

Dynamic interaction between chiral currents and surface waves in topological superfluid

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Despite extensive experimental efforts over the past two decades, topological superconductivity and Majorana fermions have remained elusive. By shifting focus from superconductors—whose topological nature is often debated—to the well-established p -wave superfluid ^3He , we propose an experimental approach designed to identify quantum Hall nature of the chiral superfluid $^3\text{He-A}$ phase. The method exploits the interaction between surface waves and the induced current transverse to varying density of the superfluid. It allows one to quantitatively measure this Hall current and thus to directly evaluate whether the Hall current is quantized or not. As the quantum Hall effect is a hallmark of topological condensed matter, such an experiment would offer direct evidence of nontrivial topology of superfluid ^3He . With the advent of microscopic superfluid resonators coupled to optical cavities, this proposal is both realistic and promising, potentially opening the door to future experiments ranging from topological acoustics to fault-tolerant quantum computing.

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I. INTRODUCTION

One of the major breakthroughs in modern condensed matter physics is the incorporation of topology in material classifications, beyond traditional classifications based on symmetry [1]. However, despite the successful experimental realization of topological insulators [2,3] and subsequently topological semimetals [4–6], topological superconductors remain an elusive goal [7–10]. Perhaps the most intriguing prediction about topological superconductors is that their boundaries are a potential host to a quasiparticle known as the Majorana fermion [7–9]. Direct synthesis of such a superconductor has been found to be controversial [10]. A more subtle strategy of combining two or more materials [11], each exhibiting nontrivial topology and superconductivity separately, led to observations that suggested the existence of Majorana fermions [12–16]. However, the interpretation of these observations has since been questioned [17–21].

While nearly all of these efforts have been focused on topological superconductors [10], the advent of topology also shed light on a long-known system: superfluid helium-3. Unlike in

any known superconductor, the p -wave nature of superfluid ^3He 's order parameter is unequivocally established in its A and B phases. Thus, it is expected to contain Majorana fermions, and a number of experiments report the existence of surface states in superfluid ^3He [22–24]. However, direct experimental confirmation of the topological nature of the bulk order parameter or the Majorana nature of these boundary states is still an outstanding challenge. One of the most powerful ways in which nontrivial topology reveals itself in a condensed matter system is in the form of a quantized conductance whose quantization is strongly tied to the topological invariants of a system. Given the analogy between the quantum Hall system and topological chiral superconductivity, quantized Hall conductivity is expected with the chiral edge current carried by Majorana fermions. However, quantum Hall measurements rely on the presence of charge in the system and are thus inapplicable to liquid ^3He .

In this work, we propose an experiment that is designed to probe this quantized Hall effect and nontrivial topology in a charge-neutral superfluid ^3He . The mass current density in superfluid $^3\text{He-A}$ at zero temperature is given by [25,26]

$$\vec{g} = \rho \vec{v} + \frac{\hbar}{4M} \vec{\nabla} \times \rho \vec{l} - \frac{\hbar}{2M} c_0 \vec{l} (\vec{l} \cdot \vec{\nabla} \times \vec{l}), \quad (1)$$

where $\vec{v}$ is the superfluid velocity, $\hbar$ is the reduced Planck constant, $M = 5.0 \times 10^{-27}$ kg is the mass of a helium-3 atom, ρ is the local superfluid density, $\vec{l}$ is the local anisotropy axis, and c_0 is an integral over single-particle momentum states [27]. In the case of a helium thin film, ρ can be understood as the bulk density ρ_0 of superfluid helium-3, locally scaled with the film thickness variation: $\rho = \rho_0 \times (1 + \eta/h_0)$, with

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$\rho_0 = 82 \text{ kg/m}^3$, h_0 the equilibrium film height, and η the film height perturbation due to the acoustic wave. The first term corresponds to an irrotational mass current, such as that arising from center-of-mass motion or acoustic waves, and the third term generates an essentially out-of-plane current, which we can neglect as its coupling to the in-plane flow due to sound modes is minimal. The second term in Eq. (1) describes the chiral nature of the $^3\text{He-A}$. The chiral current can be thought of as analogous to the bound current density j in a conductor of varying magnetization $\vec{M}$ ($j \propto |\vec{\nabla} \times \vec{M}|$), where the Cooper pairs' orbital angular momentum plays the role of the magnetization [26].

II. CONCEPT

To investigate the coupling between chiral edge current and sound modes, we consider the case of a helium-3 surface-sound resonator and focus on the role of the second term in Eq. (1). Spatial variations in $\vec{l}$ or in ρ lead to net mass currents. The direction of the current depends on the preferred orientation of the orbital angular momentum $\hat{l}$, which is to first order either aligned or antialigned with the resonator's surface normal $\hat{z}$ [28]. In such a case, one can rewrite the second term as

$$\vec{g}_{\text{chi}} = \pm \frac{\hbar}{4} \vec{\nabla} n \times \hat{z}, \quad (2)$$

where n is the number density of ^3He . What immediately stands out is that a transverse current is induced by the density gradient. This can be interpreted as dissipationless Hall viscosity [29] with the strain rate $\vec{\nabla} n$ and subsequent momentum current, $\vec{g}_{\text{chi}}$, orthogonal to each other. We identify the Hall viscosity as the coefficient linking $\vec{\nabla} n$ and $\vec{g}_{\text{chi}}$, given by $\hbar/4$ [30]. Despite significant research efforts, quantum Hall viscosity has been obscured in electronic quantum Hall systems by the presence of the lattice [31–33]. Measuring the transverse current could quantify the quantum Hall viscosity of a superfluid and thereby identify its topology.

Here, we propose a technique to measure this quantum Hall current by leveraging the spatially distributed current, which arises in response to local density variations in a p -wave superconductor or superfluid. We show how this effect arises dynamically due to sound modes in superfluid $^3\text{He-A}$, and how it can be detected via its dispersive interaction with these very sound modes [34]. In a two-dimensional picture, sound-induced deformations in the surface can be thought of as fluctuations in planar density. Such fluctuations induce a transverse current in the direction perpendicular to the density gradient according to Eq. (2). This dynamical effect is a manifestation of the dissipationless quantum Hall viscosity [29], an effect reminiscent of the transverse current induced in response to a voltage gradient in a Hall system. We term this spatially and temporally varying effect the *dynamic* chiral current, as opposed to the conventionally predicted *static* chiral edge current at the superfluid boundary [35].

This transverse current couples orthogonal pairs of superfluid sound modes and as a result lifts the degeneracy between them. We find that the magnitude of this mode splitting scales inversely with the resonator size, so that to obtain a clear signature, microscale superfluid resonators are essential. Recent

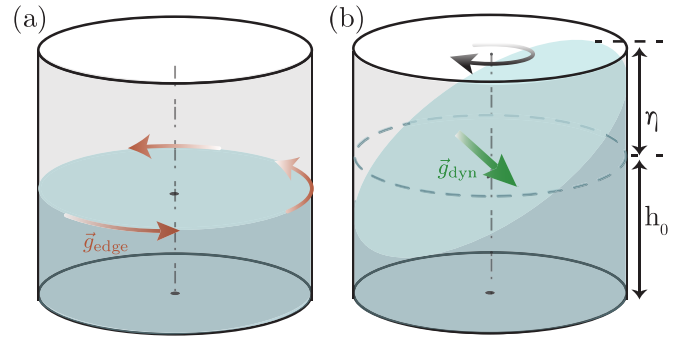


FIG. 1. Illustration of static and dynamic chiral currents. (a) Static chiral edge current at the resonator boundary. Red arrows: direction and location of the induced flow; black arrow: rotation direction of the fluid antinode. (b) Dynamic chiral current, present throughout the resonator. Green arrow: direction of the induced flow; black curved arrow: rotational direction of the clockwise rotating sound mode. h_0 and η (not to scale): equilibrium film height and acoustic perturbation, respectively.

advances in miniaturized superfluid thin-film resonators, in which superfluid surface modes are read out optically [36–40], hence open the door to detecting this deformation-induced Hall current.

III. RESULTS

A. Coupling between chiral edge current and sound modes

For surface-sound resonators, we consider the large wavelength, small amplitude limit, where the acoustic wavelength λ is much larger than the film thickness h , which is in turn much larger than the out-of-plane displacement amplitude η of the film. This implies that the motion of the fluid is essentially in plane and out-of-plane motion is negligible (see Appendix A). Such a condition arises naturally in experiments, where a resonator substrate is coated with a film of superfluid helium (see, e.g., Refs. [38,41]). Under these assumptions, Eq. (2) can be expanded as $\vec{g}_{\text{chi}} = \vec{g}_{\text{edge}} + \vec{g}_{\text{dyn}}$. Here, $\vec{g}_{\text{edge}}$ is the conventional static chiral edge current [35,42], whereas $\vec{g}_{\text{dyn}}$ is the dynamic current induced by the sound modes. Our proposed method allows us to probe the latter. These two contributions to the in-plane mass current are illustrated in Fig. 1. The static mass current, shown in Fig. 1(a), results from the steady-state variations in the film's thickness due to physical confinement within the resonator geometry. In the case where the film thickness and density are uniform, and the thickness drops to zero at an abrupt specular boundary, it corresponds to the ground-state angular momentum $N\hbar/2$ [42,43], with N the number of helium nuclei contained in the film.

The dynamic mass current, shown in Fig. 1(b), is proportional to the gradient of the (temporally and spatially varying) superfluid film height due to the sound mode.

Both the static and the dynamic contributions to the chiral edge current couple to acoustic waves via their direct overlap with the flow fields of the sound modes. In the static case, the chiral mass current leads to a time-independent flow field that superposes with the flow field of acoustic waves. The energy of an acoustic mode increases if the interference is

constructive, and vice versa. On the other hand, the dynamic chiral current arising from the density perturbation due to one acoustic mode can interfere with the flow field of another mode, leading to a “dynamic coupling” between acoustic modes.

B. Dynamic coupling

We investigate the effect of the dynamic coupling in the analytically tractable case of a circular superfluid surface-wave resonator [44]. Note that while we consider surface waves here, this discussion is also valid for density waves, such as first sound, with the appropriate transformations (see Appendix B). For a thin-film superfluid surface-wave resonator with equilibrium film height h_0 , assuming the distortion of the acoustic mode shape due to the chiral current to be small, the surface deflection profile $\eta(r, \theta, t)$ due to a surface-sound wave of frequency $f = \Omega/2\pi$ confined to a circular domain of radius R can be expressed in cylindrical coordinates as [45,46]

$$\eta(r, \theta, t) = \eta_0 J_m\left(\zeta \frac{r}{R}\right) e^{i(m\theta \pm \Omega t)}, \quad (3)$$

where η_0 is the wave amplitude, J_m is the Bessel function of the first kind, with m the mode’s azimuthal order, and ζ is a parameter that depends on the mode radial order n and the boundary conditions [45]. This approximation is justified if the ratio of the mass currents due to the chiral flow to that of the acoustic mode is small. In our case, $\frac{g_{\text{dyn}}}{g_s} = \frac{v_{\text{dyn}}}{v_s} \propto \frac{\kappa_3}{R c_s}$, with g_s the mass current associated with the surface wave [corresponding to the first term in Eq. (1)], c_s the speed of sound, and $\kappa_3 = \frac{\pi\hbar}{M}$ the circulation quantum in ^3He . For the resonators considered here, this ratio does not exceed a few percent (see Appendix A). We can directly identify the chiral current velocity as $\vec{v}_{\text{dyn}} = \vec{g}_{\text{dyn}}/\rho_0$. The fluid-flow velocity associated with the surface waves, corresponding to the first term in Eq. (1), can be calculated from Eq. (3) to be $\vec{v}_s = \pm i c_s^2 \vec{\nabla} \eta / (h_0 \Omega)$, where “+” and “−” correspond to clockwise (CW) and counterclockwise (CCW) modes, respectively [45].

The energy difference ΔE between the two sound eigenmodes arising from a background flow is, using the polarization identity and noting that $\vec{v}_s$ and $\vec{v}_{\text{dyn}}$ are independent of z [46],

$$\begin{aligned} \Delta E(t) = 2\rho_0 \int_{\theta=0}^{2\pi} \int_{r=0}^R \vec{v}_s(r, \theta, t) \cdot \vec{v}_{\text{dyn}}(r, \theta) \\ \times (h_0 + \eta(r, \theta, t)) r dr d\theta. \end{aligned} \quad (4)$$

Since $\eta \ll h_0$, we find

$$\Delta E(t) = 2\rho_0 h_0 \int_{\theta=0}^{2\pi} \int_{r=0}^R \vec{v}_s \cdot \vec{v}_{\text{dyn}} r dr d\theta. \quad (5)$$

Expressing the velocities as a function of height perturbation (cf. Ref. [46]),

$$\vec{v}_s = \pm \frac{ic_s^2}{\Omega h_0} \vec{\nabla} \eta \quad \text{and} \quad \vec{v}_{\text{dyn}} = \frac{\hbar}{4M h_0} \vec{\nabla} \times (\eta \hat{z}), \quad (6)$$

the dynamic part of the chiral current can be written as [cf. Eq. (2)]

$$\vec{g}_{\text{dyn}} = -\rho_0 \frac{\hbar}{2M h_0} \hat{z} \times \vec{\nabla} \eta + g_{\text{torque}}, \quad (7)$$

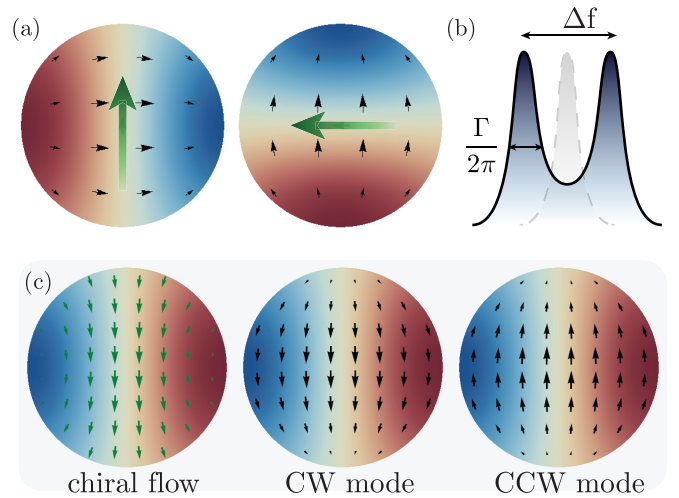


FIG. 2. Dynamic coupling mechanism for the lowest-frequency ($m = 1, n = 1$) sound eigenmode of a circular resonator with free boundary conditions [45]. Background color: height profile $\eta(r, \theta)$; black arrows: superfluid surface-sound flow field $\vec{v}_s$; green arrows: dynamically induced chiral flow $\vec{v}_{\text{dyn}}$. (a) standing wave basis. (b) Illustration of spectral sound mode frequency splitting. (c) Rotating mode basis. Left panel: mass current vector field $\propto \vec{v}_{\text{dyn}}$. The chiral flow is orthogonal to the height gradient. Middle and right panels: Flow fields for the clockwise and counterclockwise sound modes, respectively.

where $\eta \ll h_0$ is the surface deflection due to the sound wave (see Fig. 1) and g_{torque} is the out-of-plane component of the chiral mass current, which is negligible for the resonators considered here (see Appendix A). Without loss of generality, we have assumed the anisotropy axis $\hat{l}$ to be aligned with the normal to the surface $\hat{z}$, i.e., $\hat{l} = \hat{z}$.

To illustrate the effect of the chiral current on the acoustic eigenmodes, we first consider the lowest-frequency excitation of a circular resonator of radius R with free boundary conditions, where the fluid oscillates back and forth from one side of the resonator to the other, corresponding to the lowest-frequency Bessel mode [45]. In the absence of chiral currents, this type of excitation can be described in the standing wave basis formed by the two frequency-degenerate normal modes shown in Fig. 2(a), where the sound velocity field $\vec{v}_s$ is represented by the black arrows. The presence of the sound-induced density gradient η leads via Eq. (7) to an additional mass current $\vec{g}_{\text{dyn}}$ (and associated flow field $\vec{v}_{\text{dyn}}$) in a direction orthogonal to the height gradient, represented by the green arrows. This net mass flow will interact either constructively or destructively with the sound-induced flow, leading to a coupling between previously orthogonal standing wave modes. The coupling introduces a frequency splitting between these waves, in a mechanism that is reminiscent of the third-sound splitting due to the presence of quantized vortices [38,41,46–48]. This splitting provides a mechanism by which the chiral current can be detected experimentally, e.g., by optomechanical readout [38].

This lifting of the degeneracy between normal modes is experimentally resolvable if its rate Δf exceeds the acoustic mode’s linewidth $\Gamma/(2\pi)$ [38,46,48], as shown in Fig. 2(b).

In the presence of this coupling, CW and CCW propagating surface-sound waves become the diagonal mode basis [see Fig. 2(c)]. Here, we see that the induced flow is, respectively, aligned or antialigned with the CW and CCW sound flow, leading to an effective Doppler shift. This implies that the acoustic mode with the same rotation direction as the induced chiral mass current is frequency up-shifted, while that of opposite handedness is frequency down-shifted. This can alternatively be viewed as a consequence of the sound wave (despite creating an irrotational flow field) carrying angular momentum, as more fluid advances under the peaks than recedes under the troughs [46].

We estimate the magnitude Δf of the splitting due to the dynamic coupling using a perturbative approach, following Ref. [46] using Eq. (3), and performing straightforward vector calculus we obtain

$$\Delta E(t) = \frac{\rho_0 \hbar c_s^2 m}{2M \Omega h_0} \int_0^R \int_0^{2\pi} \eta(r) \eta'(r) dr d\theta, \quad (8)$$

where $\eta(r) = \eta_0 J_m(\zeta r/R)$ and $\eta' = \partial \eta / \partial r$. Noting that ΔE is now independent of time, this simplifies to

$$\Delta E = \frac{\pi \rho_0 \hbar c_s^2 m}{M \Omega h_0} \int_0^R \eta(r) \eta'(r) dr. \quad (9)$$

The energy E stored in the sound mode is given for $m < 0$ modes by [45]

$$E = \frac{\pi \rho_0 c_s^2}{2h_0} \int_0^R \eta^2(r) r dr. \quad (10)$$

The fractional change in energy is thus

$$\frac{\Delta E}{E} = \frac{2 \hbar m}{M \Omega} \frac{\int_0^R \eta(r) \eta'(r) dr}{\int_0^R \eta^2(r) r dr}. \quad (11)$$

For a harmonic oscillator, E is proportional to Ω^2 and $\frac{\Delta E}{E} = 2 \frac{\Delta \Omega}{\Omega}$; hence the frequency splitting Δf between the two counterrotating sound modes is given by

$$\Delta f_{\text{dyn}} = \frac{\Omega}{4\pi} \frac{\Delta E}{E} = \frac{\kappa_3 m}{4\pi^2} \frac{\eta^2(R)}{\int_0^R \eta^2(r) r dr}. \quad (12)$$

The splitting is independent of acoustic mode amplitude and vanishes for fixed boundary conditions, for which $\eta(R) = 0$. The existence of an induced transverse flow in the hydrodynamic oscillation of a superfluid $^3\text{He-A}$ film had been noted [49], and alluded to in the context of superfluid third-sound resonators [34]. However, the splitting of sound modes scales inversely with resonator area, and hence a measurable effect only becomes possible with the advent of the microscopic superfluid resonators [36,38,39] considered in this work.

The inverse scaling of the mode splitting with resonator area can be understood as follows: For a given wave amplitude, the magnitude of the surface gradient, and hence of the induced chiral flow velocity, is inversely proportional to R . The frequency shift Δf_{dyn} is proportional to R^{-2} . This can be seen from Eq. (12) as, on average, $\eta^2(r)$ and $\eta^2(R)$ are independent of R , while $\int_0^R r dr \propto R^2$.

C. Static chiral edge current

The static chiral edge current, which resides within a layer of a depth of the healing length $\xi \approx 100$ nm from the resonator boundary, in the case of an abruptly vanishing superfluid density is given by Refs. [35,42]. The calculation is similar to the above calculation of the dynamic chiral current. We start from the second term in Eq. (1) to get

$$\vec{g}_{\text{edge}} = -\frac{\hbar}{4M} \vec{l} \times \vec{\nabla} \rho_{\text{edge}}, \quad (13)$$

where $\vec{\nabla} \rho_{\text{edge}}$ quantifies the dropoff of the condensate density at the boundary. We assume that ρ drops from ρ_0 to 0 within a length scale of the healing length ξ . For a thin film of single domain, the l vector remains constant perpendicular to the film, i.e., $\vec{l} = \begin{pmatrix} 0 \\ 1 \end{pmatrix}$.

Then, the chiral current density at the resonator boundary given by Eq. (13) results in the superfluid velocity

$$\vec{v}_{\text{edge}} = \vec{g}_{\text{edge}} / \rho_0 \simeq \frac{\hbar}{4M\xi} \hat{e}_\theta \approx 5 \text{ cm/s}, \quad (14)$$

independently of the resonator geometry, with $\xi \sim 100$ nm for the case of an abrupt and specular boundary of superfluid ^3He at 0 bar.

The macroscopic angular momentum L_z associated with this mass current density $\vec{g}_{\text{edge}}$ is, with $R \gg \xi$,

$$L_z = m v R = g_{\text{edge}} (2\pi R) \xi R = \frac{\hbar \pi R^2 \rho_0}{2M} = \frac{\hbar}{2} N, \quad (15)$$

where N is the number of particles in the fluid. The presence of the edge current leads to a background flow, which will lead to an energy difference ΔE between CW and CCW modes. We follow Ref. [46] in calculating the resulting frequency splitting:

$$\Delta E = 2\rho_0 \int_{\theta=0}^{2\pi} \int_{r=0}^R (h_0 + \eta) \vec{v}_s \cdot \vec{v}_{\text{edge}} r dr d\theta. \quad (16)$$

Due to symmetry, $\oint \vec{v}_s \cdot \vec{v}_{\text{edge}} = 0$, and hence the energy difference averaged over an oscillation period T is given by

$$\begin{aligned} \langle \Delta E \rangle &= \frac{1}{T} \int_0^T \Delta E(t) dt \\ &= 2\rho_0 \int_0^{2\pi} \int_0^R r (\vec{v}_{\text{edge}} \cdot \hat{e}_r \langle \vec{v}_s \cdot \hat{e}_r \eta \rangle_T \\ &\quad + \vec{v}_{\text{edge}} \cdot \hat{e}_\theta \langle \vec{v}_s \cdot \hat{e}_\theta \eta \rangle_T) dr d\theta. \end{aligned} \quad (17)$$

From Eqs. (3) and (6), we see that $\vec{v}_s \cdot \hat{e}_r$ and η are out of phase and $\vec{v}_s \cdot \hat{e}_\theta$ and η are in phase. Hence, the first part of the integral in Eq. (17) is zero and the remaining integral becomes $\frac{1}{2} |\vec{v}_s \cdot \hat{e}_\theta| |\eta|$.

Performing the integral and inserting Eqs. (3) and (6), we find the same form as in Eq. (9), and hence the same magnitude of splitting as for the dynamic case:

$$\Delta f_{\text{edge}} = \frac{\kappa_3 m}{4\pi^2} \frac{\eta^2(R)}{\int_0^R \eta^2(r) r dr}. \quad (18)$$

It is important to note here that Δf_{edge} is calculated with respect to CW and CCW rotating modes, whereas Δf was computed in the standing mode basis.

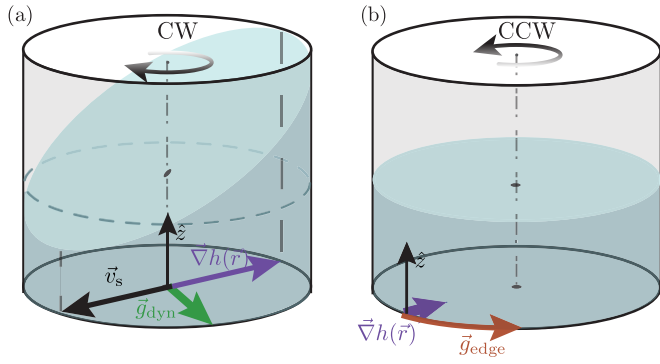


FIG. 3. (a) Illustration of the directionality of dynamic chiral current. $\vec{\nabla}h(\vec{r}) = \vec{\nabla}\eta(\vec{r})$ (purple arrow) points toward the peak of the acoustic wave. The direction of the dynamical chiral current (green arrow) is given by $-\hat{z} \times \vec{\nabla}h(\vec{r})$, so that it points toward the right-hand side of the peak, increasing the energy of the CW rotating acoustic mode (curved black arrow). The trough and peak of the acoustic mode are indicated by dashed vertical lines, and the sound velocity is indicated by a black arrow. All arrows are shown on the bottom of the superfluid resonator for ease of illustration, although the fluid flow is distributed through the whole film. (b) Illustration of the directionality of the static edge current. $\vec{\nabla}h(\vec{r}) = -\hat{e}_r h_0 \delta(r - R)$ (purple arrow) points toward the inside of the resonator at its boundary. The direction of the edge current (curved red arrow) is given by $-\hat{z} \times \vec{\nabla}h(\vec{r})$ and points in the CCW direction, hence increasing the energy of the CCW acoustic mode. By analogy, if the film thickness increases at the resonator boundary (so that $\vec{\nabla}h(\vec{r})$ points outward, not shown in figure), the static part of the chiral current increases the energy of the CW mode. While panels (a) and (b) refer to the same resonator, in panel (b) the deformation due to the sound mode is not shown to clarify that it does not cause any static current.

For a circular resonator with abruptly vanishing superfluid density (see, e.g., Refs. [50–52]), we find $\Delta f_{\text{edge}} = -\Delta f_{\text{dyn}}$, so that dynamic and static contributions cancel out. An illustrative explanation is shown in Fig. 3.

On the other hand, if the film thickness increases at the boundary, the contributions add constructively. Both boundary conditions can be engineered and are found experimentally [34,36,38,53]. The first observation of ^3He third sound employed an $R \sim 2$ cm resonator, for which the magnitude of the splitting Δf would have been in the hundreds of microhertz and thus undetectable. Only recently, advances in miniaturization and the fabrication of micrometer-sized third-sound resonators [36,38,39] have put this observation within reach. For a circular resonator of $20\ \mu$ diameter and a film thickness of 100 nm, the magnitude of the dynamic splitting is appreciable: 47 Hz for the $(m = 1, n = 1, \Omega/(2\pi) = 5.1\ \text{kHz})$ mode and 555 Hz for the $(m = 6, n = 1, \Omega/(2\pi) = 70\ \text{kHz})$ mode. This splitting would be detectable with sound quality factors on the order of 10^2 in both cases. For comparison, in superfluid ^4He resonators of the size and shape considered here, quality factors in excess of 10^3 are common [36,39].

As a concrete example, we propose silicon-on-insulator pillar designs, which we already fabricate for optomechanical experiments with surface waves in helium-4, with standard pillar height of 220 nm. These resonators have recently been shown to be able to confine superfluid surface-sound modes

and allow for optomechanical readout [54]. For this geometry and film thickness, sound waves arise from both van der Waals and surface tension in comparable magnitudes (see Appendix C). The acoustic mismatch at the boundary leads to strong confinement for the higher-order Bessel modes defined by the resonator geometry [38,39]. A range of readout mechanisms may be able to resolve the chiral-current-induced splitting of sound waves, such as capacitive readout [48] or torque sensing [35]. As a concrete example, we focus on optical readout via an optical whispering-gallery mode confined on the perimeter of the pillar [54,55].

IV. CHALLENGES

Optomechanical experiments with ^3He will be inherently challenging because of its low superfluid transition temperature. Heating issues can be minimized through techniques such as pulsed optical operation [56].

Recently, through optimization of the superfluid fountain pressure interaction and high-precision optomechanical detection techniques, the ability to detect and strongly drive superfluid ^4He third-sound modes with optical powers down to the femtowatt range [40], corresponding to less than one average intracavity photon, has been realized. Mechanical modes are optically resolvable if the optomechanical cooperativity is larger than the inverse of the average thermal phonon number $C = 4\bar{n}_{\text{cav}}g_0^2/(\kappa\Gamma) < \bar{n}_{\text{th}}^{-1}$, where $\bar{n}_{\text{cav}}$ is the average intracavity photon number, g_0 is the single-photon optomechanical coupling rate, and κ is the optical decay rate [57]. A temperature of 1 mK yields $\bar{n}_{\text{th}} = \hbar\Omega/(k_B T) = 4.7 \times 10^3$ for the $(m = 1, n = 1)$ mode and parameters considered above. With values from the experiment in Ref. [39], we find $C = \bar{n}_{\text{cav}} \times 2.5 \times 10^{-3}$, allowing the mechanical modes to be resolved with an average of less than one intracavity photon.

It is crucial to unambiguously distinguish signatures of chiral current from other effects [17,19]. Known effects that can produce a sound-mode splitting in superfluid helium resonators are geometric splitting and vortex-induced splitting [38]. As the hydrodynamic equations describing momentum density for the A and B phases of superfluid ^3He are identical except for the terms relating to the chiral current [58], these spurious sources of splitting can be calibrated by transitioning the film between the A and B phases, by tuning temperature and film thickness [51] or magnetic field. Estimates of the vortex density generated during the A – B transition span the range 10^6 – $10^{10}\ \text{m}^{-2}$. The proposed device consists of a disk approximately 100 nm thick with a diameter of order $20\ \mu\text{m}$, corresponding to a footprint area of roughly $10^{-10}\ \text{m}^2$. Under these conditions, the presence of vortices in significant numbers is unlikely.

In addition, vortex formation is inherently stochastic, leading to a distribution of induced shifts across repeated realizations of the transition. In contrast, the generation of chiral currents during the B – A transition is expected to be fully deterministic. Consequently, by repeatedly cycling the system through the B – A transition, vortex-related effects can be either excluded as the origin of the observed splitting or, if present, clearly distinguished from the chiral current contribution. A statistical analysis over multiple transitions would thus yield an unambiguous signature separating

deterministic chiral-current effects from stochastic vortex-induced contributions.

V. CONCLUSIONS

Experimental observation and quantitative measurement of the dynamical chiral current in superfluid ^3He would enable a direct test of whether the Hall response in a topological superfluid is quantized. Moreover, the proposed experiment would provide a practical method for studying dissipationless quantum Hall viscosity. ^3He -A's intrinsic chiral nature may have additional implications such as enabling topological acoustics [59,60] and time-reversal symmetry breaking, without the need for any gain or drive of the medium [61]. Since the static chiral edge current is expected to be carried by Majorana fermions (whereas this is disputed for the dynamic current), balancing or unbalancing the dynamic current with the static current can also provide a signature of Majorana fermions. This could open the door to employing these elusive quasiparticles to study non-Abelian statistics, to simulate exotic fundamental particles and cosmological phenomena [62,63], and may even help enable fault-tolerant topological quantum computing [64–67].

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DATA AVAILABILITY

The data that support the findings of this article are openly available [68], embargo periods may apply.

APPENDIX A: VALIDITY OF APPROXIMATIONS

Perturbative approach for the frequency shift. The perturbative approach used in the estimation of the frequency shift induced through both the dynamic and the static coupling between edge current and surface-sound modes is valid in the limit $v_s \gg v_{\text{dyn}}$. For a given film displacement amplitude η_0 , we find

$$v_{\text{dyn}} \simeq \frac{\hbar \eta_0}{4M d R} \quad \text{and} \quad v_s \simeq \frac{c_s^2 \eta_0}{\Omega d R} = \frac{c_s \eta_0}{\zeta d}. \quad (\text{A1})$$

The ratio between the two is given by $v_{\text{dyn}}/v_s = \frac{\kappa_3 \zeta}{4\pi R c_s}$, which is on the order of 1% for the dimensions considered here, justifying the use of a perturbative approach to consider its effect on the sound waves.

In-plane and out-of-plane chiral currents. The chiral mass current is given by the second term of Eq. (1). For a thin-film

surface-wave resonator, it can be recast as

$$\vec{g}_{\text{dyn}} = \rho_0 \frac{\hbar}{4M h_0} \vec{\nabla} \times (h_0 + \eta) \vec{l}. \quad (\text{A2})$$

Simple vector calculus yields

$$\vec{\nabla} \times (h_0 + \eta) \vec{l} = (h_0 + \eta)(\vec{\nabla} \times \vec{l}) - \vec{l} \times \vec{\nabla} \eta. \quad (\text{A3})$$

We consider, without loss of generality, that $\vec{l}$ is approximately aligned with $\hat{z}$. As $\vec{l}$ remains normal to the film surface, we find

$$\vec{l} = \begin{pmatrix} -\frac{\partial \eta}{\partial r} \\ -\frac{1}{r} \frac{\partial \eta}{\partial \theta} \\ 1 \end{pmatrix} \quad \text{and} \quad \vec{\nabla} \times \vec{l} = \begin{pmatrix} 0 \\ 0 \\ -\frac{\partial}{\partial r} \left(\frac{1}{r} \frac{\partial \eta}{\partial \theta} \right) + \frac{1}{r} \frac{\partial}{\partial \theta} \left(\frac{\partial \eta}{\partial r} \right) \end{pmatrix} \quad (\text{A4})$$

at the film surface. Taking into account that $\vec{l}$ is aligned with the substrate at $z = 0$, whereas it is aligned with the film surface at $z = h_0 + \eta \approx h_0$, we find the anisotropy axis below the film's surface:

$$\vec{l} = \begin{pmatrix} -\frac{\partial \eta}{\partial r} \frac{z}{h_0} \\ -\frac{1}{r} \frac{\partial \eta}{\partial \theta} \frac{z}{h_0} \\ 1 \end{pmatrix}, \quad \vec{\nabla} \times \vec{l} = \frac{1}{h_0} \begin{pmatrix} \frac{\partial \eta}{\partial y} \\ -\frac{\partial \eta}{\partial x} \\ -\frac{\partial}{\partial r} \left(\frac{1}{r} \frac{\partial \eta}{\partial \theta} \right) z + \frac{1}{r} \frac{\partial}{\partial \theta} \left(\frac{\partial \eta}{\partial r} \right) z \end{pmatrix}. \quad (\text{A5})$$

Inserting Eq. (A5) into Eq. (A3) and neglecting terms proportional to η^2 , we find

$$\vec{\nabla} \times (h_0 + \eta) \vec{l} \simeq -2\hat{z} \times \vec{\nabla} \eta. \quad (\text{A6})$$

Inserting the above equation into Eq. (A2), we obtain Eq. (7). The out-of-plane chiral current $\vec{g}_{\text{torque}}$ is suppressed by a factor of h_0/R (which is about 1% for the resonators considered in this work) compared to the in-plane current and will hence be neglected here.

APPENDIX B: TRANSFORMATIONS FOR FIRST SOUND

The method for detection of chiral current proposed in this work could also be applied to density waves in a three-dimensional superfluid ^3He resonator. Using current technology, this approach appears significantly less promising than using surface waves for the following reasons: First, a resonator without boundaries cannot be realized, so that the unpredictable effect of the edge current has to be taken into account when interpreting experimental results. Second, our preliminary calculations show that the expected frequency shift to sound modes in the bulk would be several orders of magnitude smaller than that for surface waves, and hence very hard to detect. Nevertheless, it might be useful to consider this effect in future experiments, including long-term goals such as braiding of Majorana fermions or topological acoustics. A comparison between surface- and bulk-acoustic quantities is shown in Table I.

TABLE I. Transformations for chiral current dynamics between surface waves and bulk density waves. $\hat{n}$ is the surface normal of the three-dimensional resonator and K_1 is the bulk modulus of first sound in superfluid ^3He .

Surface acoustics	Bulk acoustics
Film height perturbation $\eta(\vec{r}, t)$	Density perturbation $\delta\rho(\vec{r}, t)$
Unperturbed film height h_0	Static density ρ_0
Dynamic chiral current [one dimensional (1D)] $\vec{g}_{\text{dyn, 1D}} = -\rho_0 \frac{\hbar}{2M h_0} \hat{z} \times \vec{\nabla} \eta + g_{\text{torque}}$	Dynamic chiral current [two dimensional (2D)] $\vec{g}_{\text{dyn, 2D}} = -\frac{\hbar}{2M} \hat{l} \times \vec{\nabla} \rho + g_{\text{torque}}$
Static chiral current (1D) $\vec{g}_{\text{edge}} = \frac{\hbar \rho_0}{4M \xi} \hat{e}_\theta$	Static chiral current (2D) $\vec{g}_{\text{surface}} = \frac{\hbar \rho_0}{4M \xi} \hat{l} \times \hat{n}$
Surface speed of sound $c_s = \sqrt{c_\sigma^2 + c_3^2 + c_g^2}$	Bulk speed of sound $c_{\text{bulk}} = \sqrt{K_1/\rho_0}$

APPENDIX C: SOUND MODES IN A SUPERFLUID THIN FILM

For the superfluid resonators considered here, the relevant restoring forces are surface tension, van der Waals force, and gravity. The respective speeds of sound are given by

$$c_\sigma = \frac{\zeta}{R} \sqrt{\frac{h_0 \times \sigma_{\text{ST}}}{\rho_0}}, \quad c_3 = \sqrt{\frac{3\alpha_{\text{vdW}}}{h_0^3}}, \quad \text{and} \quad c_g = \sqrt{g_g h_0}, \quad (\text{C1})$$

with $\sigma_{\text{ST}} = 1.56 \times 10^{-4} \text{ J/m}^2$ the surface tension of superfluid helium [69], $\rho_0 = 82 \text{ kg/m}^3$ the density of superfluid ^3He [70], $\alpha_{\text{vdW}} = 3.5 \times 10^{-24} \text{ m}^5/\text{s}^2$ the van der Waals

coefficient of silicon, and $g_g = 9.8 \text{ m/s}^2$ the gravitational acceleration. The total speed of sound is given by $c_s = \sqrt{c_\sigma^2 + c_3^2 + c_g^2}$ and the sound frequency is

$$\nu = \frac{\Omega}{2\pi} = \frac{\zeta c_s}{2\pi R}. \quad (\text{C2})$$

For the resonator dimensions and film thickness considered in this work (radius $R = 10 \text{ }\mu\text{m}$ and film thickness $d = 100 \text{ nm}$), we find $f = 5.1 \text{ kHz}$ for the $(m = 1, n = 1)$ mode and $f = 70 \text{ kHz}$ for the $(m = 6, n = 1)$ mode, with the gravitational contribution being negligible in this case [68].

- [1] C.-K. Chiu, J. C. Y. Teo, A. P. Schnyder, and S. Ryu, Classification of topological quantum matter with symmetries, *Rev. Mod. Phys.* **88**, 035005 (2016).
- [2] B. A. Bernevig, T. L. Hughes, and S.-C. Zhang, Quantum spin Hall effect and topological phase transition in HgTe quantum wells, *Science* **314**, 1757 (2006).
- [3] M. Kono, S. Wiedmann, C. Brune, A. Roth, H. Buhmann, L. W. Molenkamp, X.-L. Qi, and S.-C. Zhang, Quantum spin Hall insulator state in HgTe quantum wells, *Science* **318**, 766 (2007).
- [4] A. A. Burkov, Topological semimetals, *Nat. Mater.* **15**, 1145 (2016).
- [5] Z. Liu, B. Zhou, Y. Zhang, Z. Wang, H. Weng, D. Prabhakaran, S.-K. Mo, Z. Shen, Z. Fang, X. Dai, Z. Hussain, and Y. Chen, Discovery of a three-dimensional topological Dirac semimetal, Na_3Bi , *Science* **343**, 864 (2014).
- [6] S.-Y. Xu, I. Belopolski, N. Alidoust, M. Neupane, G. Bian, C. Zhang, R. Sankar, G. Chang, Z. Yuan, C.-C. Lee, S.-M. Huang, H. Zheng, J. Ma, D. Sanchez, B. Wang, A. Bansil, F. Chou, P. Shibaev, H. Lin, and M. Z. Hasan, Discovery of a Weyl fermion semimetal and topological Fermi arcs, *Science* **349**, 613 (2015).
- [7] X.-L. Qi and S.-C. Zhang, Topological insulators and superconductors, *Rev. Mod. Phys.* **83**, 1057 (2011).
- [8] M. Sato and Y. Ando, Topological superconductors: A review, *Rep. Prog. Phys.* **80**, 076501 (2017).
- [9] A. Yazdani, F. Oppen, B. Halperin, and A. Yacoby, Hunting for Majoranas, *Science* **380**, eade0850 (2023).
- [10] M. Mandal, N. C. Drucker, P. Siriviboon, T. Nguyen, A. Boonkird, T. N. Lamichhane, R. Okabe, A. Chottrattanapituk, and M. Li, Topological superconductors from a materials perspective, *Chem. Mater.* **35**, 6184 (2023).
- [11] L. Fu and C. L. Kane, Superconducting proximity effect and Majorana fermions at the surface of a topological insulator, *Phys. Rev. Lett.* **100**, 096407 (2008).
- [12] V. Mourik, K. Zuo, S. M. Frolov, S. R. Plissard, E. P. A. M. Bakkers, and L. P. Kouwenhoven, Signatures of Majorana fermions in hybrid superconductor-semiconductor nanowire devices, *Science* **336**, 1003 (2012).
- [13] S. Nadj-Perge, I. K. Drozdov, J. Li, H. Chen, S. Jeon, J. Seo, A. H. MacDonald, B. A. Bernevig, and A. Yazdani, Observation of Majorana fermions in ferromagnetic atomic chains on a superconductor, *Science* **346**, 602 (2014).
- [14] Q. L. He, L. Pan, A. L. Stern, E. C. Burks, X. Che, G. Yin, J. Wang, B. Lian, Q. Zhou, E. S. Choi, K. Murata, X. Kou, Z. Chen, T. Nie, Q. Shao, Y. Fan, S.-C. Zhang, K. Liu, J. Xia, and K. L. Wang, Chiral Majorana fermion modes in a quantum anomalous Hall insulator–superconductor structure, *Science* **357**, 294 (2017).
- [15] R. M. Lutchyn, E. P. A. M. Bakkers, L. P. Kouwenhoven, P. Krogstrup, C. M. Marcus, and Y. Oreg, Majorana zero modes in superconductor–semiconductor heterostructures, *Nat. Rev. Mater.* **3**, 52 (2018).
- [16] Microsoft Azure Quantum, M. Aghaee, A. A. Ramirez, Z. Alam, R. Ali, M. Andrzejczuk, A. Antipov, M. Astafev, A.

- Barzegar, B. Bauer, *et al.*, Interferometric single-shot parity measurement in InAs-Al hybrid devices, *Nature (London)* **638**, 651 (2025).
- [17] M. Kayyalha, D. Xiao, R. Zhang, J. Shin, J. Jiang, F. Wang, Y.-F. Zhao, R. Xiao, L. Zhang, K. M. Fijalkowski, P. Mandal, M. Winnerlein, C. Gould, Q. Li, L. W. Molenkamp, M. H. W. Chan, N. Samarth, and C.-Z. Chang, Absence of evidence for chiral Majorana modes in quantum anomalous Hall-superconductor devices, *Science* **367**, 64 (2020).
- [18] P. Yu, J. Chen, M. Gomanko, G. Badawy, E. P. A. M. Bakkers, K. Zuo, V. Mourik, and S. M. Frolov, Non-Majorana states yield nearly quantized conductance in proximatized nanowires, *Nat. Phys.* **17**, 482 (2021).
- [19] M. Valentini, M. Borovkov, E. Prada, S. Martí-Sánchez, M. Botifoll, A. Hofmann, J. Arbiol, R. Aguado, P. San-Jose, and G. Katsaros, Majorana-like Coulomb spectroscopy in the absence of zero-bias peaks, *Nature (London)* **612**, 442 (2022).
- [20] Y. Tanaka, S. Tamura, and J. Cayao, Theory of Majorana zero modes in unconventional superconductors, *Prog. Theor. Exp. Phys.* **2024**, 08C105 (2024).
- [21] D. Castelvecchi, Microsoft claims quantum-computing breakthrough—But some physicists are sceptical, *Nature (London)* **638**, 872 (2025).
- [22] Y. Aoki, Y. Wada, M. Saitoh, R. Nomura, Y. Okuda, Y. Nagato, M. Yamamoto, S. Higashitani, and K. Nagai, Observation of surface Andreev bound states of superfluid ^3He by transverse acoustic impedance measurements, *Phys. Rev. Lett.* **95**, 075301 (2005).
- [23] H. Choi, J. P. Davis, J. Pollanen, and W. P. Halperin, Surface specific heat of ^3He and Andreev bound states, *Phys. Rev. Lett.* **96**, 125301 (2006).
- [24] H. Ikegami, Y. Tsutsumi, and K. Kono, Chiral symmetry breaking in superfluid ^3He -A, *Science* **341**, 59 (2013).
- [25] M. C. Cross, A generalized Ginzburg-Landau approach to the superfluidity of helium 3, *J. Low Temp. Phys.* **21**, 525 (1975).
- [26] N. D. Mermin and P. Muzikar, Cooper pairs versus Bose condensed molecules: The ground-state current in superfluid ^3He -A, *Phys. Rev. B* **21**, 980 (1980).
- [27] R. Combescot and T. Dombre, Superfluid current in ^3He -A at $T = 0$, *Phys. Rev. B* **28**, 5140 (1983).
- [28] V. Ambegaokar, P. G. deGennes, and D. Rainer, Landau-Ginsburg equations for an anisotropic superfluid, *Phys. Rev. A* **9**, 2676 (1974).
- [29] J. E. Avron, R. Seiler, and P. G. Zograf, Viscosity of quantum Hall fluids, *Phys. Rev. Lett.* **75**, 697 (1995).
- [30] N. Read, Non-Abelian adiabatic statistics and Hall viscosity in quantum Hall states and $p_x + ip_y$ paired superfluids, *Phys. Rev. B* **79**, 045308 (2009).
- [31] M. Barkeshli, S. B. Chung, and X.-L. Qi, Dissipationless phonon Hall viscosity, *Phys. Rev. B* **85**, 245107 (2012).
- [32] S. T. Ramamurthy and T. L. Hughes, Patterns of electromagnetic response in topological semimetals, *Phys. Rev. B* **92**, 085105 (2015).
- [33] P. Rao and B. Bradlyn, Hall viscosity in quantum systems with discrete symmetry: Point group and lattice anisotropy, *Phys. Rev. X* **10**, 021005 (2020).
- [34] A. M. R. Schechter, R. W. Simmonds, R. E. Packard, and J. C. Davis, Observation of third sound in superfluid ^3He , *Nature (London)* **396**, 554 (1998).
- [35] H. S. Byun, J. Jeong, K. Kim, S. G. Kim, S.-B. Shim, J. Suh, and H. Choi, Measuring angular momentum of $p_x + ip_y$ topological superfluids: A proposal, *Phys. Rev. B* **98**, 024518 (2018).
- [36] G. I. Harris, D. L. McAuslan, E. Sheridan, Y. Sachkou, C. Baker, and W. P. Bowen, Laser cooling and control of excitations in superfluid helium, *Nat. Phys.* **12**, 788 (2016).
- [37] D. L. McAuslan, G. I. Harris, C. Baker, Y. Sachkou, X. He, E. Sheridan, and W. P. Bowen, Microphotonic forces from superfluid flow, *Phys. Rev. X* **6**, 021012 (2016).
- [38] Y. P. Sachkou, C. G. Baker, G. I. Harris, O. R. Stockdale, S. Forstner, M. T. Reeves, X. He, D. L. McAuslan, A. S. Bradley, M. J. Davis, and W. P. Bowen, Coherent vortex dynamics in a strongly interacting superfluid on a silicon chip, *Science* **366**, 1480 (2019).
- [39] X. He, G. I. Harris, C. G. Baker, A. Sawadsky, Y. L. Sfondla, Y. P. Sachkou, S. Forstner, and W. P. Bowen, Strong optical coupling through superfluid Brillouin lasing, *Nat. Phys.* **16**, 417 (2020).
- [40] A. Sawadsky, R. A. Harrison, G. I. Harris, W. W. Wasserman, Y. L. Sfondla, W. P. Bowen, and C. G. Baker, Engineered entropic forces allow ultrastrong dynamical backaction, *Sci. Adv.* **9**, eade3591 (2023).
- [41] F. M. Ellis and L. Li, Quantum swirling of superfluid helium films, *Phys. Rev. Lett.* **71**, 1577 (1993).
- [42] J. A. Sauls, Surface states, edge currents, and the angular momentum of chiral p -wave superfluids, *Phys. Rev. B* **84**, 214509 (2011).
- [43] M. Stone and R. Roy, Edge modes, edge currents, and gauge invariance in $p_x + ip_y$ superfluids and superconductors, *Phys. Rev. B* **69**, 184511 (2004).
- [44] K. R. Atkins, Third and fourth sound in liquid helium II, *Phys. Rev.* **113**, 962 (1959).
- [45] C. G. Baker, G. I. Harris, D. L. McAuslan, Y. Sachkou, X. He, and W. P. Bowen, Theoretical framework for thin film superfluid optomechanics: Towards the quantum regime, *New J. Phys.* **18**, 123025 (2016).
- [46] S. Forstner, Y. Sachkou, M. Woolley, G. I. Harris, X. He, W. P. Bowen, and C. G. Baker, Modelling of vorticity, sound and their interaction in two-dimensional superfluids, *New J. Phys.* **21**, 053029 (2019).
- [47] I. Rudnick, H. Kojima, W. Veith, and R. S. Kagiwada, Observation of superfluid-helium persistent current by Doppler-shifted splitting of fourth-sound resonance, *Phys. Rev. Lett.* **23**, 1220 (1969).
- [48] F. M. Ellis and H. Luo, Observation of the persistent-current splitting of a third-sound resonator, *Phys. Rev. B* **39**, 2703 (1989).
- [49] G. Volovik, An analog of the quantum Hall effect in a superfluid ^3He film, *Z. Eksp. Teor. Fiz.* **94**, 123 (1988) [*Sov. Phys. JETP* **67**, 1804 (1988)].
- [50] N. Zhelev, T. S. Abhilash, E. N. Smith, R. G. Bennett, X. Rojas, L. Levitin, J. Saunders, and J. M. Parpia, The A-B transition in superfluid helium-3 under confinement in a thin slab geometry, *Nat. Commun.* **8**, 15963 (2017).
- [51] L. V. Levitin, R. G. Bennett, A. Casey, B. Cowan, J. Saunders, D. Drung, T. Schurig, and J. M. Parpia, Phase diagram of the topological superfluid ^3He confined in a nanoscale slab geometry, *Science* **340**, 841 (2013).
- [52] L. V. Levitin, R. G. Bennett, E. V. Surovtsev, J. M. Parpia, B. Cowan, A. J. Casey, and J. Saunders, Surface-induced order

- parameter distortion in superfluid ^3He — B measured by nonlinear NMR, *Phys. Rev. Lett.* **111**, 235304 (2013).
- [53] A. Vorontsov and J. A. Sauls, Spectrum of third sound cavity modes on superfluid 3, *J. Low Temp. Phys.* **134**, 1001 (2004).
- [54] W. W. Wasserman, R. A. Harrison, G. I. Harris, A. Sawadsky, Y. L. Sfondla, W. P. Bowen, and C. G. Baker, Cryogenic and hermetically sealed packaging of photonic chips for optomechanics, *Opt. Express* **30**, 30822 (2022).
- [55] V. Almeida, C. Barrios, R. Panepucci, and M. Lipson, All-optical control of light on silicon chip, *Nature (London)* **431**, 1081 (2004).
- [56] S. M. Meenehan, J. D. Cohen, G. S. MacCabe, F. Marsili, M. D. Shaw, and O. Painter, Pulsed excitation dynamics of an optomechanical crystal resonator near its quantum ground state of motion, *Phys. Rev. X* **5**, 041002 (2015).
- [57] M. Aspelmeyer, T. J. Kippenberg, and F. Marquardt, Cavity optomechanics, *Rev. Mod. Phys.* **86**, 1391 (2014).
- [58] D. Vollhardt and P. Wölfle, *The Superfluid Phases of Helium 3* (Taylor & Francis, London, 1990).
- [59] Z. Yang, F. Gao, X. Shi, X. Lin, Z. Gao, Y. Chong, and B. Zhang, Topological acoustics, *Phys. Rev. Lett.* **114**, 114301 (2015).
- [60] A. Souslov, B. C. van Zuijden, D. Bartolo, and V. Vitelli, Topological sound in active-liquid metamaterials, *Nat. Phys.* **13**, 1091 (2017).
- [61] R. Fleury, D. L. Sounas, C. F. Sieck, M. R. Haberman, and A. Alù, Sound isolation and giant linear nonreciprocity in a compact acoustic circulator, *Science* **343**, 516 (2014).
- [62] J. T. Mäkinen, V. Dmitriev, J. Nissinen, J. Rysti, G. Volovik, A. Yudin, K. Zhang, and V. Eltsov, Half-quantum vortices and walls bounded by strings in the polar-distorted phases of topological superfluid ^3He , *Nat. Commun.* **10**, 237 (2019).
- [63] G. E. Volovik, *The Universe in a Helium Droplet*, International Series of Monographs on Physics (Oxford University Press, Oxford, 2009).
- [64] C. Nayak, S. H. Simon, A. Stern, M. Freedman, and S. Das Sarma, Non-Abelian anyons and topological quantum computation, *Rev. Mod. Phys.* **80**, 1083 (2008).
- [65] B. Lian, X.-Q. Sun, A. Vaezi, X.-L. Qi, and S.-C. Zhang, Topological quantum computation based on chiral Majorana fermions, *Proc. Natl. Acad. Sci. USA* **115**, 10938 (2018).
- [66] D. A. Ivanov, Non-Abelian statistics of half-quantum vortices in p -wave superconductors, *Phys. Rev. Lett.* **86**, 268 (2001).
- [67] S. Das Sarma, M. Freedman, and C. Nayak, Majorana zero modes and topological quantum computation, *npj Quantum Inf.* **1**, 15001 (2015).
- [68] C. G. Baker and S. Forstner, Calculation of the interaction between chiral current and surface waves in the A-phase of helium-3, Zenodo, 2024, <https://doi.org/10.5281/zenodo.10867829>.
- [69] M. Iino, M. Suzuki, A. J. Ikushima, and Y. Okuda, Surface tension of liquid ^3He down to 0.3 K, *J. Low Temp. Phys.* **59**, 291 (1985).
- [70] A. M. R. Schechter, Third sound in superfluid helium-3, Ph.D. thesis, University of California, Berkeley, 1999.