torus, and the beam rotates around it with a progressively smaller radius.

This paper reports on the injection and trapping of an ultrahigh-current electron beam into a modified betatron accelerator. For computational convenience the various parameters in the simulation were such that the induced fields dominate the external fields. This parameter regime should be avoided in an actual device because during the acceleration it is possible that the low frequency of rotation (bounce) will change sign and thus the beam will become unstable. However, the main features of the proposed injection are not sensitive to the relative magnitude of the fields but rather to the magnitude of their difference. Both analysis and simulation are based on the cold-beam approximation. Presently, work is in progress with finite-emittance beams. Finally, it has been assumed that a hard vacuum is continuously maintained inside the confining chamber and thus the plasma formation and its effect on the beam[10] were neglected.

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Sound Attenuation Measurements in Superfluid ^{3}He-A Well below T_c: An Anomalous Behavior at Low Pressure

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Sound attenuation has been measured at low and medium pressures down to $T/T_c \sim 0.4$ in the A phase of superfluid ^{3}He stabilized by a magnetic field. Damping is due to Cooper pair breaking and constitutes a useful probe of the superfluid gap. The gap parameter is found to correspond to the weak-coupling value at high pressure but to be anomalously small at low pressure.

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We have measured the attenuation of zero sound in the A phase of superfluid ^{3}He at low pressure down to temperatures less than $0.4T_c$. Since the dominant damping mechanism is Cooper pair breaking, these measurements yield direct information on the microscopic structure of the A phase.

As is now well established,[1] the A phase owes its existence above the polycritical pressure in zero magnetic field to strong-coupling effects and, more specifically, to the spin-fluctuation mechanism put forward by Brinkman and Anderson.[2] Its domain of existence is extended down to absolute zero by the application of magnetic fields of a few kilogauss[3] which suppress the $S_z = 0$ component of the B phase while leaving the gap param-

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eter $\Delta(T)$ of the A phase unaltered to order $(\mu H/E_F)^2$.[4] Hence, a meaningful experimental study of $\Delta(T)$ to very low temperatures can be attempted. The use of zero-sound attenuation as a probe of superfluid pairing was originally suggested by Wölfle[5] and relies on the fact that the pair-breaking condition $\hbar\omega \geq 2\Delta_0(T)\sin\theta$, where $\Delta_0(T)$ is the equatorial value of the gap and θ the angle between $\vec{1}$ and a given direction on the Fermi surface, can be met whatever the energy $\hbar\omega$ of the incoming phonon. The attenuation α of zero sound with wave vector $\vec{q}$ due to this mechanism can be expressed as[6]

$$\alpha = q[(c_0 - c_1)/c_1]\,\mathrm{Im}\xi_0(\omega). \tag{1}$$

The quantity $(c_0 - c_1)/c_1$ is the measured relative difference between the zero- and first-sound velocities. The complex function $\xi_0(\omega)$, whose real part is linked to sound velocity, has been evaluated by a number of authors[6] in the framework of weak-coupling theory. Fermi-liquid corrections involving Landau parameters $F_l{}^s$ of order $l \geq 2$ should be included in ξ_0. However, these terms are very cumbersome to calculate and their effect is not thought to be large. They have been neglected in the theory. As a manifestation of the A-phase anisotropy, $\xi_0(\omega)$ depends on the angle between $\vec{1}$ and $\vec{q}$. The expression of ξ_0 in the case of interest here, where $\vec{q}$ is perpendicular to $\vec{1}$, is given by

$$\xi_\perp{}^0 = \frac{45}{16}\left(\frac{3}{2}\langle\lambda\sin^4\theta\rangle - \frac{\langle\lambda\sin^2\theta\rangle^2}{\langle\lambda\rangle} - \frac{1}{4}\,\frac{\langle\lambda\sin^2\theta\rangle^2}{\langle\lambda\rangle - 2(\Delta_0/\omega)^2\langle\lambda\sin^2\theta\rangle}\right). \tag{2}$$

The angular brackets denote angular averages over the Fermi surface (θ varying between 0 and $\pi/2$) and λ is a function of T/ω and ω/Δ which can be viewed as an effective superfluid fraction $\rho_s(\omega)/\rho$ at the frequency ω and is expressed by

$$\lambda = \int_{-\infty}^{+\infty} d\epsilon\,\frac{\tanh E/2T}{E^2 - \omega^2/4}\,\frac{|\Delta^2|}{2E}, \tag{3}$$

where $E = (\epsilon^2 + |\Delta^2|)^{1/2}$ and $\Delta = \Delta_0\sin\theta$. The integrand of Eq. (3) exhibits a pole at $\omega = 2\Delta_0\sin\theta$. The integral has to be taken as the sum of a real (Cauchy) principal value and an imaginary residue, which, in turn, yields an imaginary contribution to $\xi_\perp{}^0$. This contribution represents the pair-breaking mechanism.

The last term of Eq. (2) also exhibits a pole, known as the clapping mode[7] which occurs at $\omega \simeq 1.23\Delta(T)$. In different orientations of $\vec{q}$ with respect to $\vec{1}$, there exist other modes, known as the normal flapping and superflapping modes which do not appear in $\xi_\perp{}^0$. All these modes are largely smeared into broad resonances by pair breaking, even at $T = 0$, as shown analytically by Dombre and Combescot.[8]

Wölfle and Koch[9] have evaluated numerically the sound attenuation in a number of experimental situations in small fields, close to T_c, above the polycritical pressure. These authors have also included the effect of collisions between quasiparticles which is quite sizable in the vicinity of T_c. They report a fair agreement between theory and experiment which lends support to the use of the weak-coupling model, at least close to T_c and at high pressure. This conclusion is supported by our findings at medium pressure but no more so at low pressure.

The effect of a field on sound propagation has been evaluated theoretically by Tewordt and Schopohl in the case $\vec{q} \parallel \vec{1}$.[10] For this orientation, the mode frequencies ω_i are shifted by the field to $(\omega_i{}^2 + \gamma'^2 H^2)^{1/2}$, where γ' is a renormalized gyromagnetic ratio. The off-resonance contribution to α is unaltered. In our cell, $\vec{q}$ is parallel to $\vec{H}$ and hence perpendicular to $\vec{1}$ (except, possibly, in the immediate vicinity of the cell walls). For this geometry, the flapping modes are not excited and we do not expect an effect of the field on $\xi_\perp{}^0$.

The sonic cell and spectrometer, the cryogenics, and the thermometry are the same as used in our previous work.[11] Absolute values of the attenuation have been calibrated by two independent methods. In the first, the normal-phase attenuation has been forced to a T^2 law. In the second, advantage has been taken of the fact that α in the B phase[11] vanishes as $T \to 0$ for frequencies which are far off the squashing and real squashing modes. Both calibration methods agree to ~0.2 cm^{-1}. Reproducibility from run to run is of the same order. The results for α, plotted against T/T_c in Fig. 1, do not depend on the magnitude of the applied field from 4 to 8 kG. At 15.0 bars, the direction of H was tilted by 7° by application of a transverse field. A definite variation of α by about 1 cm^{-1} was noted. The alignment of $\vec{H}$ and $\vec{q}$ is achieved by construction to better than 3°. Although great care was exercised in the measurements to avoid nonlinear effects,[12] the true zero amplitude may be higher than the values in Fig. 1, although by less than 0.5 cm^{-1}. We estimate the absolute uncertainty on temperature to

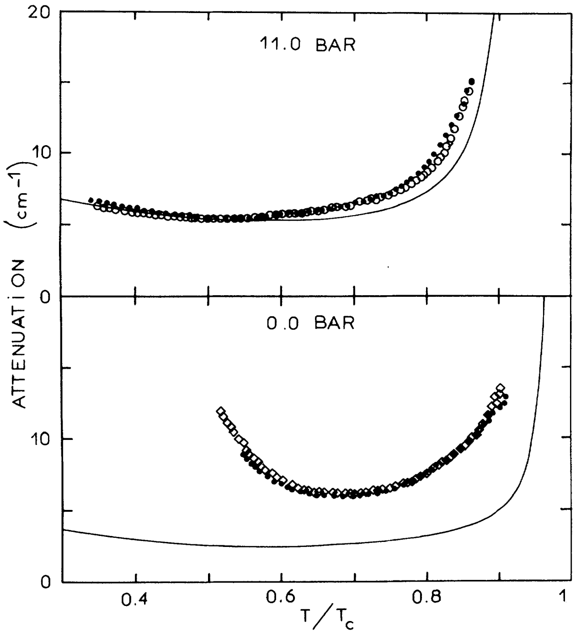

FIG. 1. Amplitude attenuation of zero sound in inverse centimeters vs T/T_c in superfluid ^{3}He-A at the pressures of 0.0 bar (bottom) and 11.0 bars (top) for two values of the applied magnetic field (diamonds, 4 kG; open circles, 5 kG; closed circles, 8 kG). The plain curves are the weak-coupling theoretical results obtained from Eqs. (1)–(3), with use of the gap function and the values of the parameters given in Fig. 2 and in its caption.

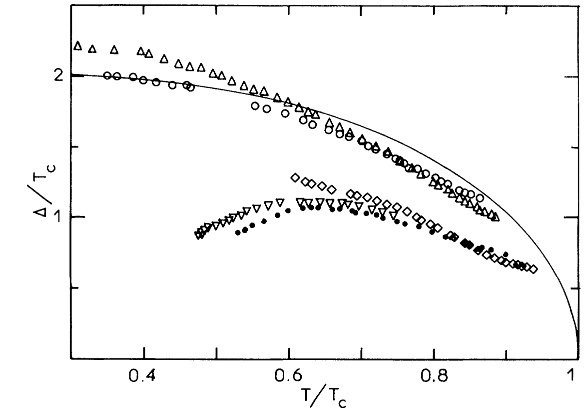

FIG. 2. Normalized gap vs T/T_c obtained from the zero-sound attenuation as explained in the text at various pressures (closed circles, 0.0 bar; inverted triangles, 0.37 bar; diamonds, 1.0 bar; open circles, 11.0 bars; triangles, 15.0 bars). The values of $(c_0-c_1)/c_1$ and T_c used in the analysis at these pressures are, respectively, 0.037, 1.05 mK; 0.0345, 1.10 mK; 0.0312, 1.185 mK; 0.0126, 2.01 mK; and 0.010, 2.19 mK. The plain curve represents the predictions of weak-coupling theory for the Anderson-Brinkman-Morel state (Ref. 13).

be less than 4% at T_c and less than 10% at the lowest temperatures.

With much the same numerical techniques as in Ref. 9, we have computed α for given values of T/T_c and ω/Δ using Eqs. (1)–(3). Comparison with the observed attenuation yields an experimentally determined value of the gap parameter. The resulting gap functions $\Delta(T_c, T/T_c)$ are plotted in Fig. 2 normalized by k_BT_c. The experimental Δ deviates strongly from the weak-coupling value for the Anderson-Brinkman-Morel state[13] at low pressure. The discrepancy between the weak-coupling theoretical attenuation and the observed values can be seen in Fig. 1 to be well outside experimental errors. Thus, some unidentified factor, not accounted for in existing theories, plays an important role in zero-sound attenuation at low pressure.

This conclusion is based on the following assumptions:

(1) Damping by quasiparticle collisions is indeed negligible for high frequencies at $T \ll T_c$ as predicted by theory.[9,14]

(2) $\vec{q}$ and $\vec{l}$ are perpendicular throughout the sample and there is no textural effect. The lack of dependence of α on H, as discussed above, and the reproducibility of the measurements support this assumption. Another indication that we are dealing with a single damping mechanism lies in the nonlinear behavior, reported in Ref. 12. When the excitation level is increased, the medium becomes gradually transparent to sound. It is hard to envision that several different damping mechanisms (i.e., pair breaking and either collisions or textures, etc.) would exhibit the same pattern of nonlinearity.

(3) Fermi-liquid corrections of high order are negligible. The contribution of F_2^s is unknown in the A phase where anisotropy complicates the picture. In the isotropic case, this correction is of the order of $\mathrm{Re}\xi_\perp{}^0 F_2^s/5$, which may amount to $\sim$20% at $T\sim 0$ if $|F_2^s|$ is taken to be of the order of 1. This uncertainty weakens our conclusion although $|F_2^s|$ is not expected to be much larger than 1.

Strong-coupling corrections on the value of the gap parameter in the B phase have been found in a number of experiments[15-17] to be of the order of 20% at high pressure and much less at low pressure. Such a behavior, which we would expect to find also in the A phase, does not point toward strong coupling as a possible cause of the depression of Δ. The discrepancy reported here would be more consistent with an admixture of f waves in the predominantly p-wave superfluid. In this respect, it is worth noting that $\xi_\perp^0$, given by Eq. (2), is obtained from the subtraction of nearly equal angular averages and may be altered quite significantly even by a small change in these averages.

Finally, it cannot be ruled out entirely that the flapping mode, or some more exotic mode such as the out-of-phase flapping mode,[18] is excited despite the fact that the resulting attenuation is quite insensitive to magnetic fields.

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Experimental Evidence for Spin Depolarization in Photoemission

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Nonconservation of photoelectron spin polarization during photoemission has been observed. Initially polarized electrons photoexcited in a germanium single crystal are completely depolarized by paramagnetic gadolinium atoms deposited on the surface of the sample.

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The power of photoemission as a technique to study the electronic properties of solids relies on the fact that a sizable fraction of the excited electrons leave the crystal without having undergone any scattering event that changes the typical observables of the outgoing electron, namely, energy and momentum parallel to the surface. Because of the strong interactions between electrons the mean free path between inelastic collisions is limited: As function of energy it follows a well-established universal curve which applies with good accuracy to all materials.[1] The momentum parallel to an ideal surface is conserved as well[2]; however, if the order of the surface atomic ar-

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