Probing the *A*-*B* Phase Interface in Superfluid ^{3}He by Andreev Reflection of a Quasiparticle Beam

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Using a ballistic excitation beam generated by a quasiparticle blackbody radiator, we have probed the gap structure of a *B*-*A* interface in superfluid ^{3}He, stabilized magnetically. We see Andreev reflection both from the interface and from the field-distorted *B*-phase gap. As Andreev processes return excitations to the radiator enclosure, the fraction reflected governs the radiator temperature, from which we infer the maximum spin-dependent quasiparticle gap as a function of magnetic field. These measurements are the first to probe the superfluid ^{3}He phase interface with quasiparticles. [S0031-9007(96)02033-9]

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Superfluid ^{3}He shows the highest level of broken symmetry of any ordered condensed matter system. Consequently the superfluid may exist in several phases with different internal structures. The two most common, the *A* phase and the *B* phase, have very different order parameters and gap structures. Since these phases may coexist, their phase boundary provides a remarkably unique interface. This is the only such high-symmetry structure between two highly ordered but different Bose condensates to which we have access. Despite the inherent interest, practical difficulties in studying this structure have limited experimental study. The interfacial energy was measured long ago [1]. Various aspects of the interface propagation have been studied [2], but very little else.

In this Letter we report the first results of an experiment where we exploit the properties of Andreev reflection to measure the transmission across the *B*-*A* phase interface of a quasiparticle beam. The difference in the gap structures of the two phases ensures that at low temperatures ($kT \ll \Delta$) the available quasiparticle states in the two phases are very different, making quasiparticles ideal probes for observing the interface. We observe very clearly the sudden jump in the gap when the beam is blocked by a region of the *A* phase via the attendant spectacular fall in the quasiparticle transmission. Further, to demonstrate the quantitative as well as the qualitative possibilities of this method we have also measured the maximum quasiparticle gap as a function of magnetic field, both in the *A* phase and in the *B* phase, which we find to be consistent with currently accepted values. At the temperatures of the experiment ($\approx 120\ \mu$K) the energy flux emitted in the beam is very sensitive to changes in the excitation gap, a 5% increase in gap yielding a factor of 2 reduction in flux.

By directing a thermal beam of excitations through a region of changing gap we can infer the maximum gap height along the path from the fraction Andreev reflected, provided that the excitation mean free paths are long compared with the experimental dimension. Excitations with energies greater than the largest effective gap along the beam path can pass, whereas excitations with smaller energies cannot pass and must be reflected by Andreev processes, which accurately return them back along their previous trajectory.

To create a static *B*-*A* interface we apply a magnetic field to a small region of the superfluid, since above some critical field (0.339 T for 0 bar [3]) the *A* phase is preferred. The two superfluid phases respond to a magnetic field in different ways. The gap in the pseudoisotropic *B* phase is distorted by the field, decreasing along the axis parallel to the field and increasing in the perpendicular directions. The parallel gap falls with increasing field [4,5] approximately as $\Delta_{B0}\sqrt{1 - 1/2(\gamma_L B/\Delta_{B0})^2}$, where γ_L is the Fermi-liquid corrected value of the gyromagnetic ratio γ. For quasiparticles with momenta along the field the quasiparticle minimum energy is further split by $\pm\gamma_L B/2$ depending on the quasiparticle spin. Conversely, in the *A* phase the order parameter is inherently anisotropic. The gap has nodes along the axis of the Cooper pair orbital angular momentum vector. Apart from creating a small difference in gap between the up-spin and down-spin pairs, an external field has almost no effect in the *A* phase other than to align the nodal direction perpendicular to the field. There is no spin splitting of the excitation spectrum.

To probe the maximum gap we direct an excitation beam through the region exposed to the magnetic field. The principle of the experiment is illustrated in Fig. 1. A thermal beam of excitations is directed at the *B*-*A* interface [Fig. 1(a)]. Only those excitations with energies above the maximum gap can proceed. The rest are Andreev reflected. If we can measure the energy flux transmitted and we know the beam temperature then we can infer the maximum gap "barrier." For an interface stabilized by a parallel magnetic field the situation is more complicated as the *B*-phase gap distortion depends on the excitation spin [Figs. 1(b), 1(c)].

The experimental arrangement is shown in the inset of Fig. 2. The beam is created by a quasiparticle black-

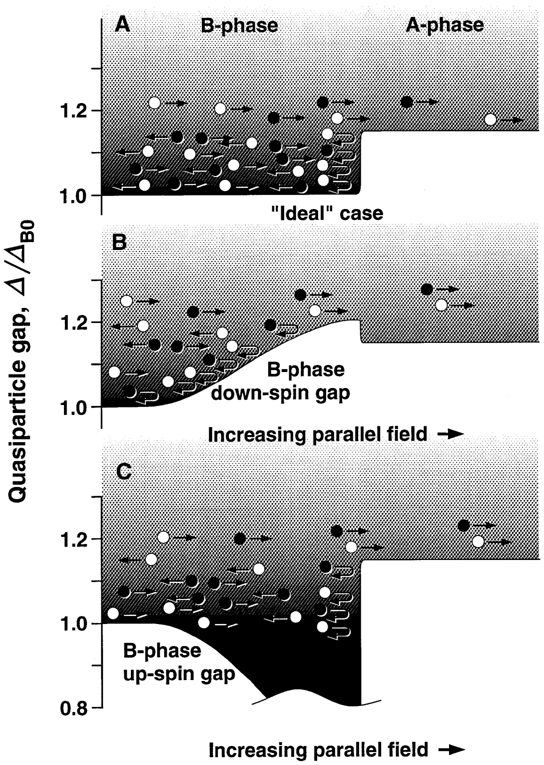

FIG. 1. Interaction of a beam of quasiparticles incident on a *B*-*A* phase interface from the *B* phase. (a) The "ideal" case, both the *B*-phase and *A*-phase gaps are constant. Excitations with energies greater than the *A*-phase gap pass. Those with lower energies are Andreev reflected. In practice, when the *A* phase is stabilized by a parallel magnetic field, there is also a *B*-phase gap distortion which depends on the excitation spin. (b) The case for down-spin excitations. The *B*-phase gap rises toward the phase boundary. Low energy down-spin excitations are therefore Andreev reflected from the *B*-phase gap. (c) In the up-spin case, the *B*-phase gap falls towards the phase interface. Up-spin excitations therefore reach the interface unimpeded.

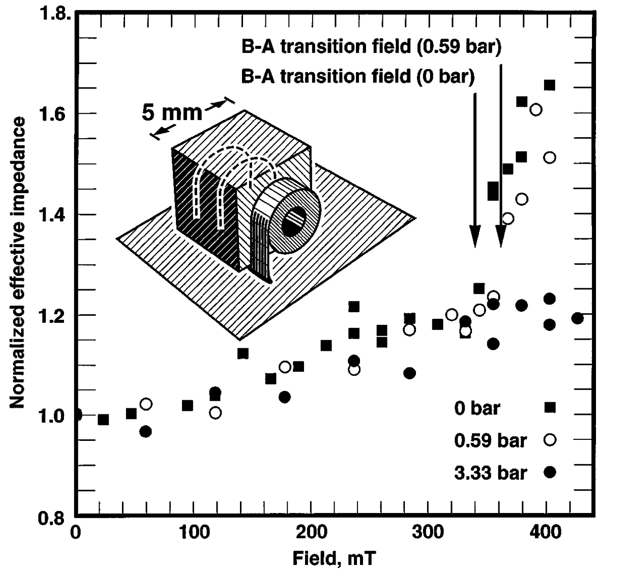

FIG. 2. The measured effective impedance of the blackbody radiator (see text) as a function of magnetic field in the solenoid for 0, 0.59, and 3.33 bar. The vertical arrows indicate the *B*-*A* transition field for the two lower pressures. At the field where the *A* phase first appears the impedance climbs rapidly owing to the sudden increase in Andreev reflection (see text). The inset shows the experimental arrangement of the blackbody radiator box. Around the orifice is situated the miniature superconducting solenoid which produces the magnetic field and is encircled by a silver foil heat sunk to the outer cell.

body radiator [6]. This consists of a 5 mm $\times$ 5 mm $\times$ 5 mm box, made of epoxy-impregnated paper with a $\phi \approx$ 0.3 mm hole in one wall. The box is immersed in superfluid ^{3}He in the inner cell of a double-chambered demagnetization cell and is placed in a void cut in the copper refrigerant in the form of thin sheets coated with sintered silver, a configuration which ensures that the excitation density outside the radiator is negligible compared with that inside. Inside the radiator are two vibrating wire resonators (VWRs) of Nb-Ti filamentary wire [7]. A heater resonator of 13 μm wire, driven hard, heats the liquid in the box by pair breaking, thereby increasing the excitation density/temperature, which is detected by the increase in the damping of the second thermometer resonator of 4.5 μm wire. Two further VWRs outside the radiator monitor the background temperature. The excess quasiparticle density inside the box gives rise to the thermal beam of excitations leaving the hole.

Around the radiator orifice is placed a miniature solenoid with a 1 mm inside diameter. The field direction ensures that the maximum *B*-phase gap distortion and the maximum *A*-phase gap are parallel to the beam direction. The coil is wound with 364 turns of 0.05 mm diameter Nb-Ti filamentary wire. Since it is difficult either to measure or to calculate accurately the field profile of such a small solenoid, we know the central axial field only to within 10%. To produce the large fields needed to stabilize the *A* phase the coil has to carry currents where there is some dissipation. Fortunately the Kapitza resistance between the windings and the liquid is large enough that most of the heat generated can be conducted to the outer cell by a silver heat sink wrapped around the coil. However, there remains a significant heat leak into the inner cell liquid which limits the maximum usable field to about 400 mT.

When the solenoid is energized those thermal excitations in the beam emitted by the radiator which have sufficient energy to escape over the field-generated gap barrier are lost to the bulk liquid outside, whereas the others are Andreev reflected back into the blackbody radiator. To maintain steady state conditions the excitation density inside the radiator adjusts itself so that the energy lost in the beam balances the applied heating. Therefore, for the same heating level, the effect of increasing the barrier is to increase the excitation density inside the radiator.

The measurements are made as follows. A steady heat input is applied to the radiator with the heater VWR. The heat input is calculated from the in-phase product of the ac voltage and current across the heater VWR.

At equilibrium, the frequency width of the thermometer VWR is measured. This is done for a range of heating levels varying typically from 1 to 200 pW. The solenoid is then magnetized and the process repeated.

To put this on a quantitative footing, first consider the case of zero applied field where the gap is everywhere Δ_0, the undistorted B-phase gap. The power entering the radiator comprises the applied power Q_{app} and the heat leak Q_{leak}. At equilibrium these equal the power leaving the radiator in the beam. The beam power Q_{beam} is proportional to the product of quasiparticle number density, energy, and group velocity, i.e., $Q_{\text{beam}} = A\int_{\Delta_0}^{\infty}\exp(-\epsilon/kT)\epsilon\, d\epsilon$, where A is a constant. Therefore we find

$$\begin{aligned} Q_{\text{beam}} &= Q_{\text{app}} + Q_{\text{leak}} \\ &= AkT\exp(-\Delta_0/kT)(\Delta_0 + kT), \end{aligned} \tag{1}$$

where $AkT\exp(-\Delta_0/kT)$ is the number flux leaving the box and $(\Delta_0 + kT)$ is the mean energy per excitation. We measure Q_{app} and the temperature inside the radiator T. We know Δ_0. Therefore by measuring T as a function of applied power we can fit Eq. (1) to deduce the values of the background heat leak Q_{leak} and of the calibration constant A. The data fit Eq. (1) to better than 10% over the whole range of applied power.

Assume we now apply a field somewhere along the beam trajectory and increase the effective gap to a new higher value Δ'. Now only those excitations in the beam with energies greater than Δ' can still escape. The energy lost in the output beam is thus reduced to $AkT\exp(-\Delta'/kT)(\Delta' + kT)$. Since this flux must balance the applied heating, the temperature in the box will rise to a new value T' to compensate. We can think of the hole as having an impedance which increases when the beam trajectory is restricted. We can define a relative impedance I as the ratio of the potential beam energy flux at the higher temperature $AkT'\exp(-\Delta_0/kT')(\Delta_0 + kT')$ to that which we would obtain for the same input power with no beam restriction. That is $I = [T'\exp(-\Delta_0/kT')(\Delta_0 + kT')]/[T\exp(-\Delta_0/kT)(\Delta_0 + kT)]$.

In Fig. 2, we show the measured impedance as a function of the magnetic field for the three pressures 0, 0.59, 3.33 bars. The steady increase in impedance with field at the lower fields arises from Andreev scattering from the gradually increasing down-spin B-phase gap. At a field of approximately 340 mT for 0 bar we see a sharp increase in the impedance, signaling the formation of the A phase within the solenoid. The increase arises since on the formation of the A phase, excitations of both spin species see the maximum A-phase gap barrier, whereas with B phase alone, only the down-spin excitations are restricted by an enhanced gap. We utilize this A-phase signature to provide a more accurate calibration of the solenoid field. The data for 0.59 bar show the sharp increase at a higher field, consistent with the value of 363 mT for B_{AB} given in Ref. [3]. At 3.33 bars we see no sharp rise, since at this pressure the A phase is stable only above 448 mT, beyond the range of our solenoid.

To infer the maximum quasiparticle energy gap $\Delta_{\max}$ as a function of field from the measured impedance we separate the power emitted from the radiator into that carried by the two spin components $Q_\downarrow$ and $Q_\uparrow$. We note that down-spin quasiparticles with $\epsilon < \Delta_{\max}$ must be Andreev reflected and accurately retrace their paths back into the radiator and therefore do not contribute to the beam power. Therefore,

$$Q_\downarrow = A/2kT\exp(-\Delta_{\max\downarrow}/kT)(\Delta_{\max\downarrow} + kT).$$

The up-spin quasiparticles are virtually unaffected by the magnetic field in the B phase, since their effective energy gap is *reduced* by the field. Therefore in the B phase

$$Q_\uparrow(B) = A/2kT\exp(-\Delta_0/kT)(\Delta_0 + kT).$$

However, once the A phase is present, these excitations are also presented with the barrier of the A-phase gap, and therefore above the B-A field we assume that $Q_\downarrow$ corresponds to the maximum B-phase value and that $Q_\uparrow$ is now given by

$$Q_\uparrow(A) = A/2kT\exp(-\Delta_A/kT)(\Delta_A + kT).$$

In an ideal one-dimensional geometry $Q_\uparrow$ and $Q_\downarrow$ comprise the only contributions to the power emitted. However, since the radiator orifice is situated in the fringing field of the solenoid, a small volume of the B phase within the radiator is exposed to the field (the field at the orifice is about half the maximum). Owing to the depression of the up-spin gap, this region is heavily populated with up-spin quasiparticles. In the absence of inelastic processes, these quasiparticles cannot contribute to the beam since they are trapped in the potential well created by the field. However, in practice a large fraction of these excitations can leave the hole and scatter off the inside walls of the solenoid. Once such a quasiparticle has reached the solenoid wall we may assume that there is a significant probability of one or more inelastic collisions with the wall in which the quasiparticle can gain energy of order kT. If, as a result, the energy of the quasiparticle rises above Δ_0 then it can escape into the bulk and will contribute to the power. We must include a contribution from this source $Q_\uparrow^-$ by assuming that a fraction f of the up-spin quasiparticles with energies within kT of Δ_0 are able to contribute, thus yielding

$$\begin{aligned} Q_\uparrow^- = fA/2kT[&\exp(-\epsilon_{\min}/kT)(\epsilon_{\min} + kT) \\ &- \exp(-\Delta_0/kT)(\Delta_0 + kT)], \end{aligned}$$

i.e., a contribution from all such states from $\epsilon_{\min}$ upwards within kT' of Δ_0 where T' is the ^{3}He temperature outside the radiator.

The maximum energy gap seen by the beam $\Delta_{\max}$ can now be calculated. We equate $Q = Q_{\text{app}} + Q_{\text{leak}} = Q_\downarrow + Q_\uparrow + Q_\uparrow^-$ with f as the one free parameter. The value of f is chosen to give the expected value of the maximum gap just below B_{AB} at 0 bar. This yields a value of $f = 0.7$ which we then assume for all three pressures.

The values of $\Delta_{\max}$ inferred as described above are shown in Fig. 3 for three pressures as a function of field, along with the expected gap profile as calculated by Ashida and Nagai [8]. Since the values of $\Delta_{\max}$ depend on whether we assume the beam trajectory is blocked by the *A* phase or only passes through the undistorted *B* phase, we obtain two results for each field. Concentrating first on the data for zero bar, with increasing magnetic field we see a steady increase in the calculated value of the maximum *B*-phase gap (open circles), until at a field of around 340 mT, when the *A* phase blocks the path, the calculation breaks down and the value rapidly rises to infinity. Points calculated on the assumption that the *A* phase is present yield values of Δ_A virtually identical to Δ_{B0} until the transition when they rise to yield, at around 380 mT, a value of Δ_A of about 1.14 Δ_{B0}, slightly lower than the accepted figure. The results for 0.59 bar show very similar behavior except that as expected the jump to the *A* phase appears at higher fields. At 3.33 bars the data clearly show that the *B*-*A* transition is not reached. This is consistent with the expected transition at 448 mT,

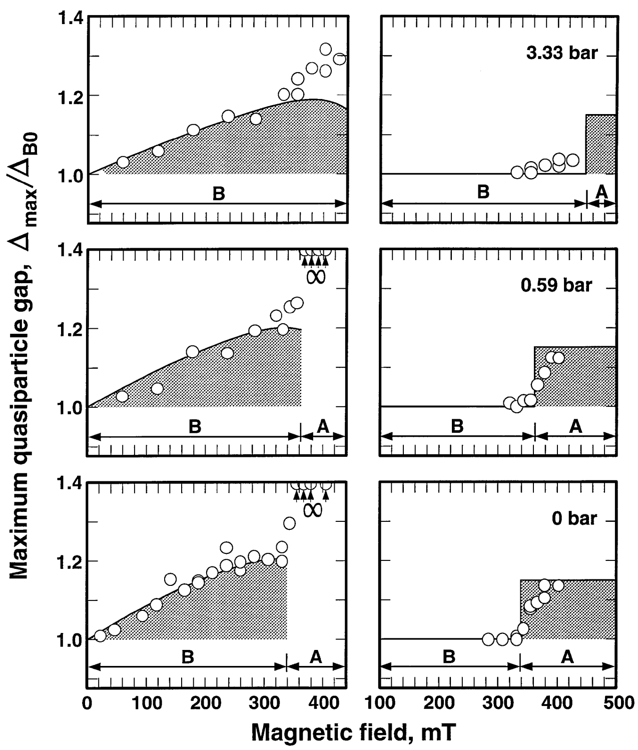

FIG. 3. The effective gap calculated as described in the text. Left hand side figures represent points calculated on the assumption of only the *B* phase present, and right hand figures that the *A* phase has blocked the beam. The calculated effective gaps [8] are also plotted. Above the *B*-*A* transition points calculated on the assumption of only the *B* phase yield infinite values for the down-spin gap.

which is beyond the range of our miniature solenoid. We should stress that while the calculation relies on one free parameter *f*, this is fixed by the data taken at zero bar and the higher pressure values are then consistent with accepted values with no further adjustment.

We have assumed above that any excitation with an energy greater than the barrier can escape freely. However, since the *B*-*A* interfacial width is of the order of the coherence length, there should also be "above-the-barrier" Andreev scattering [9] which would increase the effective barrier height. We see no clear sign of such behavior but since our system only approximates to the model we are not able to make an any more definite statement. To see this more clearly we would need instead to direct the beam from the *A* phase to the *B* phase.

To conclude, we have made the first direct investigation of a static *B*-*A* interface with a thermal beam of quasiparticles. We observe Andreev reflection both from the gap discontinuity at the interface and also from the field-distorted *B*-phase down-spin gap. We also present preliminary measurements of the effective quasiparticle energy barriers. This is the first time that a beam of excitations has been used as a spectroscopic tool in superfluid ^{3}He. More quantitative measurements will follow when we devise better beam collimation. It is also worth noting that, in a field just below that needed to create the *A* phase, the spin splitting means that we can produce beams of excitations with polarizations of 80% or more.

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[1] D. D. Osheroff and M. C. Cross, Phys. Rev. Lett. **38**, 905 (1977).
[2] D. S. Buchanen, G. W. Swift, and J. C. Wheatley, Phys. Rev. Lett. **57**, 341 (1986).
[3] Inseob Hahn, Y. H. Tang, H. M. Bozler, and C. M. Gould, Physica (Amsterdam) **194B–196B**, 815 (1994).
[4] Nils Schopohl, J. Low Temp. Phys. **49**, 347 (1982).
[5] R. Movshovich, N. Kim, and D. M. Lee, Phys. Rev. Lett. **64**, 431 (1990).
[6] S. N. Fisher, A. M. Guénault, C. J. Kennedy, and G. R. Pickett, Phys. Rev. Lett. **69**, 1073 (1992).
[7] C. A. M. Castelijns, K. F. Coates, A. M. Guénault, S. G. Mussett, and G. R. Pickett, Phys. Rev. Lett. **55**, 2021 (1985).
[8] Masami Ashida and Katsuhiko Nagai, Prog. Theor. Phys. **74**, 949 (1985); (calculated values of $\gamma_L(B)$ were provided in a private communication).
[9] N. Schopohl and D. Waxman, Phys. Rev. Lett. **63**, 1696 (1989).