Helium three

Magnetic superfluid phases in the low-millikelvin region, a Josephson "ringing" effect and fluid anisotropy are among the exotic properties of this surprising substance.

John C. Wheatley

The past two decades have seen great progress in our understanding of the properties of condensed helium 3. Research on He^3 in our lab (see figure 1 for its modern form) began in 1958. At that time the properties of the more abundant isotope, helium 4, were rather well known: A phase transition at about 2 K transforms it from a normal liquid to a superfluid, a liquid that under some conditions will flow without viscous drag. The explanation for this behavior is rooted in concepts from the condensation of an ideal Bose–Einstein gas. But He^3 obeys Fermi–Dirac statistics; would it show any likeness to a Fermi gas or liquid? Would it also exhibit a transition to superfluidity?

We shall see below that He^3 can indeed be described as a Fermi liquid and that it exhibits transitions to as least three additional phases but—in contrast to He^4—at temperatures of only a few millikelvins. Experiments on spin diffusion, nuclear susceptibility, ultrasonic propagation and other properties have clarified the nature of the normal fluid. The newly discovered low-temperature phases are magnetic superfluids that exhibit such interesting properties as the Josephson effect and fluid anisotropy.

A clue from spin diffusion

The work on He^3 that influenced us most at the beginning was that of four groups:

- William Fairbank, King Walters and their co-workers, who had studied the nuclear susceptibility down to about 0.1 K, found a nearly T^{-1} temperature dependence above 1 K, but an approach to temperature independence at the lowest temperatures. Their results at low pressure were not badly approximated by the susceptibility of an ideal Fermi gas with a degeneracy temperature of 0.45 K, about ten times lower than that expected from an ideal gas of the same number density.
- Douglas Brewer, John Daunt and their colleagues, who had measured the specific heat down to about 0.08 K and found that the slowly changing heat capacity observed above 0.2 K changed at lower temperatures to one that could quite reasonably be said to approach a linear temperature dependence with an effective mass ratio of about two, in quite good agreement with a microscopically based calculation by Keith Brueckner and John Gammel.
- K. N. Zinov'eva, who had measured the viscosity of liquid He^3 down to 0.35 K and found it to rise rapidly with decreasing temperature at the lowest temperatures, but not according to the T^{-2} law expected from Lev Landau's theory of a Fermi liquid, as elaborated by A. A. Abrikosov and I. M. Khalatnikov.[1]
- Richard Garwin and Haskell Reich who, on the other hand, had measured the spin-diffusion coefficient down to about 0.5 K and found it to approach a relatively large but temperature-independent value.

We had been working on a nonresonant ringing method for measuring nuclear spin temperatures in metals at the low ambient magnetic fields that accompany temperatures achieved by electronic magnetic cooling. Although the first measurements Howard Hart and I made on liquid He^3 in the fall of 1958 utilized this method, we soon abandoned it in favor of Erwin Hahn's spin-echo method as developed by H. Y. Carr and Edwin Purcell, because with it the spin-diffusion coefficient as well as the nuclear spin susceptibility could be measured. Insisting that the spin-echo measurements be made in a field of about 20 gauss to facilitate obtaining low temperatures, Hart and I began our spin-diffusion work, with these results:[2] In the vicinity of the lowest temperatures reached by Garwin and Reich the spin-diffusion coefficient begins to rise, and by our then-lowest temperature, 0.07 K, has increased by a factor of ten.

Many liquids have a viscosity that increases with decreasing temperature (oil, for instance), but the large and rapidly rising spin-diffusion coefficient we observed, taken together with the other results mentioned, suggests that at low enough temperature He^3 may well be a Fermi liquid. This observation and the techniques on which it was based formed the cornerstone for our later work.

Pure helium 3 as a Fermi liquid

In subsequent years we learned a great deal more about He^3 as a possible normal Fermi liquid. With Bill Abel, Ansel Anderson, Bill Black, John Connolly, Moyses Kuchnir, Bill Reese, Gerhard Salinger, Ray Sarwinski and Bill Steyert, I studied a variety of properties of pure He^3 at progressively lower temperatures during this period. Additional motivation to lower the temperature was meanwhile provided by the theoretical predictions of a pairing transition to a superfluid state, such as that in Bardeen–Cooper–Schrieffer theory, by L P. Pitaevskii, by Brueckner, Toshio Soda, Philip Anderson and Pierre Morel, and by Victor Emery and Andrew Sessler. For a while the predicted transition temperature decreased with

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time, always seeming to be just a bit below the low-temperature limit of the existing experiments.

However, the elusive superfluidity only served to intensify our interest in the properties of the normal liquid. These are perhaps epitomized by experimental data published[3] in 1965 for the spin-diffusion coefficient D and nuclear susceptibility χ of pure liquid He^3 at lower pressure. The lowest temperature was decreased by a factor of twenty from that achieved in our first work in 1959. The temperature scale used, T^*, is based on the susceptibility of powdered cerous magnesium nitrate in the shape of a right circular cylinder with diameter equal to height. Much to our surprise we found that we could obtain a rather consistent picture by setting the absolute temperature T equal to T^* over most of the temperature range. Measurements by noise thermometry made recently suggest that we can take T equal to T^* down to about 5 mK. The observed T^{*-2} dependence of D on temperature is just what would be expected for the mutual scattering of Landau's quasiparticles in the Fermi surface and strongly supports his concept of the Fermi liquid. The essential temperature independence of χ also fits with the picture of a degenerate gas of Fermi quasiparticles.

Experimental data on the specific heat C^*, however, do not lead to such a clear-cut conclusion. Rather than reaching a constant value, the quantity C^*/nRT^* for pure He^3 continues to rise as the temperature drops. In a short but important paper, Philip Anderson called attention to these experimental observations. This led to the concept of He^3 as a strongly interacting Fermi liquid and to the development by Sebastian Doniach and Stanley Engelsberg of the spin-fluctuation model, also known as the paramagnon model. This model has been of great importance not only in describing thermodynamic and transport properties but also in understanding many features of the newly discovered superfluid phases.

Magnetically assisted heat flow

In the course of the above measurements we accidentally found a remarkable property of the thermal contact between pure He^3 and cerous magnesium nitrate, which plays a crucial role in the success of many experiments on pure He^3 below melting pressure. In our specific-heat measurements we added pulses of heat to this mixture and observed the average temperature as a function of time. The measured time constant for equilibrium, rather than increasing rapidly with decreasing temperature as our experience with the Kapitza thermal-boundary resistance suggested, actually reached a maximum and decreased with decreasing temperature, reaching the unbelievably small value of 20 seconds at the lowest temperature.

When the time constant is analyzed in terms of two reservoirs connected by a thermal resistance, we find that the resistance is linear in temperature. Since the thermal resistance of liquid He^3 should be linear in temperature, we thought at first that the He^3 itself was the source of the resistance and that the boundary resistance was orders of magnitude less than we expected. But when we later found that the linear resistance did not occur for dilute solutions of He^3 in He^4 with comparable bulk thermal conductivity, we guessed that a magnetic coupling between the pure He^3 and the magnetic cerous-magnesium-nitrate substrate was responsible for this very effective energy transfer. A short time later Anthony Leggett and Matti Vuorio theoretically explained the linear temperature dependence of the thermal resistance on the basis of magnetic energy coupling.

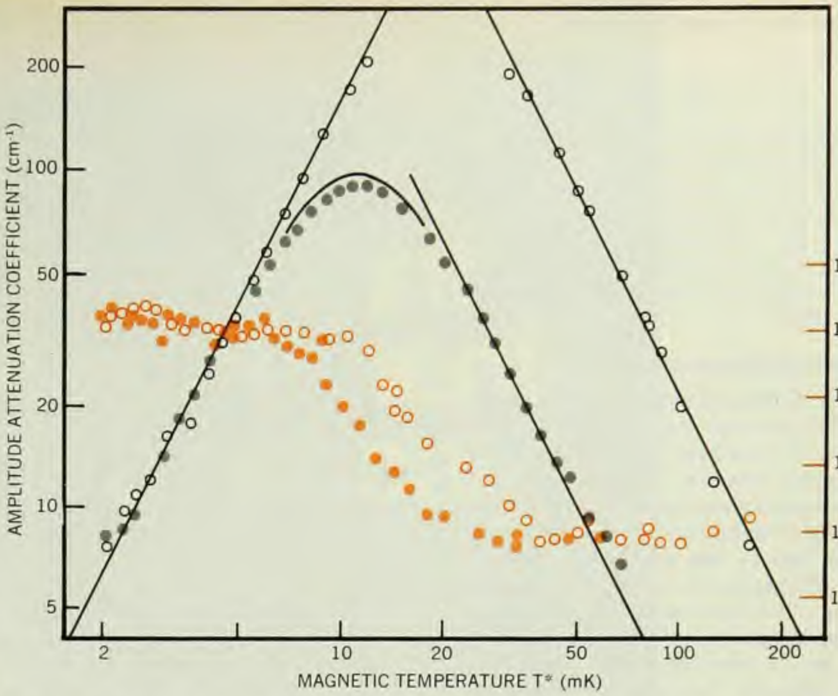
Landau's Fermi liquid consists of a strongly degenerate gas of quasiparticles with a number density equal to that of the liquid. These quasiparticles scatter one another around the Fermi surface with a scattering rate proportional to T^2. We have seen from specific-heat and transport measurements that He^3 fits this picture to a considerable extent.

Another feature of Landau's theory is the concept of the effective interaction between a given quasiparticle excitation and others, described in terms of a set of so-called "Landau parameters." As Leggett describes it,[4] the interaction can be visualized and quantified in terms of a number of molecular fields. These can be classified according to the

A dilution refrigerator with its pumping lines emerging above (left foreground) hangs over an open pit in Wheatley's copper-screened lab. Another one (behind it and to the left) is enclosed in its Dewar. Wheatley's assistants shown in this photo are (left to right) Douglas Paulson, Ronald Sager, Evelin Pichelman, Robert Kleinberg and Matti Krusius. Photograph by Douglas Paulson. Figure 1

experiment and to the deformation of the Fermi surface. For a simple change of density we have a field proportional to the Landau parameter F_0; for a uniform flow we have another field proportional to F_1; for the application of a magnetic field and a spherically deformed Fermi surface we have a molecular magnetic field proportional to Z_0, and so on.

These parameters serve a very useful purpose in describing experimental data on, for example, first-sound velocity, specific heat and nuclear magnetic susceptibility. But perhaps more importantly the possible existence of such effective fields led Landau[5] to predict in 1957 the existence of collisionless sound, referred to as zero sound. In a Fermi liquid at such a low temperature and such a high frequency that collisions would not maintain thermal equilibrium, zero sound would propagate under the action of the effective molecular fields. A transition from first to zero sound in He^3 was first observed by John Wilks and his collaborators in some elegant experiments on acoustic impedance. In 1966 Abel, Anderson and I studied the progation of 15- and 45-MHz ultrasound in pure He^3 at low pressure and found the results on attenuation and velocity shown in figure 2. The frequency-independent, T^2-dependent attenuation predicted for zero sound and the ω^2/T^2 dependence of attenuation for first sound were semi-quantitatively verified, while the observed velocity increase of zero over first sound could be understood in terms of Landau parameters deduced from first-sound velocity and specific heat measurements. We conclude from the results on figure 2 that pure He^3 can be described quite well in terms of Landau's concept of a Fermi liquid.



The transition from zero sound to first sound in He^3 is indicated by these graphs. The amplitude-attenuation coefficient (grey, 15.4 MHz; black, 45.5 MHz) is independent of the frequency for the lower temperatures, as predicted for zero sound. The sound velocity (in color; closed circles, 15.4 MHz; open circles, 45.5 MHz) decreases in the transition. Figure 2

Dilute solutions of He^3 in superfluid He^4

In 1965, near the end of the period I have just described, two events of great scientific and technological importance to the study of He^3 occurred:

- David Edwards analyzed measurements of the phase-separation temperature of dilute solutions of He^3 in superfluid He^4 down to about 0.1 K. He concluded that at the absolute zero about 6% of He^3 will dissolve in He^4. A single He^3 atom is therefore more strongly bound to superfluid He^4 at low temperatures than it is to pure He^3 itself, but a concentration of 6% is sufficient for the Fermi kinetic energy of the degenerate gas to balance this potential-binding-energy difference at absolute zero. This suggested that for concentrations less then 6%, it should be possible to study new strongly degenerate Fermi systems with degeneracy temperatures that can be varied by changing the concentration.
- Following pioneering work at Leiden University, Henry Hall in Manchester described his version of Heinz London's dilution refrigerator, which had reached a temperature of 65 mK. A parallel development by Neganov in the Soviet Union was unknown to us at the time. The implication of Edwards's suggestion for dilution refrigeration was that significant cooling should be possible down to very low temperatures—much lower then the 65 mK reported by Hall.

We were extremely interested in these developments, and a short time later Anderson, Ronald Roach, Sarwinski and I had measured the heat capacity of 1.3% and 5.0% dilute solutions down to below 10 mK and found not only an essentially ideal (though degenerate) gas behavior with specific heat linear in temperature near absolute zero but also no phase separation.

We followed these with magnetic measurements by the spin-echo method and additional thermal conductivity measurements to which Abel, David Edwards, Richard Johnson and Bill Zimmerman also contributed. In obtaining these data we found the thermal contact between cerous magnesium nitrate and dilute solutions to be very much worse than that to pure He^3. The nuclear susceptibility was characterized by that of a nearly ideal gas having weak attractive interactions (a positive Z_0). Spin diffusion measurements of both 5.0% and 1.3% solutions show a diffusion coefficient much greater than that of pure He^3. At the lowest temperatures they conform reasonably well to the T^{*-2} law characteristic of degenerate Fermi gases or liquids, with the 5% solution having a larger diffusion coefficient than the 1.3% solution. The size of the diffusion coefficient at the lowest temperature for the 5% solution is staggering: over 20 cm^2/sec, nearly 10^6 times greater than the smallest diffusion coefficient observed in pure liquid He^3 at low pressure.

Thermal conductivity measurements indicated competition between He^3 and He^4 excitations in heat conduction at higher temperatures, but were consistent with the T^{-1} dependence expected for the conductivity of a degenerate Fermi gas at the lowest temperatures.

As our experimental measurements unfolded, John Bardeen, Gordon Baym and David Pines developed a theory of the effective interactions between the He^3 quasiparticles in dilute solutions, which has been very important in understanding the various observed phenomena. Interactions between the particles may be understood in terms of the scattering of particles near the Fermi surface: The magnitude and direction of the momentum transfer **q** can be quite different for direct and exchange scattering. Because its magnitude q varies from zero to twice the Fermi momentum and the latter is a function of the He^3 concentration, the dependence of the effective interaction $V(q)$ on q can in principle be deduced from measurements at a variety of concentrations. This interaction is attractive for small q and must increase to account for the higher transport in the more concentrated solutions. Bardeen, Baym and Pines also estimated the transition temperature to a BCS-like paired s-wave superfluid state by the formula

$$T_c \simeq T_f \exp|N(0)V|^{-1}$$

The quantity $N(0)\langle V \rangle$ estimated from the measurements became maximally negative at a concentration of 1.6%, leading to a critical temperature T_c of about 2 μK. At least at this stage the hope of observing superfluid Fermi liquid in a superfluid Bose liquid appeared pretty dim.

Technological innovations

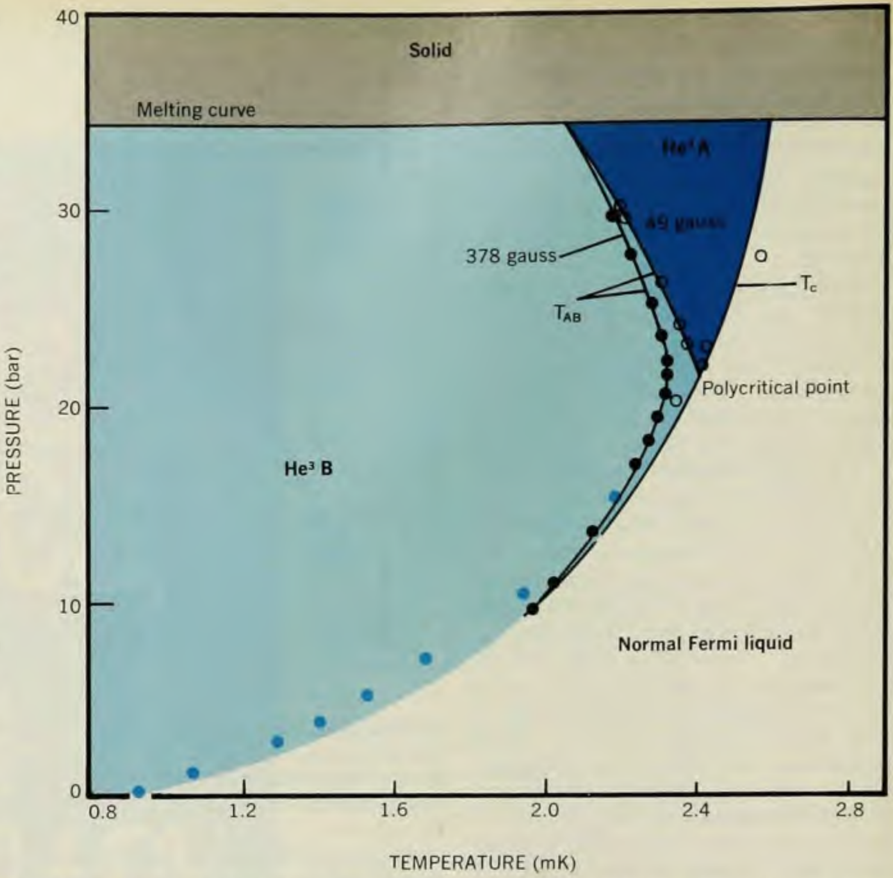
With all the detailed scientific knowledge that we now had about pure He^3 and dilute solutions, new technological innovations were not difficult to make. Great progress was made in 1966 in developing an improved dilution refrigerator to replace the potassium chrome alum refrigerator used as a precooling means. The alum magnetic refrigerator was located above the experimental cells, was of substantial size, required a large magnetic field and could not be maintained cold for very long periods of time. The dilution refrigerator changed all that. The temperatures Oscar Vilches and I achieved were well below 10 mK—very favorable for the development of the type of refrigerator that is common today.

Another type of nonmagnetic cooling was pioneered by Y. D. Anufriyev in 1965, who first found the cooling effect predicted by I. Pomeranchuk to accompany the adiabatic conversion of liquid to solid He^3. In Anufriyev's experiments an indicated low temperature of 20 mK was reached. However, if the frictional heating caused by the mechanical motion needed to reduce refrigerator volume could be eliminated, the resultant cooling power at low temperatures should be greater than that possible with the dilution refrigerator. A thermodynamic argument shows that the refrigeration available on adiabatic conversion of liquid to solid He^3 along the melting curve is linear is the temperature, at least above about 10 mK.

In 1968 Johnson and I built an improved compressional cooling device based on this method. Working with Ralph Rosenbaum and Orest Symko, we found that it worked beautifully, with scarcely a trace of frictional heating. We produced temperatures in the low-millikelvin region, which at that time were the lowest temperatures ever achieved by purely mechanical means. This refrigerator and others like it have also been used to study properties of solid He^3, but that is beyond the scope of this article.

The new phases

Following the explosion of refrigeration and measurement technology in the late 1960's and early 1970's there was a qualitative change in our capability to perform experiments at millidegree temperatures and below. The new technology played a crucial role in subsequent developments. The most spectacular of these was the discovery by Douglas Osheroff, Robert Richardson and David Lee of the "A" and "B" features on the pressurization curve of He^3 in an adiabatic compressional cooling cell. In a conference address three years ago, Lee described the experiments at Cornell to uncover the mysteries behind these features. We were fascinated by the nuclear-magnetic-resonance properties, which showed that these A and B features had revealed transitions to new liquid phases. Reference 7 reviews the experimental work up to the fall of 1974 and reference 4 is an excellent theoretical review.



The phase diagram of He^3 in the low-millikelvin region. Below the solid phase (shown in grey) are three liquid phases. These are, for decreasing temperatures, the normal Fermi liquid (white), the A phase (dark color) and the B phase (light color). The transition temperature between the latter two, T_{AB}, is strongly affected by the magnetic field; data for 49 gauss and 378 gauss are shown. The line of second-order transitions T_c, between He^3 A and the normal Fermi liquid, intersects the T_{AB} line at the so-called "polycritical point." Figure 3

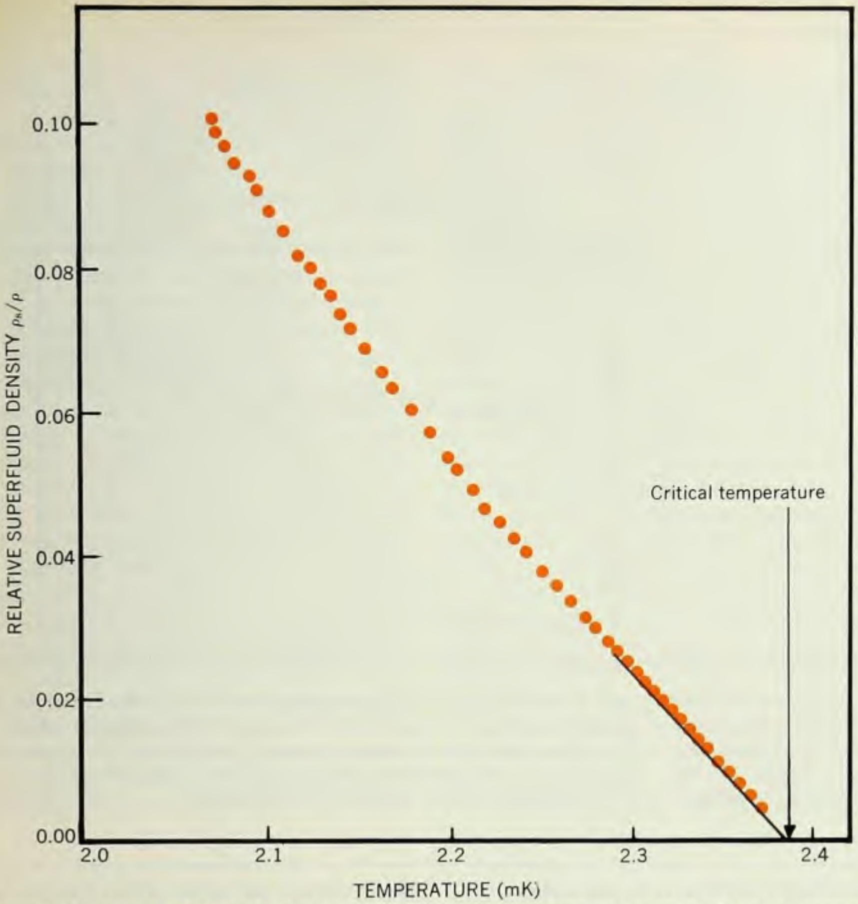
The phase diagram in the region of the new phases is shown in figure 3. There is a line of Ehrenfest-type second-order transitions called T_c separating what is called He^3 A from normal Fermi liquid. A line of first-order transitions T_{AB} at a lower temperature separates He^3 A from He^3 B. In zero magnetic field the line T_{AB} appears to intercept the line T_c at a point known as the polycritical point, but which has never been precisely identified. Relatively small magnetic fields have a profound effect on the line T_{AB} near the polycritical point, and below it have the effect of interposing a thin slice of He^3 A between He^3 B and the normal fluid.

From the probable existence of the polycritical point and the profound effect of the magnetic field we can conclude qualitatively that the thermal properties of the A and B phases near T_c are very similar, while other properties, such as magnetic susceptibility, are very different. In a magnetic field the line T_c is also split into two second-order transitions at T_{c1} and T_{c2}, as Wilfrid Gully, Osheroff, Dewey Lawson, Richardson and Lee showed. The splitting is small and essentially linearly proportional to the field H:

$$T_{c1} - T_{c2} = \frac{6\mu K}{kG} H$$

Between T_{c1} and T_{c2} the fluid is called He^3 A_1. As far as we know now, there are thus three bulk liquid phases in addition to the normal liquid: He^3 A_1, He^3 A and He^3 B.

The form that the second-order transition took in our first experiments with Richard Webb, Thomas Greytak and Johnson shows that the transition at T_c is sharp, reflecting a discontinuity but



Superfluid density relative to total density as a function of temperature for He^3 B at a pressure of 20.7 bar. The density is measured by fourth-sound transmission through a superleak of 50-micron parallel plates. BCS-type theories predict such a linear dependence. Figure 4

not a divergence of the specific heat at T_c. The specific-heat discontinuity was measured from roughly melting pressure to about 23 bar, and was always larger than expected for a weak-coupling BCS-like transition. The T_{AB} line was first found in heat flow studies with Greytak, Johnson and Douglas Paulson, and was confirmed by measurements of static magnetism. As the temperature decreases through T_c, the magnetization increases slightly and then becomes constant within the A phase. The slight increase is presumed to occur as the He^3 cools through the A_1 phase. Thus He^3 A_1 and He^3 A are about equally magnetic with the normal liquid. But as the liquid cools through the temperature T_{AB} into the B phase the magnetism drops suddenly to a value, relative to the nuclear paramagnetism of the normal state, that is determined mainly by the ratio T_{AB}/T_c.

It is significant that the relative magnetization of He^3 B fits a nearly universal curve when plotted against T/T_c, where T_c is the transition temperature from He^3 A to normal liquid. This is one basis for suggesting that B and A fluids have the same T_c. At much lower temperatures, measurements in the B phase by Osheroff at melting pressure and by Antti Ahonen, Mikko Haikala, Matti Krusius and Olli Lounasmaa at lower pressures show that the resonant magnetism has become nearly independent of temperature at a value about 0.3 of the nuclear paramagnetism in the normal state.

The B-phase *static* magnetism decreases at pressures near the polycritical point about twice as rapidly with decreasing temperature as the *resonant* magnetism measured by Krusius's group (and with a qualitatively similar comparison with Osheroff's result at melting pressure). Further experimental work is needed to resolve this discrepancy.

Before leaving the phase diagram I want to point out a remarkable property of the second-order line. If the transition at T_c is to a BCS-like pairing state, we might expect that an approximate formula for the transition temperature would be

$$T_c \simeq T_{\text{eff}} \exp\left[-\frac{1}{N(0)V_{\text{eff}}}\right]$$

It might be appropriate to set T_{eff} equal to the "magnetic" Fermi temperature, $(3/2)T^*$, where T^* is the temperature needed to calculate the low-temperature paramagnetic susceptibility of the normal liquid with Curie's law. If we use this formula together with $N(0)$ and T_c at melting pressure to deduce V and then *assume* that V is independent of the pressure, we find that the formula predicts a T_c of about 1 mK at 21 bar and of 14 μK at zero pressure, the latter over sixty times lower than the observed T_c. The effective interaction V_{eff} evidently depends on pressure. If we adopt a purely empirical approach and calculate $N(0)V_{\text{eff}}$ from

$$N(0)V_{\text{eff}} = \left[\ln\left(\frac{3}{2}T^*/T_c\right)\right]^{-1}$$

as a function of pressure, we find that it increases from zero to melting pressure by 37%.

Motivated by a suggestion of Leggett, we noted that

$$\frac{|Z_0|}{1 + Z_0/4}$$

increases by about the same factor. An experimental comparison of these factors indicates that T_c can be calculated with reasonable accuracy with only Fermi-liquid parameters.

Magnetic superfluids

We can now answer the question whether or not the new phases are superfluid. Experiments have been performed in Helsinki on the resonant behavior of a wire vibrating in the liquid, in La Jolla on the hydrodynamic flow of heat—experiments somewhat analogous to the early experiments of Jack Allen on He^4 and in both La Jolla and Cornell on fourth sound, the propagation of sound through a superleak. In the case of fourth sound there is no signal at all unless there is a superfluid present. We now know that fourth sound propagates through both A and B phases. Quantitative observations of fourth-sound velocities C_4 yield the relative superfluid density by the formula

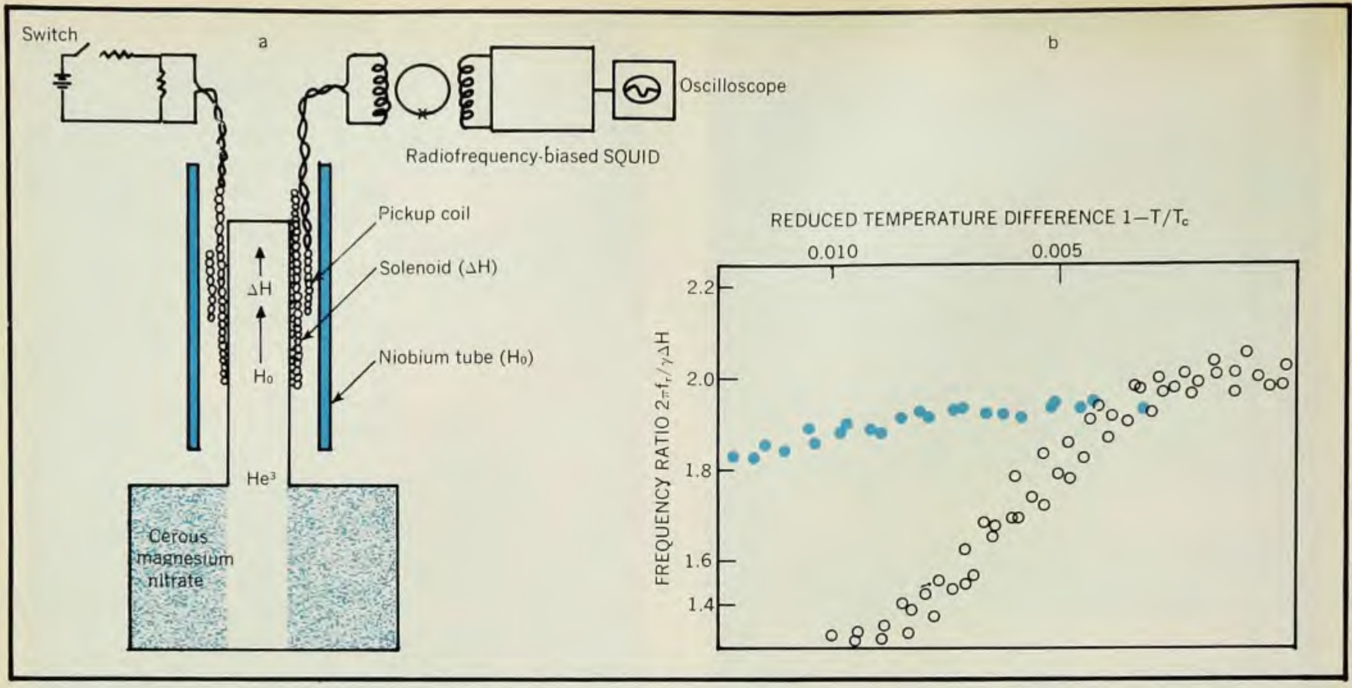
$$\rho_s/\rho = C_4^2/C_1^2$$

where C_1 is the first-sound velocity. In figure 4 I show some results for superfluid density obtained with Haruo Kojima and Paulson in this way using as a "superleak" a stack of parallel plates with a 50-micron spacing between plates. Such a large spacing is possible because the normal viscosity of the liquid is in excess of 0.1 poise, as opposed to the 10–20 micropoise that characterizes the normal viscosity of superfluid He^4.

I particularly draw your attention to the temperature dependence of ρ_s/ρ near the critical temperature:

$$\rho_s/\rho \propto 1 - \frac{T}{T_c}$$

This is the temperature dependence expected for the square of the order parameter in BCS-like theories of the superfluid state. The same temperature dependence is observed for several other quantities that can be measured precisely. The normal viscosity of He^3 below T_c is also very interesting, show-



Josephson-effect measurements on the superfluid phases of He^3 support their explanation in terms of weakly coupled interpenetrating superfluids. The schematic shows the superconducting niobium tube that provides a constant trapped magnetic field, the coaxial ringing and pickup coils and the CMN refrigeration stage; the dilution refrigerator is not shown. The graphs show the ratio of the driven-mode ringing frequency to $\gamma\Delta H/2\pi$ as a function of the reduced temperature difference for a rectangular cavity. The colored circles are for He^3 A at 22.0 bar; the black, He^3 B at 20.7 bar. The ratio of two near the critical temperature indicates tunnelling by He^3. Figure 5

ing a precipitous drop with a temperature dependence of $[1 - T/T_c]^{1/2}$.

We have seen that the new phases are magnetic superfluids with a superfluid density near T_c that varies with the BCS-like temperature dependence. Let us discuss in more detail some questions concerning the BCS pairing hypothesis for superfluid He^3. In superconductors the electrons pair in nonmagnetic singlet $\uparrow\downarrow$ states, and the corresponding superfluid is nonmagnetic with respect to spins. But in He^3 the experimental data suggest that pairing is in triplet states, in which case we can imagine, quantizing along a field, that we can have magnetic $\uparrow\uparrow$ pairs, magnetic $\downarrow\downarrow$ pairs and nonmagnetic triplet $\uparrow\downarrow$ pairs. Superfluids formed with either $\uparrow\uparrow$ or $\downarrow\downarrow$ correlated pairs would be about equally magnetic with the normal Fermi liquid while the superfluid component for $\uparrow\downarrow$ pairs would be nonmagnetic.

There are, then, three possible spin arrangements of the correlated pairs and, interestingly enough, three known new phases: the A_1 phase, the A phase and the B phase. As I pointed out in discussion of the phase diagram, both the A_1 and the A phases are about equally magnetic with the normal liquid, while the magnetism of the B phase decreases below T_c, the resonant magnetism becoming nearly temperature independent at about 0.3 the normal-liquid value as T approaches zero. This immediately suggests that the A_1 and A phases have no $\uparrow\downarrow$ pairs while the B phase contains some nonmagnetic $\uparrow\downarrow$ pairs—but not all, to account for the nonzero magnetism as $T \to 0$.

We can go further regarding the A_1 phase. It was suggested theoretically by Vinay Ambegaokar and David Mermin, and then proved most conclusively in experiments by Osheroff and Anderson, that the A_1 phase contains only one magnetic superfluid, either $\uparrow\uparrow$ or $\downarrow\downarrow$ although it is not yet known which. Thus, in addition to the quasiparticle excitations of the normal fluid we imagine that the A_1 phase has a one-component ($\uparrow\uparrow$ or $\downarrow\downarrow$) superfluid, that the A phase has two magnetic interpenetrating superfluids—one being the $\uparrow\uparrow$ fluid and the other the $\downarrow\downarrow$ fluid, and the B phase has three interpenetrating superfluids—magnetic $\uparrow\uparrow$ and $\downarrow\downarrow$ fluids and nonmagnetic $\uparrow\downarrow$ fluid. As emphasized by Leggett and by Kazumi Maki and his collaborators, these superfluids cannot act independently since they are weakly coupled to one another by a coherent nuclear dipole–dipole interaction.

Ringing

The idea of interpenetrating, weakly coupled superfluids of correlated pairs can be experimentally tested, and here I will refer to some recent experiments I have done with Richard Webb and Ronald Sager. Let me begin by reminding you of Josephson's effect for two weakly coupled superconductors. If I apply a voltage difference V across a suitable junction, then the phase of the superconducting electron state on one side of the junction advances with respect to that on the other at the rate $2\Delta\mu/h$, where $\Delta\mu = eV$ is the change in chemical potential experienced by a single electron on crossing the barrier, and the 2 comes in because what tunnels across the barrier is a pair. This is the same factor of 2 that appears in the expression $h/2e$ for the flux quantum. There result ac currents at frequency $2eV/h$ with a strength depending on how the energy coupling across the barrier depends on phase.

In He^3 A we imagine that there are interpenetrating $\uparrow\uparrow$ and $\downarrow\downarrow$ fluids homogeneously coupled by the coherent dipolar interaction. The superfluids are not separated in space as for the customary Josephson effect. Furthermore, the coupling between the superfluids can not be manipulated—except to some extent with pressure and magnetic field. It does not appear to be possible to establish a chemical-potential difference between the two superfluids that is both homogeneous and essentially constant in time. However, if we can do our experiments in a time short compared with the time scale on which the chemical potential does change, then perhaps we can observe Josephson supercurrents in superfluid He^3 at constant $\Delta\mu$ and test the pairing hypothesis, as Leggett recently suggested. Note that the analog of a supercurrent flow across a barrier is the transfer of pairs from, say, the $\downarrow\downarrow$ to the $\uparrow\uparrow$ fluids, corresponding to a change of magnetization. The analog of an ac supercurrent is thus an oscillating magnetism.

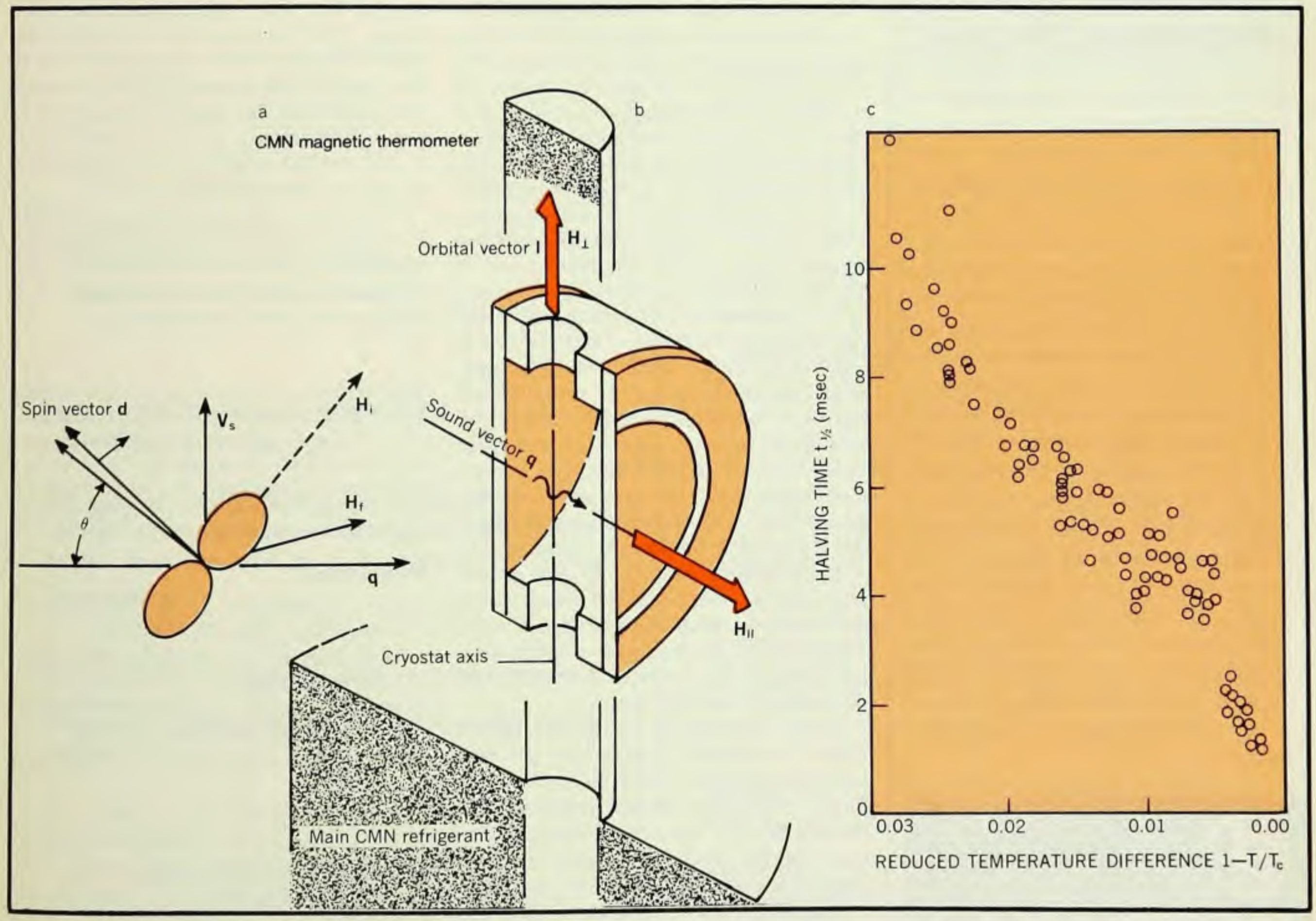
Experiments that we have arranged to test these ideas are based on the Varian–Packard field turn-off method but, taking a cue from Leggett's prediction of parallel (or longitudinal) NMR, with the flux pickup coils coaxial with the field that is changed. A schematic diagram of the CMN-refrigerated He^3 method that we used is shown in figure 5a. A tower above the refrigerant contains He^3, which can be subject to both a constant field $\mathbf{H}_0$ trapped in a superconducting niobium tube and an incremental field $\mathbf{\Delta H}$ that can be changed rapidly. The magnetization of the He^3 parallel to $\mathbf{\Delta H}$ is sensed by pickup coils and a broad-band detector. If the field $\mathbf{\Delta H}$ is suddenly applied to the He^3, the chemical potential of ↑ spins goes down by $\frac{1}{2}\gamma\hbar\Delta H$ (we take ↑ along the magnetic moment, not the spin) while that of ↓ spins goes up by $\frac{1}{2}\gamma\hbar\Delta H$, the difference being $\gamma\hbar\Delta H$. Thus, the chemical potential difference for transferring a *pair* is $2\gamma\hbar\Delta H$. The corresponding angular frequency is $2\gamma\Delta H$. If we are dealing with He^3 B, the chemical potential of the ↑↓ fluid is not changed by the field, so we expect both $\gamma\Delta H$ and $2\gamma\Delta H$ frequencies to be present. Note in particular that these frequencies do not depend on temperature.

As a practical matter it is now known that when a small enough field change ΔH is applied there is a subsequent parallel ringing of the magnetization under the action of the coherent dipolar coupling. The square of this ringing frequency is proportional to $1 - T/T_c$, like the superfluid density, so we can make it quite small by letting T approach T_c. What happens in this temperature-dependent ringing is that, after the field change ΔH, pairs are transferred from one superfluid to another—rather analogously to the so-called "Josephson plasma oscillation." We conclude that the chemical potential difference oscillates at this temperature-dependent frequency. But since the frequency $2\gamma\Delta H$ is independent of temperature we can in principle make it as large as we please with respect to the frequency that characterizes chemical-potential changes induced by dipolar coupling by making observations close to T_c.

In the experiments we made ΔH as large as possible consistent with experimental limitations and then observed the ringing frequency following sudden field turnoff ΔH. I show in figure 5b the results obtained with Webb and Sager for a rectangular cavity 1.0 mm across with the field parallel to the boundaries. As T approaches T_c, the ringing frequency does approach twice $\gamma\Delta H/2\pi$, within our present experimental accuracy, for both A and B fluids. Similar results were found with a stack of parallel plates of 0.5-mm separation with the field perpendicular to the boundaries.

Anisotropy

We have spent some time discussing interesting properties of the magnetic superfluids. Another property of these remarkable fluids is their anisotropy, and in these concluding paragraphs I will concentrate on orbital and mechanical as opposed to spin anisotropy. We will understand the word "anisotropy" to imply that the momentum density corresponding to a given normal fluid or superfluid velocity depends on the orientation of some internal variable of the fluid with respect to that velocity. The superfluid density therefore might be an anisotropic tensor. Indeed, this is expected for He^3 A, but no one has ever done an experiment in which the com-



Geometry, apparatus and results of an orbital anisotropy experiment. The helium is subjected to ultrasound and a magnetic field that is suddenly changed. Diagram **a** shows the orientations of the relevant vectors, including flow field $\mathbf{v}_s$ and sound vector $\mathbf{q}$. Diagram **b** is a view of the zero-sound cell used to study orbital dynamics in He^3 A. Graph **c** shows the time for half the attenuation change to take place after the perpendicular field changes from 9.3 G to 2.6 G, with a parallel field of 8.6 G and at a pressure of 23.0 bar. Figure 6 ponents of its density tensor have been measured. Although we do not know *how* anisotropic He^3 A is, we do know *that* it is anisotropic, particularly on the basis of experiments performed at Cornell University by Lawson, Hans Bozler and Lee and at the Argonne Laboratory by Pat Roach, Bernard Abraham, Paul Roach and John Ketterson. They observed anisotropy of the propagation of collisionless sound in He^3 A with respect to the angle between a magnetic field and the wave vector of the sound.

One of the first phenomena studied in the new phases of He^3 was the propagation of collisionless—or zero—sound. It was found that, starting at T_c, the attenuation increased dramatically as the temperature dropped, by a greater amount the higher the frequency, reached a peak, and then decreased to frequency-independent values. This peak is interpreted to represent absorption of sound into collective modes of the order parameter. In an axial p-state model for He^3 A, the absolute square of the gap has the axially symmetric two-lobed form shown in polar coordinates in figure 6a, the unit vector **l** representing this symmetry axis.

Coupling of sound with wave vector **q** to collective modes depends on the relative orientation of **l** and **q**. If at constant temperature the orientation of **l** with respect to **q** can be made to change, the attenuation should also change. At both Cornell and Argonne a dependence of attenuation on the relative orientation of **q** and a magnetic field **H** was observed. Sound attenuation is a probe of the orientation of the orbital state of He^3 A, while He^3 B shows no such anisotropy.

It is known that both a magnetic field and a velocity field can orient He^3 A. A superfluid velocity field $\mathbf{v}_s$ orients **l** parallel to $\mathbf{v}_s$ with an energy proportional to $(\mathbf{l}\cdot\mathbf{v}_s)^2$ owing to the anisotropy of the superfluid density tensor. A magnetic field does not orient **l** directly. . I mention here that the spin order is described in part in terms of a unit vector **d**. A magnetic field tends to orient **d** in a plane perpendicular to **H** with an energy proportional to $(\mathbf{d}\cdot\mathbf{H})^2$ owing to the susceptibility anisotropy of the A fluid. Then the coherent dipolar energy, proportional to $(\mathbf{l}\cdot\mathbf{d})^2$, tends to orient **l** and **d** parallel to one another.

In the diagram in figure 6a some of these vectors are displayed for the case in which a weak $\mathbf{v}_s$ and **H** lie in the same plane. The angle θ that **l** makes with respect to **q** is the complement of the angle made by the field with **q** under field-dominated equilibrium conditions.

This suggests an experiment to study the dynamics of the orbital superfluid state of He^3 A, a subject about which very little is now known. In an exploratory experiment Paulson, Robert Kleinberg and I have arranged a zero-sound cell as shown in figure 6b to use the dependence of ultrasound attenuation on the orientation of **l** as a probe to study the motion of **l** following a sudden change of **H**.

The He^3 studied lies between two circular quartz transducers and is subjected to a field $\mathbf{H}_\parallel$ parallel to **q** and another field $\mathbf{H}_\perp$ which is perpendicular to **q** and along what is probably the direction of a heat-flow-induced superfluid velocity field $\mathbf{v}_s$. In typical experiments $\mathbf{H}_\perp$ is rapidly changed in such a way as to take advantage of a range of angle θ for which attenuation and θ are nearly linearly proportional.

The action is illustrated by the diagram in figure 6a. The field is suddenly changed from $\mathbf{H}_i$ to $\mathbf{H}_f$. For the field range possible in these experiments the **l·d** coupling is much stronger than the **d·H** coupling, so after an initial transient **l** and **d** are probably nearly parallel. The **d·H** coupling then acts to orient **d**, and thus **l** via the **l·d** coupling, perpendicular to the final field $\mathbf{H}_f$. For the angle range used, the orientation of **l** changes approximately exponentially with time, at least if the temperature is not too close to T_c. Figure 6c shows, as a function of $(1 - T/T_c)$, the time $t_{1/2}$ needed for half the change to take place. We find that the times are quite short and that they decrease toward zero as T approaches T_c. Increasing the field does shorten the time at constant temperature, so in the vicinity of T_c the orbital response time appears to be getting uncomfortably short.

These are just preliminary results. Nevertheless I present them here to emphasize my belief that helium 3 will continue to find ways to surprise and challenge us in the years ahead.

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