Phase Diagram of Superfluid ^{3}He-A_1

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(Received 30 July 1984)

The A_1 and A_2 transitions, which delineate the A_1 phase of superfluid ^{3}He, are measured in magnetic fields up to 3 T over a wide range of pressures. At zero pressure the A_1-A_2 splitting is found to be smaller than at melting pressure by a factor of 5. Strong-coupling effects are similarly smaller by an order of magnitude.

PACS numbers: 67.50.Fi

The thermodynamic properties of the A and B phases of superfluid ^{3}He are widely measured and relatively well understood. On the other hand, despite the fact that the A_1 phase was discovered over ten years ago,[1] experimental difficulties have precluded precise thermodynamic measurements in this phase except at melting pressure.[2,3] Even though it was understood from the beginning that the A_1 phase resulted from BCS pairing in only one spin population,[4] only recently has it been shown that it is the down spins which pair.[5]

The A_1 phase of superfluid ^{3}He exists only in a magnetic field and in a narrow temperature region around the zero-field transition, T_c, from normal fluid to superfluid. The A_1 and A_2 transitions, which delineate the A_1 phase, split linearly in field and asymmetrically about T_c, A_1 moving up and A_2 moving down. The total thermal width of this phase is only a few tens of microkelvins per tesla.

We have measured the magnetic field dependence of the A_1 and A_2 transitions in fields up to 3 T by their sudden and dramatic effect on the attenuation of zero sound.[6] More importantly, we have accurately measured the distance of each transition from the zero-field transition, not just the total splitting. It is this feature which strains the capability of ultralow-temperature thermometry. An error in locating the zero-field transition of only a few microkelvins would give substantial errors in the derived magnitude of strong-coupling effects.

The experimental difficulties in any A_1 experiment arise directly from the presence of the magnetic field required to stabilize this phase. Thermometry in a high field is not a solved problem though it is crucial to this experiment. There are three possible solutions. First, one could try to develop a thermometer which is truly field independent, yielding at most 1 μK/T offset. Second, one could place a thermometer far from the high-field region and connect it to the sample by a strong thermal link, typically silver bars. Both approaches allow very large magnetic fields to be used, but place stringent constraints on the long-term reproducibility of the thermometers, and are incapable of discriminating against field-dependent heat leaks, field-dependent Kapitza resistances, and field-dependent conductivity in a thermal link.

We have developed a third alternative, shown in Fig. 1. We apply the large magnetic field only where it is required at the top of one of our experimental towers (the "A_1 cell") occupied by one pair of zero-sound crystals. These detect the A_1 and A_2 transitions. Depending upon the pressure, the crystals are driven at either 7 or 21 MHz to maximize our sensitivity to the transitions. A short tube of ^{3}He connects the A_1 cell to the nuclear demagnetization stage. The adjacent tower contains a conventional lanthanum-cerium magnesium nitrate

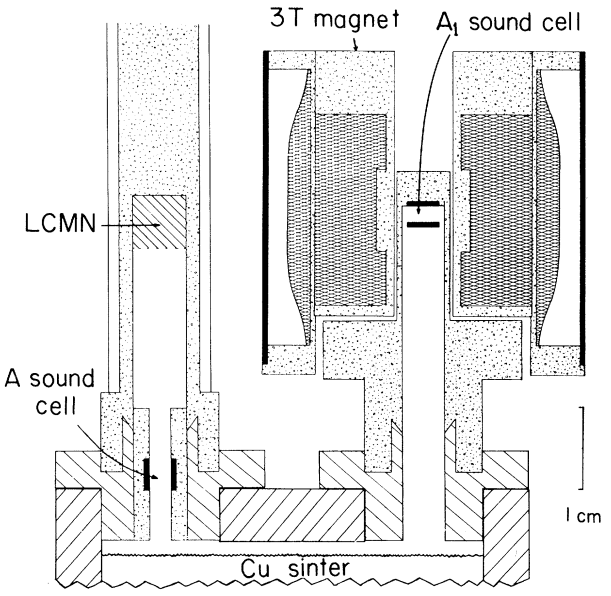

FIG. 1. Schematic drawing of experimental apparatus. A continuous liquid path connects the A_1 cell in a high field to the LCMN thermometer and the A cell in zero field. The speckled regions are epoxy.

(LCMN) thermometer and a second pair of zero-sound crystals (the "A cell") which mark the A transition, T_c. Both cells operate concurrently. The A-cell marker anchors our high-field points to the zero-field transitions, insulating them from long-term drifts in thermometry. Since all heat flows required to establish thermal equilibrium are through ^{3}He only, field-dependent offsets to T_c should be minimized. The calculated time constant for thermal equilibrium within the liquid (not to the copper stage) is typically 10 sec. In fact, the dominant time constant in our experiment is the relaxation time of the LCMN thermometer itself, which is roughly 1 min.

Our magnet generates fields up to 3 T which are homogeneous to 0.1% over the A_1 cell and known to an accuracy of 0.1%. Importantly, this magnet confines the field so that adjacent sensors are protected. For example, we have never observed any interaction of this magnet with our thermometer. This magnet has already been described by Israelsson and Gould.[7] It is thermally linked to our mixing chamber and carefully isolated from the experimental tower.

At a given pressure data were usually taken in two fields, 2 T and 3 T. At each field all three transitions were measured within 6 to 9 h. Changing the magnetic field by 1 T warmed the sample chamber by 60 μK, presumably by eddy currents on the A_1 sound crystals. Both fields at a single pressure could be completed in 18–24 h. The pressure was then slowly reduced by 3 bars, resulting in more warming of 50–100 μK. Several pressures could thus be completed in a single demagnetization run.

Every transition at every pressure and field was measured a minimum of three times by sweeping the temperature in both directions at typical rates of 5–10 nK/sec. Plotting the apparent transition temperature versus sweep rate yielded a straight line whose slope corresponds to the LCMN relaxation time of 1 min, and whose intercept is the equilibrium transition. At this point a small but reproducible dependence of our thermometer on the nuclear demagnetization field was corrected. This moved the A_1 and A_2 transitions typically 0.4 μK. (The A_1 magnet showed no such effect.) Fitting two straight lines to the A_1 and A_2 transitions revealed a zero-field T_c offset between the A_1 and A cells which was pressure independent and equal to 3.2 ± 0.9 μK (systematic). This 0.9 μK uncertainty ultimately limits our accuracy. The slopes of the two fitted straight lines comprise our results.

Our results are shown in Fig. 2 and Table I. The

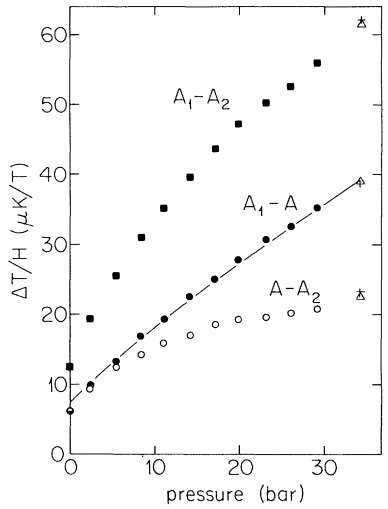

FIG. 2. Field dependence of the A_1 and A_2 transitions. Our results are given for the A_1 slope (solid circles), the A_2 slope (open circles), and the width of the A_1 phase (squares). These are compared to melting-pressure results of Ref. 2 (crosses) and Ref. 3 (triangles). The curve is a smoothed version of the theoretical prediction for η from Ref. 16 (ignoring one theory point at 17 bars).

strong pressure dependence of the splittings of the A_1 and A_2 transitions is clearly evident. Our values extrapolated to melting pressure are in reasonable agreement with previous work.[2,3,8] The second and third columns of the table give values for the Landau-Ginzburg coefficients described below.

The last three columns of the table give the total A_1-A_2 splitting according to three different temperature scales. Although LCMN is an excellent secondary thermometer, we do not have an absolute scale. Rather, we have fitted our observed zero-field T_c values (assuming the usual Curie-Weiss susceptibility) to values of T_c at various pressures as tabulated by Paulson *et al.*,[9] by Alvesalo *et al.*,[10] and by Greywall.[11] After measuring our thermometer's sensitivity to pressure we fit each temperature scale to our data with maximum errors of ± 4 μK. Outside of Table I all of our results for splittings use the scale of Ref. 9. The ratio of the A_1 and A_2 splittings is essentially independent of scale.

The thermodynamics of the A_1 phase is well described by the mean-field Landau-Ginzburg expansion. The general expression for the free energy in a p-wave superfluid was given by Mermin and Stare[12] in terms of five undetermined coefficients $\beta_1, \ldots, \beta_5$. In the A_1 and A phases, the free energy including the usual Ambegaokar-Mermin[4]

TABLE I. Measured pressure dependence of the Landau-Ginzburg parameters, and the width of the A_1 phase according to three different temperature scales as discussed in the text.

Pressure (bars)	$10^3\eta$	$-\beta_5/\beta_{245}$	$(T_{A_1}-T_{A_2})/H$ (μK/T) La Jolla scale	Helsinki scale	Greywall scale
0.0	4.26	0.976	13.44	14.93	14.32
2.3	6.31	1.044	19.27	21.39	20.31
5.3	8.55	1.085	25.59	28.38	26.69
8.3	10.77	1.186	30.95	34.26	32.03
11.0	12.37	1.216	35.13	38.85	36.15
14.1	14.49	1.321	39.67	43.84	40.58
17.1	16.07	1.343	43.70	48.29	44.53
19.8	17.86	1.441	47.16	52.06	47.91
23.1	19.77	1.573	50.40	55.59	51.06
26.0	20.84	1.611	52.64	58.06	53.17
29.1	22.60	1.697	55.97	61.73	56.47

parameter η is

$$F=-[(t-\eta h)\Delta_\downarrow^2+(t+\eta h)\Delta_\uparrow^2]+\tfrac{1}{2}\beta_{24}(\Delta_\downarrow^4+\Delta_\uparrow^4)+\beta_{2455}\Delta_\downarrow^2\Delta_\uparrow^2,$$

where $t=1-T/T_c$, $h=\mu H/kT_c$, $\Delta_{\uparrow(\downarrow)}$ is proportional to the up- (down)-spin order parameter, and multiple subscripts to β imply summation. The advantage of the β parametrization is that it applies to all superfluid phases. Information about the parameters in the A_1 phase can be directly used in the B phase or in the cores of vortices in rotating ^{3}He.[13] This does not apply to other parametrizations.[3,14]

The transition temperatures are

$$1-T_{A_1}/T_c=t_1=-\eta h$$

$$1-T_{A_2}/T_c=t_2=\eta h\beta_{245}/(-\beta_5).$$

In ^{3}He it turns out that $\beta_5<0<\beta_{245}$ always. All strong-coupling effects are contained in the β parameters and so we isolate them by plotting in Fig. 3 the up/down ratio $-t_1/t_2=-\beta_5/\beta_{245}$. As a result of our systematic uncertainty in the T_c offset, the ratios are accurate to $\pm 3\%$, approximately independent of pressure. Strong-coupling effects are greatly reduced at low pressures. Extrapolating these results to melting pressure again yields reasonable agreement with prior measurements.[2,3]

Rainer and Serene[15] point out that strong-coupling corrections should vary as T_c/T_F, which is an order of magnitude smaller at zero pressure than on the melting curve as in our data. They also demonstrate that strong coupling corrections as measured by the β parameters depend upon quasiparticle scattering amplitudes in the normal state. Our results, therefore, provide a test for all models of interactions in ^{3}He. All current theories are insufficiently accurate to explain our results.

Recently Bedell and Quader[16] have calculated the field dependence of the A_1 transition, η. We have included their tabulated results in Fig. 2. The agreement is surprisingly good, especially considering the uncertainties in temperature scales.

Simultaneous with this experiment Sagan *et al.*

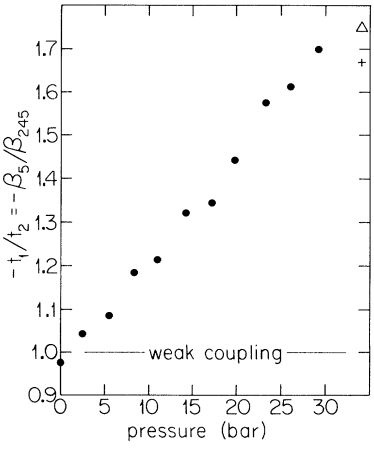

FIG. 3. The up/down ratio (solid circles). This depends only upon strong-coupling corrections. We again compare to the melting-pressure results of Ref. 2 (cross) and Ref. 3 (open triangle). The weak-coupling value is also shown.

have been performing similar measurements[17] which yield comparable A_1-A_2 splittings but systematically smaller up/down ratios. The primary differences between our apparatuses are that Sagan *et al.* reach higher fields than we do, but they have no concurrent indicator of the zero-field transition and further require a distant thermometer. This may account for the discrepancy between our values for the up/down ratio.

We would like to thank D. C. Sagan and D. M. Lee for several communications during the performance of our experiments. One of us (C.M.G.) would like to thank R. C. Richardson for a question asked nine years ago which led to this experiment. This work was supported by the National Science Foundation through Grants No. DMR82-08760 and No. DMR 82-00661. One of us (C.M.G.) is an Alfred P. Sloan Research Fellow.

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[14]The early spin-fluctuation theory of W. F. Brinkman and P. W. Anderson [Phys. Rev. A **8**, 2732 (1973)] introduced the spin-fluctuation parameter δ which turns out to give $-t_1/t_2=(1+\delta/2)/(1-\delta/2)$. S. Takagi [Prog. Theor. Phys. **51**, 1998 (1974)] took the same form of the free energy but changed the definition of δ so that $-t_1/t_2=(1+\delta)/(1-\delta)$. Finally, W. F. Brinkman, J. W. Serene, and P. W. Anderson [Phys. Rev. A **10**, 2386 (1974)] improved the spin-fluctuation calculations which yield $-t_1/t_2=(1+0.35\delta)/(1-0.52\delta)$. The fact that a value of δ in one experiment cannot be carried to others is documented in Ref. 3.

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