and the ability to maintain temperatures for very long periods of time.[11] The sound apparatus was contained in a magnetically shielded tower above the nuclear refrigerator, the sound being propagated across a 12.7-mm gap using 5-MHz–fundamental quartz transducers. Propagation could be studied at 5, 15, and 25 MHz. The temperature was measured magnetically using 0.21 g of powdered $Ce_{0.05}La_{0.95}$ magnesium nitrate in the form of a right circular cylinder with diameter equal to height. The thermometer and sound cell were connected thermally by a $5\frac{1}{2}$-mm-diam, $3\frac{1}{2}$-cm-long column of ^{3}He which gave excellent thermal contact and allowed a high reproducibility of the attenuation-temperature data: T_c could be reproduced to several parts in 10^5. The thermometer is not only precise but apparently nearly as accurate, regarding Kelvin temperatures, as other thermometers.[12] This point is crucial. For example, in the work of Paulson, Johnson, and Wheatley,[4] where a concentrated cerium magnesium nitrate (CMN) thermometer was used, it was not possible to separate unambiguously temperature scale from ^{3}He effects.

Our raw measurements of sound signal amplitude were corrected for two effects. First, because of the long ^{3}He sound path, attenuation of the signal can be high and a significant amount of sound can propagate to the receiver crystal through non-^{3}He paths. This effect was measured by setting the temperature at a value for which ^{3}He attenuation was very large and corrected for by assuming this signal to be independent of the directly propagated ^{3}He signal. We rejected all data for which the ^{3}He sound signal was less than 10 to 15 dB above this background. Second, especially at higher pressures where the excess attenuation effect is largest, zero sound may not be fully developed in the temperature range of the measurements, which must extend well above T_c to account for the normal–Fermi-liquid zero-sound attenuation. Accordingly we further corrected the amplitude-attenuation-coefficient data to zero-sound attenuation α_0 from measured (corrected as above) attenuation α using the formula[3] $\alpha_0^{-1} = \alpha^{-1} - \alpha_1^{-1}$, where α_1 is a calculated first-sound attenuation. For the latter we used data given by Wheatley.[13] We also considered the further correction to this formula by Egilsson and Pethick.[14] The data to be presented in this Letter are at 25 MHz, where the correction for α_1 is smallest (a maximum of +0.46%) and hence where the results are probably most accurate.

A graph which shows best the *departures* of the zero-sound attenuation from its normal–Fermi-liquid behavior is given in Fig. 1 where $[\alpha_0/(\alpha_0)_c](T/T_c)^{-2}$ is plotted against T_c/T, where $(\alpha_0)_c$ is the zero-sound attenuation at T_c. This is a second generation graph in that we first plotted $[\alpha_0/(\alpha_0)_c](T^*/T_c^*)^2$ vs T_c^*/T^*, since if Kelvin and magnetic temperatures are related by the simple formula $T = T^* + \Delta$ then such a plot should at high enough temperatures be a straight line with slope to intercept ratio of $2\Delta/T_c^*$. We were not able to distinguish, especially in the absence of a theory of the temperature dependence of the excess attenuation, between a small excess attenuation and a small Δ. {The data shown are for 25 MHz and 30.87 bars, with $T_c^* = 2.5865$ mK and $\alpha_c = 1.602$ cm^{-1} [$(\alpha_0)_c = 1.610$ cm^{-1}].} The magnetic transition temperature at melting pressure is 2.625 mK. Taking $\Delta = 0.1$ mK brings the temperature T_c at melting pressure into the mid-

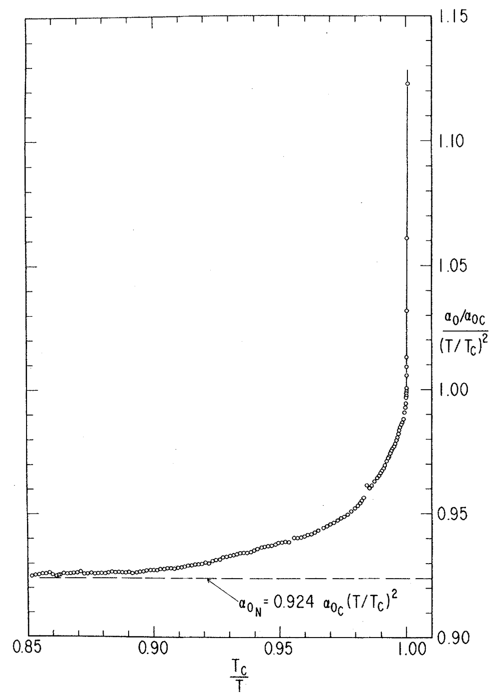

FIG. 1. Plot of $[\alpha_0/(\alpha_0)_c](T/T_c)^{-2}$ against T_c/T in liquid ^{3}He for a frequency of 25 MHz and a pressure of 30.87 bars. Here $T = T^* + 0.1000$ mK and T^* is measured magnetic temperature, α_0 is zero-sound attenuation, and $(\alpha_0)_c$ and T_c are, respectively, the zero-sound attenuation and the temperature at the transition to the superfluid state. In a normal Fermi liquid $[(\alpha_0)_N/(\alpha_0)_c](T/T_c)^{-2}$ is a constant; the dashed line at 0.924 corresponds to a possible relation $(\alpha_0)_N = 0.924(\alpha_0)_c(T/T_c)^2$, with $(\alpha_0)_c = 1.61$ cm^{-1}; $T_c^* = 2.5865$ mK, and T_c taken to be 2.6865 mK.

dle of previous determinations[13,15] and also tends to reduce the limiting slope at small T_c/T of the data in Fig. 1. We know from our thermometry studies[12] that Δ could be zero but could not be much larger than 0.1 mK. Accordingly in Fig. 1 the temperatures used are obtained from T^* by adding 0.1 mK since we believe that this is a probable correction. However, no qualitative conclusion would be altered by plotting raw data using magnetic temperatures only.

Examination of Fig. 1 shows the qualitative dramatic increase of attenuation at T_c on lowering the temperature, an increase sufficiently rapid as to enable T_c to be located to a precision of 1 in 10^5. Although experimental inaccuracy and imprecision do not permit an unambiguous determination of the excess attenuation at T_c, the data in Fig. 1 suggest that the extrapolated normal zero-sound attenuation at T_c is about 8% less than the value actually measured at T_c. The observations at 15 MHz give essentially the same result. We have also made measurements at a higher pressure, 32.56 bars; an intermediate pressure, 19.94 bars; and a low pressure, 0.05 bars. The effect is present over the whole range so that its existence does not depend on the liquid above T_c being adjacent to ^{3}He-A or ^{3}He-B, but it decreases with decreasing pressure, the excess attenuation at T_c being $6\frac{1}{2}\%$ of α_c at 19.94 bars and very

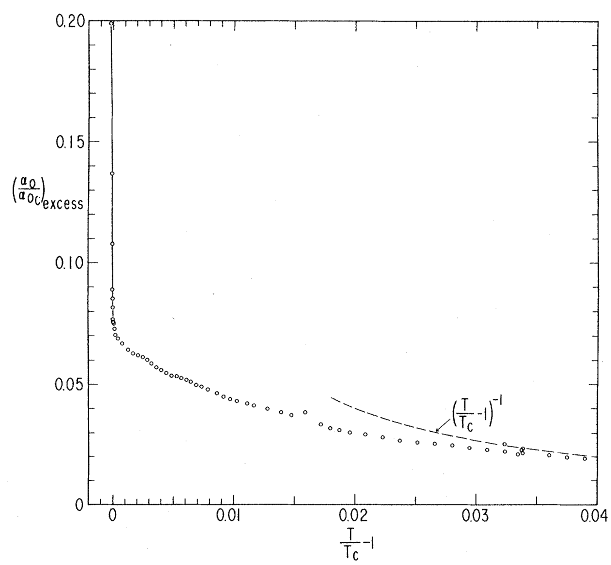

FIG. 2. Reduced excess attenuation $\alpha_0/(\alpha_0)_c - 0.924(T/T_c)^2$ as a function of reduced-temperature difference $T/T_c - 1$ for the data of Fig. 1. The line showing a $(T/T_c - 1)^{-1}$ temperature dependence is drawn for comparison purposes.

approximately 2% of α_c at 0.05 bars.

To permit a closer, though approximate, view of the excess zero-sound attenuation near T_c, we assumed that $(\alpha_0)_N/(\alpha_0)_c = 0.924(T/T_c)^2$, as in the horizontal dashed line in Fig. 1, and plotted in Fig. 2 the quantity $[\alpha_0 - (\alpha_0)_N]/(\alpha_0)_c$ from the data in Fig. 1 as a function of the reduced-temperature difference $[(T/T_c) - 1]$. Here $(\alpha_0)_N$ is normal zero-sound attenuation. This plot has no simple temperature dependence; we show a $[(T/T_c) - 1]^{-1}$ dependence on the figure. Data at 15 MHz are similar to those shown, suggesting that there is not a major frequency dependence at a given pressure. This in itself may be an important clue since the ultrasonic attenuation below T_c can have a strong frequency dependence, excepting temperatures very close to but below T_c,[16] where pair breaking is thought to dominate the attenuation.[17]

In addition to the slight drop in viscosity coefficient η observed by Parpia *et al.*[6] in a temperature interval above T_c of about $10^{-3}T_c$, these authors also reported a substantial decrease in η below the value for a normal Fermi liquid, distinguishable to temperatures perhaps as much as 30% above T_c. They used a concentrated CMN thermometer and plotted η^{-1} vs T^{*2}. We have found even with our diluted CMN thermometer that on such plots it is difficult to distinguish between an effect of Δ in $T = T^* + \Delta$ and actual deviations of the ^{3}He properties from their normally expected properties. This was certainly an impediment to interpretation of our earlier observations.[4] Parpia *et al.*[6] noted that since the effect increased as the pressure decreased it was not entirely caused by artifacts of the temperature scale but could be caused by the quasiparticle free paths becoming comparable to the plate separation in their torsional pendulum. The effect observed in the present work has the opposite pressure dependence in an open geometry and is closely associated with the second-order transition.

One may question why incipient superfluidity of ^{3}He can be seen for a considerable temperature range above T_c using zero sound while effects on susceptibility and viscosity are minor. The answer may be that in the susceptibility and viscosity observations we are dealing with changes in processes which are fully developed above T_c while for zero sound new and very powerful mechanisms for sound attenuation are known to become available below T_c. In this sense, zero-sound attenuation in liquid ^{3}He above T_c is similar to electrical resistance and diamagnetism in superconductors above their transition temperatures.

Note added.—Following a suggestion by Professor Maki, we find that an excellent fit to the experimental data shown in Figs. 1 and 2 for the excess ultrasonic attenuation over that extrapolated for normal zero sound is

$$[\alpha_0/(\alpha_0)_c]_{\text{excess}} = 0.07_6 - 0.327(T/T_c - 1)^{1/2},$$

for $(T/T_c - 1) > 0$. Below the transition temperature, $1.2\times 10^{-3} > (1 - T/T_c) > 5\times 10^{-5}$, the excess attenuation is fitted by

$$[\alpha_0/(\alpha_0)_c]_{\text{excess}} = 0.090 + 3.9\times 10^2(1 - T/T_c).$$

For $5\times 10^{-5} > (1 - T/T_c) > 0$, the excess attenuation rises somewhat more rapidly than this formula suggests.

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Thermorotation Effects in Superfluid Helium

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We have observed a variety of thermorotation effects, effects involving thermal counterflow in rotating superfluid helium. In particular, the temperature and chemical-potential gradients associated with the motion of vortex lines have been measured.

It is generally believed that dissipative superfluid flow is associated with the motion of vortices through the superfluid.[1] Anderson[2] has pointed out that vortices crossing a line connecting two points at rest in the superfluid produce a chemical-potential difference between the two points proportional to the rate of crossing. This is a consequence of an "ac Josephson effect" in helium in which a chemical-potential difference is proportional to the rate at which the order-parameter phase difference changes—each vortex passage producing a phase change of 2π. There have been a number of experiments[3] which demonstrate that vorticity is associated with turbulent counterflow or pressure-driven flow but these experiments are impossible to directly relate to Anderson's prediction. There is evidence[4] that charged vortex rings pushed towards an orifice connecting two baths can produce a level difference between the baths. This, however, is a situation in which there is no dissipative superflow and, in any event, quantitative confirmation of Anderson's prediction is lacking. In this Letter, we describe a series of experiments in which an array of parallel vortex lines is moved through the fluid in a relatively controlled manner and the resultant temperature and chemical-potential gradients are observed. In addition, we report the observation of a number of interesting effects associated with the fact that the vortex array is confined within a channel of finite dimensions.

