Majorana Fermions Bound at Vortices and Surface of Superfluid $^3$He

T. Mizushima, T. Kawakami, Y. Tsutsumi, M. Ichioka, and K. Machida
Department of Physics, Okayama University, Okayama 700-8530, Japan
E-mail: mizushima@mp.okayama-u.ac.jp

Abstract. We investigate the Majorana fermions bound at vortices and specular surface of superfluid $^3$He A- and B-phases. It is demonstrated that the statistics of the integer vortex which is competitive to the half-quantum vortex is sensitive to the orientation of the magnetic field. We also discuss the surface Andreev bound state in the B-phase, which is expected to behave as Majorana fermions at the zero field limit. Based on the quasiclassical and Bogoliubov-de Gennes theories, we clarify that in the presence of the magnetic field parallel to the surface, the stable configuration of the $\hat{n}$-vector, which reflects the spin-orbit symmetry breaking, traces the point at which the surface state has the maximum gap.

1. Introduction
Recently, much interest has focused on superfluid $^3$He confined to a restricted geometry with a sub-micron thickness, because the non-trivial topological charge is inherent in the bulk region of spin-triplet $p$-wave superfluids [1]. The phase boundary between the non-trivial bulk region and trivial vacuum always gives rise to the topological phase transition, where the quasiparticle spectrum becomes gapless.

It has been revealed that the Andreev bound state confined to integer- and half-quantum vortices (HQV’s) and specular surfaces behave as Majorana zero modes and Majorana fermions, which can be characterized with the self-Hermitian property of the quasiparticle operator and field operator. Due to the self-Hermitian property, the half-quantum vortices accompanied with the Majorana zero mode obey the non-Abelian statistics [2] and the Majorana fermion living on the specular surface gives rise to the Ising anisotropy of spin susceptibility in the superfluid $^3$He B-phase [3, 4, 5, 6].

In this paper, we clarify the interplay of the magnetic field and dipole interaction on the Majorana fermions bound at vortices and surface of superfluid $^3$He. In Sec. 2, we discuss the statistics of integer vortices, since they are energetically competitive to the HQV’s in a slab geometry under high pressure. It is emphasized that the magnetic field plays a critical role on the statistics. It is also known that a slab geometry with a sub-micron thickness [7, 8] emphasizes the surface Andreev bound state (SABS) in the B-phases, which reflects the non-trivial topological invariant defined in the bulk region. In Sec. 3, we examine the thermodynamically stable configuration of the $\hat{n}$-vector in a slab geometry. Based on the quasiclassical theory including the dipole interaction, it is found that in low temperature and high field regimes, the $\hat{n}$-vector
traces the trajectory in which the surface state has a maximum gap. Throughout this paper, we use $k_B = h = 1$.

2. Integer- and half-quantum vortices in the $^3$He A-phase

In this section, we clarify how the magnetic field affects the Majorana zero modes at integer vortex and HQV of superfluid $^3$He A-phase. The order parameter of the A-phase is composed of the $d$- and $\ell$-vectors, where the former (latter) is associated with the spin (orbital) degrees of freedom. In a slab geometry with a thickness $\ell$ much smaller than the dipole coherence length $\xi_D \approx 10 \mu m$ as shown in Fig. 1(a), only the $d$-vector remains active. The resulting order parameter can be described as $d_\mu(k_x + i k_y) e^{i \varphi}$. It is remarkable that the order parameter is invariant under the simultaneous rotation of $\varphi \mapsto \varphi + \pi$ and $d \mapsto -d$. Hence, the discrete symmetry in the A-phase gives rise to the possibility of a vortex state with half-quantum vorticity $\kappa = \frac{1}{2}$, where the vorticity $\kappa$ is defined as $\varphi \mapsto \kappa \varphi$ with the azimuthal angle $\varphi$ in the $x$-$y$ plane.

The HQV is expected to be stabilized in rotating $^3$He confined to a slab, because the free energy of the HQV, $\mathcal{F}_{\text{HQV}}$, can be simply estimated as $\mathcal{F}_{\text{HQV}} / \mathcal{F}_\text{IV} \sim (1 + \rho_{\text{spin}} / \rho_s) / 2$ within the Ginzburg-Landau regime at low pressure, where $\mathcal{F}_\text{IV}$ is the free energy of the integer vortex with $\kappa = 1$. The ratio of superfluid density $\rho_s$ and spin superfluid density $\rho_{\text{spin}}$ is below 1 for all temperatures of $^3$He, that is, $\mathcal{F}_{\text{HQV}} < \mathcal{F}_\text{IV}$. As the pressure increases, however, it is demonstrated in Ref. [9] that the strong coupling effect induces an effectively attractive interaction between HQV’s and makes them unstable toward the integer vortex.

It is important to mention that these energetically competitive vortices can be distinct in the sense of the number of the zero energy states: A single HQV (integer vortex) always accompanies with the odd (even) number of zero energy states. Since the integer vortex is composed of two separated spin sectors, say $\uparrow$ and $\downarrow$, which have same winding numbers, the zero energy states are degenerate with respect to spins, $E_\uparrow = E_\downarrow = 0$. The zero energy states are characterized with the self-Hermitian Majorana operator, $\gamma = \gamma^\dagger$ and $\{\gamma, \gamma\} = 1$, and are found to govern the statistics of their host vortices. In the presence of two zero energy states with $\gamma_{1,2}$, the linear combination introduces the so-called complex fermion as $c \equiv \gamma_1 + i \gamma_2 \neq c^\dagger$. This implies that the eigenstate of the Majorana zero mode with $\gamma_{1,2}$ is expressed as the linear combination of the vacuum and occupation of the single complex fermion $|0\rangle$ and $|1\rangle$. Since the single HQV contains the single $\gamma$, two Majorana zero modes bound at spatially separated vortices form the complex fermion. Hence, the non-locality of the complex fermion gives rise to the non-Abelian statistics of HQV’s [2]. In contrast, the spin degeneracy in integer vortices enables to form the complex fermion within a single vortex as $c = \gamma_1 + i \gamma_1$, leading to the Abelian statistics.

However, we emphasize here that in principle the statistics of integer vortices are controllable as a consequence of the magnetic field $H$ and dipole interaction. The magnetic field energy tends to align $d$ to $d \perp H$, while the dipole interaction favors $d \parallel \ell$. Since the dipole field $H_D$ is approximately $\sim 2 mT$, $d \perp H$ can be realized for $H \gg H_D$. Figure 1(b) shows the relative angle $\theta_d + \theta_H$, calculated with the Ginzburg-Landau theory [9]. Within $H_D \ll H \ll E_F$, it is found that

![Figure 1](image_url)
the magnetic field can split the spin degeneracy of the Majorana zero modes as [9, 10]

\[ E_\sigma \approx -\frac{2\Delta_0}{\pi^{3/2}} \left[ \cos(k_F D_\nu + \pi/4) \pm \frac{\mu_H}{E_F} \sqrt{k_F D_\nu} \sin \left( k_F D_\nu + \frac{\pi}{4} \right) \right] e^{-D_\nu/\xi_\sigma}, \]

where \( \xi_\sigma \) is the coherence length in the spin \( \sigma \) sector. The first term in the right-hand side reflects the quasiparticle tunneling between neighboring vortices with the distance \( D_\nu \) and \( E_F = k_F^2/2M \) is the Fermi energy of \( ^3\text{He} \) with the mass \( M \). With the lift of degeneracy of Majorana zero modes, the braiding rule of the integer vortices can obey the non-Abelian statistics.

3. Surface Andreev bound states in \( ^3\text{He} \)

3.1. Surface Andreev bound states in A- and B-phases

Here, we consider the superfluid \( ^3\text{He} \) confined to a slab geometry with a sub-micron thickness \( D \sim 10\xi_0 < L \), as shown in Fig. 1(a). As mentioned in Sec. 2, the slab geometry within the thickness \( D \ll 10\xi_0 \) energetically favors the A-phase, \( d = \Delta_0(\hat{k}_x + i\hat{k}_y)\hat{z} \), where the \( \mathcal{L} \)-vector, which points the direction of point nodes in the A-phase, is aligned to \( \hat{z} \) in the whole system. As long as we consider the finite size of the slab, the \( \mathcal{L} \)-vectors face to the circumference of the slab at \( x = \pm L/2 \), where \( \mathcal{L} \) is parallel to the circumference. In this situation, the SABS's appear at \( x = \pm L/2 \). The resulting bound states are found in Ref. [11, 12] to have the gapless spectrum. The low-energy spectrum consists of the dispersion linear on the \( k_y \) and the dispersionless with respect to \( k_z \), as \( E(k) = \pm \Delta_0 k_y/k_F \) at \( x = \mp L/2 \). The SABS is insensitive to the orientation of the magnetic field, when \( H \) is much larger than the dipole field \( \sim 2\text{mT} \). Since the flat band structure of the zero energy states can drastically enhance the low-energy density of states, the emergence of the SABS in the A-phase may be observed through the surface specific heat measurement and so on [11].

For the thickness \( D > 10\xi_0 \), the distorted B-phase becomes the ground state in the low pressure and low temperature regime. Without loss of generality, the order parameter is described as

\[ d_\mu(\hat{k}, r) = R_{\mu\nu}(\hat{n}, \varphi)d_{\nu\nu}(r)\hat{k}_\nu \]

where \( R_{\mu\nu}(\hat{n}, \varphi) \) is the matrix for a rotation around the \( \hat{n} \)-axis by the angle \( \varphi \), which reflects the spin-orbit symmetry breaking of the B-phase manifold. At the limit of the zero field, the dipole interaction favors \( \hat{n} \parallel \hat{z} \), leading to the diagonal representation, \( d_\mu(\hat{k}, r) = \delta_{\mu\nu}d_{\nu\nu}(r)\hat{k}_\nu \). Within this representation, the surface Andreev bound state is derived from the BdG equation in Refs. [3, 4, 5] and the Ising anisotropy of spin susceptibility is associated with the Majorana nature of the surface state.

Here, we derive the surface Andreev bound state from the general form in Eq. (2). First, using the relation \( R_{\mu\nu}R_{\eta\mu} = R_{\nu\eta}R_{\mu\nu} = \delta_{\mu\eta} \), the BdG equation within the Andreev approximation reduces to

\[ \begin{bmatrix} -iv_F \cos \theta \partial_z \tau_z - \mu_H \hat{\ell}_\mu(\hat{n}, \varphi)\tilde{\sigma}_\mu \tau_0 + d_{\mu\mu}(z)\sigma_\mu \hat{k}_\mu \tau_\nu \end{bmatrix} \Phi_k(z) = E\Phi_k(z) \]

where \( v_F \) is the Fermi velocity, \( k \approx k_F(\cos \phi \sin \theta, \sin \phi \sin \theta, \cos \theta) \), \( \mu = x, y, z \), and \( \tau_\mu (\sigma_\mu) \) is the Pauli matrices in the particle-hole (spin) space. Here, we define \( \hat{\ell}_\mu(\hat{n}, \varphi) \equiv H_\nu R_{\mu\nu}(\hat{n}, \varphi)/H \) and \( \tilde{\sigma}_\mu \equiv \tilde{\sigma}_\mu(\hat{n}, \varphi) = \sigma_\nu R_{\nu\mu}(\hat{n}, \varphi) \), which were first introduced in Ref. [13]. Then, using the unitary matrix which transforms \( \tilde{\sigma}_\mu \) to \( \sigma_\mu \) and \( M \equiv (\tau_0 \sigma_x e^{i\theta \sigma_z} + \tau_\sigma \sigma_x e^{i\theta \sigma_z})/\sqrt{2} \) with \( \vartheta = \frac{\theta}{2} + \frac{\pi}{4} \), Eq. (3) reduces to

\[ \begin{bmatrix} -iv_F \cos \theta \partial_z \tau_z - \mu_H \hat{\ell}_\mu(\hat{n}, \varphi)\sigma_\mu \tau_0 - \Delta_{\perp} \cos \theta \sigma_\mu \tau_z - \Delta_{\parallel} \sin \theta \sigma_0 \tau_y \end{bmatrix} M\Phi_k(z) = EM\Phi_k(z) \]

which is now block-diagonalized to two sectors, \( (\Phi_1, \Phi_3) \) and \( (\Phi_2, \Phi_4) \). Here, we introduce \( \sigma' = [-\sigma_z \sin \phi + \sigma_y \cos \phi, \sigma_z \cos \phi + \sigma_y \sin \phi, \sigma_x] \), \( \Delta_{\perp} \equiv d_{xx} \) and \( \Delta_{\parallel} \equiv d_{xx} - d_{yy} \).
From the reduced equation (4) with the specular boundary condition \( \Phi(z=0)=0 \) and \( H=0 \), the gapless cone spectrum bound on the surface can be derived as

\[
E(k_x, k_y) = \pm \frac{\Delta_0}{k_F} \sqrt{k_x^2 + k_y^2},
\]

where \( \Delta_0 = \Delta_\perp \equiv \Delta_0 \). This expression is independent of the orientation of \( \hat{n} \) and the angle \( \varphi \).

In the presence of the finite magnetic field \( H \neq 0 \), the energy gap of the surface state depends on the \( \hat{n} \)-vector as

\[
\min |E(k_x, k_y)| = \mu H \left| \hat{\varepsilon}_z(\hat{n}, \varphi) \right| = \mu H |R_{\mu z}(\hat{n}, \varphi)|. 
\]

This implies the Ising anisotropy \([3, 4]\) that the energy of the SABS depends on the orientation of \( \hat{n} \). As seen in Sec. 3.1, the SABS is sensitive to the orientation of the \( \hat{n} \)-vector and the angle \( \varphi \). For the microscopic determination of the orientation, it is necessary to take account of the dipole interaction into the gap equation. Hence, we here present the self-consistent formalism based on the microscopic determination of the orientation, it is necessary to take account of the dipole interaction into the gap equation. Hence, we here present the self-consistent formalism based on the quasiclassical theory, including the dipole interaction.

The superfluid phases in \(^3\)He under low pressure are well described with the quasiclassical Eilenberger equation \([14]\)

\[
[i\omega_n \hat{\tau}_3 - \hat{S}(\hat{k}, \mathbf{r}), \hat{g}(\hat{k}, \mathbf{r}; i\omega_n)] + ivF \hat{k} \cdot \nabla \hat{g}(\hat{k}, \mathbf{r}; i\omega_n) = 0, \quad \hat{S} = \left[ \begin{array}{cc} \sigma_j \nu_j & i\sigma_\mu \sigma_y d_\mu \sigma_j \nu_j \\ i\sigma_y \sigma_\mu d_\mu & \sigma_j \nu_j \end{array} \right],
\]

where \( j = 0, x, y, z \) and \( \omega_n = (2n + 1)\pi T \) is the Matsubara frequency with \( n \in \mathbb{Z} \). The magnetic field is taken into \( \hat{\nu} \equiv \mu H \text{diag}[\sigma_\mu, \sigma_y]/(1 + F_0^2) \). Equation (7) governs the evolution of the quasiclassical Green’s functions parameterized in the particle-hole space as

\[
\hat{g}(\hat{k}, \mathbf{r}; i\omega_n) = \left[ \begin{array}{cc} \sigma_j g_j(\hat{k}, \mathbf{r}; i\omega_n) & i\sigma_\mu \sigma_y f_\mu(\hat{k}, \mathbf{r}; i\omega_n) \\ i\sigma_y \sigma_\mu f_\mu & \sigma_j g_j \end{array} \right].
\]

Equation (7) coupled with the normalization condition \( \hat{g}^2 = -\pi^2 \hat{1} \) can be numerically solved with the Riccati parametrization \([15, 16]\).

The quasiclassical self-energy \( \hat{S} \) in Eq. (7) consists of the pair potential \( d_\mu \) and Fermi liquid correction \( \nu_j \). The diagonal self-energies are associated with the quasiclassical Green’s functions as

\[
\nu_j(\hat{k}, \mathbf{r}) = \langle A^{U}(\hat{k}, \hat{k}'), g_j(\hat{k}', \mathbf{r}; i\omega_n) \rangle_{k' \omega_n},
\]

where \( \langle \cdots \rangle_{k \omega_n} \) indicates the Matsubara sum with the cutoff \( E_c = 20\pi T_c \) and the Fermi surface average. The coefficients \( A^{(\mu \sigma \rho)} = A^\sigma \) and \( A^{(\nu \rho \sigma)} = A^\sigma \) are the quasiparticle scattering amplitude and are parametrized with the Landau’s Fermi liquid parameter \( F_0^{\mu \sigma} \) \([14]\), where \( F_0^0 = 9.3 \), \( F_0^3 = 5.39 \), \( F_0^{\perp} = -0.695 \), and \( F_0^{\perp} = -0.5 \).

The order parameter of spin-triplet \( p \)-wave superfluids can be decomposed into nine components as \( d_{\mu \nu}(\hat{k}, \mathbf{r}) = d_{\mu \nu}(\mathbf{r}) \hat{k}_\nu \), where the gap equation \( d_\mu(\hat{k}, \mathbf{r}) = \langle V(\hat{k}, \hat{k}') f_\mu(\hat{k}', \mathbf{r}; i\omega_n) \rangle_{k' \omega_n} \) gives the relation between the order parameter and the quasiclassical Green’s functions. In the superfluid \(^3\)He, the effective pairing interaction \( V(\hat{k}, \hat{k}') \) is composed of the BCS interaction \( 3\hat{k} \cdot \hat{k}' \) and the dipole interaction \([17]\). This can be summarized as follows:

\[
d_{\mu \nu}(\mathbf{r}) = (3|g| - \tilde{g}_D) \left\langle \hat{k}_x f_\mu \right\rangle_{k \omega_n} - 3\tilde{g}_D \left\langle \delta_{\mu x} \left\langle \hat{k} \cdot f \right\rangle \right\rangle_{k \omega_n} + \epsilon_{\mu \gamma} \left( \left\langle \hat{k} \cdot f \right\rangle \right)_{k \omega_n},
\]
3.3. Microscopic determination of $\mathbf{n}$-vectors

Here, we solve Eqs. (7) and (9) self-consistently, where the specular boundary condition on $z=0$ and $D$ is imposed on $\hat{g}$. In the presence of the dipole interaction, all the nine components of $d_{\mu\nu}$ interplay with each other, so that the self-consistent equations yield slow convergence. To overcome this difficulty, we solve the self-consistent equations for a given $\hat{d}$ of the resulting $\mathbf{r}$. As seen in Eq. (2), the remaining order parameters reduce to only three components $d_{\mu\nu}(z)$. The detailed procedure is as follows: (i) First, we solve Eq. (7) for a given self-consistent potential. (ii) The three order parameter $d_{\mu\nu}(z)$ is obtained from the resulting $\hat{g}$ through $d_{\mu\nu}(\mathbf{r}) = R_{\mu\nu}(\hat{g}, \varphi)d_{\mu\nu}(\hat{r})\hat{k}_\eta$, where $d_{\mu\nu}(\mathbf{r})$ is calculated from Eq. (9). (iii) Then, Eq. (7) is solved again with the resulting $d_{\mu\nu}(\hat{r})$ and a given $\hat{n}$ and $\varphi$. This iteration implies that $\hat{n}$ and $\varphi$ are fixed as spatially uniform constants during the iteration, which can be justified from the fact of $\xi_0 \gg D$. The thermodynamic potential $\delta\Omega[\hat{g}]$ within the quasiclassical approximation is described in Refs. [18, 19]. To this end, we determine the stable configuration of the $\mathbf{n}$ and $\varphi$ which minimizes $\delta\Omega$ under a fixed $D$, $H$ and $T=0.2T_0$.

In Fig. 2(a), we present the size- and field-dependence of $\varphi$ for $H \parallel \hat{z}$ which minimizes $\delta\Omega$. Since the dipole interaction tends to align $\mathbf{n}$ to the direction normal to the surface, the configuration of $\hat{n} \parallel \hat{z}$ is the ground state for $H \parallel \hat{z}$. In the thermodynamic limit of $D \gg \xi_0$, the gap equation (9) leads to the simple relation for $\varphi$, the so-called Leggett angle, $\cos \varphi = -\frac{1}{2} \frac{\Delta_{\mu\nu}}{\Delta_{\mu\nu}}$ [17]. As seen in Fig. 2(a), however, the finite size effect deviates $\varphi$ from the Leggett angle. Figure 2(b) shows the spatial profiles of $d_{\mu\nu}(z)$ with the resulting $\varphi$ and $\hat{n} \parallel \hat{z}$ at $D=20\xi_0$, $T=0.2T_0$, and $\mu H/\pi T_0 = 0.0305$. In this orientation of $H$, as discussed in Sec. 3.1, the SABS is always accompanied with the finite energy gap $\mu H$, which considerably enhances the magnetization density on the surface [3, 4, 6, 19].

Now, we microscopically determine the orientation of $\hat{n}$ in a slab geometry with a magnetic field parallel to the surface $H \parallel \hat{x}$. As seen in Sec. 3.1, in the presence of a parallel field, the low-energy spectrum of the SABS is sensitive to the orientation of $\hat{n}$ and $\varphi$. Figure 3(a) shows the minimum energy gap of the SABS at $D=20\xi_0$ and $\mu_B H=0.0305\pi T_0$, where we suppose that $\varphi$ is insensitive to the orientation of $H$. In Fig. 3(b), we present the thermodynamic potential landscape obtained from the self-consistent solutions of Eqs. (7) and (9) with a given $\hat{n}$ and $\varphi$. It is demonstrated from these figures that the thermodynamic potential landscape has the bottom at which the SABS has the maximum energy gap $\mu H$. As $H$ increases, the maximum point of the surface energy gap shifts to the $\hat{n}_y$ direction, because $\cos \varphi$ approaches 0. Then, it

![Figure 2](image-url)

**Figure 2.** (a) Field-dependence of the stable $\varphi$ for $H \parallel \hat{z}$ at $T=0.2T_0$ and $D=20\xi_0$. The first-order phase transition between B- and A-phases occurs at $\mu H=0.09\pi T_0 \approx 0.36T$. Here $\hat{n}$ is aligned to the $\hat{z}$-axis. The inset of (a) shows the $D$-dependence at $H=0$, where the broken line indicates the relation $\cos \varphi = -\frac{1}{2} \frac{\Delta_{\mu\nu}}{\Delta_{\mu\nu}}$ valid for $D \gg \xi_0$. (b) Spatial profile of $d_{\mu\nu}(z)$ with the resulting $\varphi$ and $\hat{n}=\hat{z}$ at $\mu H=0.0305T_0$, where $\xi_0 \equiv v_F / \pi T_0$ is the coherence length.
Figure 3. (a) Energy gap $\min |E|/\mu H$ of the SABS described in Eq. (6) when $H \parallel \hat{x}$ and $\mu H/\pi T_{c0} = 0.0305$, where $\varphi$ is same as that in Fig. 2(a). (b) Thermodynamic potential $\delta \Omega$ on the unit sphere of $\hat{n}$, obtained from the self-consistent calculation with parameters same as (a). The symbol “×” in both figures denotes the bottom in the thermodynamic potential landscape.

is found that the bottom in the landscape traces the trajectory on $\hat{n}$-sphere where the SABS has the maximum energy gap. In contrast, the $\hat{n}$-vector points to $\hat{z}$ when the magnetic field energy becomes competitive to the dipole interaction.

4. Concluding remarks

Here, we have investigated the interplay of the orientation magnetic field and the dipole interaction on vortices and surface Andreev bound states of superfluid $^3$He. It has been demonstrated that the statistics of the integer vortex associated with the Majorana zero modes is sensitive to the orientation of the magnetic field. In addition, the interplay of the magnetic field and dipole interaction drastically changes the surface Andreev bound states. We have clarified that at low temperatures, the ground state configuration of $\hat{n}$ traces the point at which the maximum energy gap in the surface Andreev bound states is open. However, The weak field comparable to the dipole field $\sim 2mT$ deviates the configuration of $\hat{n}$ from the maximum gap point of the surface state, where the topological charge and Majorana fermion may be realized. The detailed study in this regime remains as a future problem.

This work was supported by JSPS and the MEXT KAKENHI (No. 22103002).

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