SUPPLEMENTARY INFORMATION

I. EXPERIMENTAL DETAILS

In our experimental arrangement, the cavity spacing of $D = 31.6 \pm 0.1$ μm was measured *in situ* with longitudinal sound at a temperature of 18 mK, using the 17^{th} transducer harmonic which generates a large amplitude longitudinal signal. We note that the texture of the order parameter must be homogeneously oriented along the sound propagation and magnetic field directions, since the diameter of the transducer (0.84 cm) is much larger than the spacing of the cavity. The frequency and amplitude of the transducer resonances are extremely sensitive to the acoustic conditions of the cavity which modify the electrical impedance that we monitor with a continuous-wave, RF-bridge[17,18]. We use odd harmonics of our transducer from the 13^{th} to the 29^{th}, corresponding to frequencies from 76 to 171 MHz. Details of cooling and thermometry can be found elsewhere[18]. We use the weak-coupling-plus model for the gap[44] as tabulated by Halperin and Varoquaux[16], locked to the Greywall temperature scale[43]. This allows us to account for variations of the transverse sound velocity with changes in both pressure and temperature, although temperature effects are very slight at the low temperatures used in this experiment. The accuracy of the Greywall temperature scale[43] determines the accuracy of Δ in the weak-coupling-plus model and is estimated to be 1%.

II. CALCULATION OF THE TRANSVERSE SOUND VELOCITY

In order to calculate the transverse sound velocity from equation (1), which appears in Figs. 2b and 2d, we use the full expression for the $J = 2^-$ mode frequency in a magnetic field given by Moores and Sauls[40] $\Omega_{2^-,m_J}(H)^2 = \Omega_{2^-}^2 + 2m_J g_{2^-}\gamma_{eff} H\omega$, for $m_J = \pm 1$. Additionally, we used the quasiparticle restoring force for transverse sound[40] $\Lambda_0 = \frac{F_1^s}{15}(1 - \lambda)(1 + \frac{F_2^s}{5})/(1 + \lambda\frac{F_2^s}{5})$, which includes all quasiparticle interaction terms[16], F_l^s, for $l \leq 2$. Similarly $\Lambda_{2^-} = \frac{2F_1^s}{75}\lambda(1 + \frac{F_2^s}{5})^2)/(1 + \lambda\frac{F_2^s}{5})$.

The Tsuneto function $\lambda(\omega, T)$ can be thought of as a frequency dependent superfluid density[40], that depends on the gap amplitude, Δ. Therefore, at low temperatures the quasiparticle term is small and Λ_{2^-} is enhanced. In our calculations of velocity we use

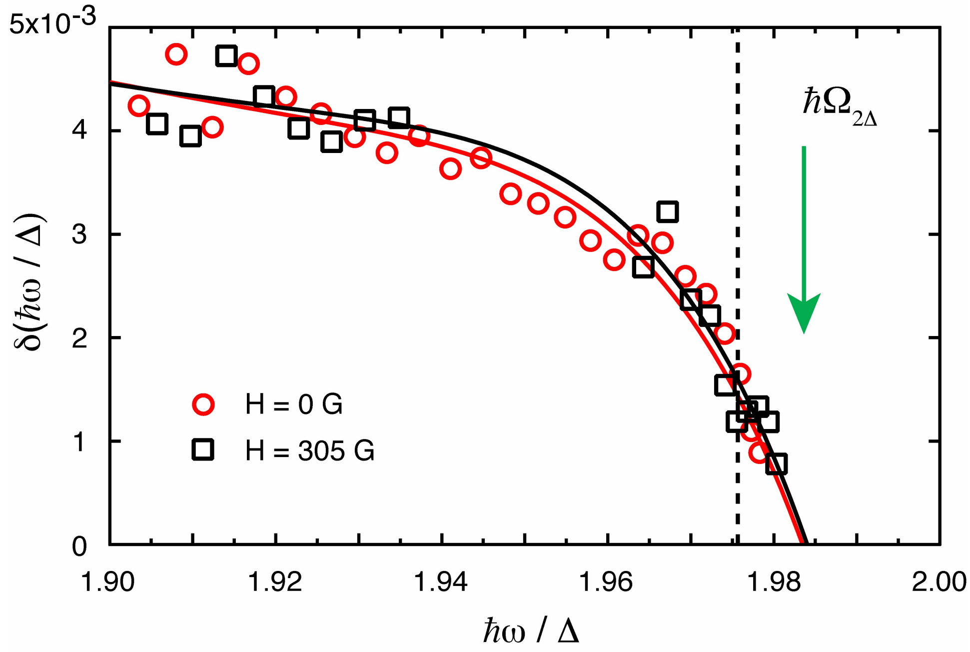

Figure S1 Period of acoustic energy oscillations from the acoustic cavity interference of Fig. 3b as the acoustic frequency approaches the 2Δ mode, green arrow, for zero magnetic field (red circles) and 305 G (black squares). We identify $\hbar\Omega_{2\Delta}$ from the extrapolation to zero (where the velocity diverges). Both traces point to the same energy within experimental resolution. For comparison the energy for the collective mode $J = 1^-$, $m_J = -1$ in $H = 305$ G with a Zeeman splitting of $g = 0.4$ is shown as a black dashed line as expected theoretically[46]. The absence of data over a small range near 1.95 for the 305 G trace corresponds to the region where Faraday rotation has decreased the amplitude of the acoustic oscillations and the period is not reliably determined.

$\lambda(\omega, T)$ adapted to incorporate the weak-coupling-plus gap as described above[44]. The dispersion relation, equation (1), holds in the long wavelength limit, $kv_F \ll \omega$. Consequently, our estimation of the coupling strength of the 2Δ-mode may be modified in a full q-dependent analysis.

III. LOCATING THE NEW MODE

As the acoustic frequency approaches that of an order parameter collective mode, the sound velocity and attenuation diverge. In order to pinpoint the frequency of the mode we plot the inverse of the signal oscillation amplitude (attenuation) and the oscillation period (velocity) and extrapolate to zero. In Fig. S1 we show this procedure for the oscillation period for two experiments with zero magnetic field and $H = 305$ G. The curves are fits to guide the eye. Nonetheless it is clear that the mode frequency can be precisely determined.

This is the method used to obtain Fig. 3a. Moreover, the 2Δ mode does not appear to have a Landé g-factor that is large enough to appear directly as a shift of the data trace in $H = 305$ G. As a point of comparison, if the Landé g-factor were to be of the magnitude theoretically predicted for the $J = 1^-$ mode a shift in frequency indicated by the dashed line would be expected[46].

IV. LONGITUDINAL SOUND AND THE $J = 1^-$ MODE

Ling *et al.*[20] reported the existence of a mode, with a splitting which they ascribed to $J = 1^-$, $m_J = \pm 1$. These signatures in longitudinal sound attenuation appeared only in the presence of an applied magnetic field, which in their case was oriented transverse to the sound propagation direction. Since then it has been shown[19,45] that only the $m_J = 0$ mode can couple to longitudinal sound, contrary to what was originally proposed[20,46]. Additionally, Ashida *et al.*[19] find that the anomalies observed in longitudinal sound attenuation in non-zero field[20] can be accounted for in terms of pair-breaking coupled to $J = 1^-$, $m_J = 0$ and $J = 2^-$, $m_J = 0, \pm 2$ modes, in some combination, but that the analysis is complicated by the inhomogeneous texture associated with the experimental conditions of a transverse magnetic field. From their analysis it appears that identification of the $J = 1^-$ mode is not yet well-established. A comparable calculation has not been performed for transverse sound; however, on symmetry grounds the $J = 1^-$ mode cannot couple to transverse sound in zero magnetic field, and in a non-zero field the coupling to transverse sound is indirect via the $J = 2^-$ mode. We estimate the coupling strength for this process to be very weak compared to what we measure for the 2Δ mode providing additional evidence that the $J = 1^-$ mode is not responsible for our observations.