Transverse magnetic field and chiral-nonchiral transition in vortex states for nearly $B \parallel ab$ in chiral $p$-wave superconductors

Masahiro Ishihara, Yujiirou Amano, Masanori Ichioka and Kazushige Machida
Department of Physics, Okayama University, Okayama 700-8530, JAPAN
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On the basis of Eilenberger theory, we study the vortex state when a magnetic field is applied nearly parallel to the $ab$ plane in a chiral $p$-wave superconductor with a large anisotropy ratio of $ab$ and $c$, as in Sr$_2$RuO$_4$. We quantitatively estimate the field dependence of the pair potential, magnetization, and flux line lattice form factor, and study the transition from the chiral $p$-state at low fields to the nonchiral $p_d$ state at high fields. Even for exactly parallel fields to the $ab$ plane, transverse fields exist in the chiral state. The chiral-nonchiral transition disappears when the magnetic field orientation is tilted within 1° from the $ab$ plane. This may be a reason why the experimental detection of this transition is difficult.

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I. INTRODUCTION

Chiral $p$-wave superconductivity attracts much attention as one of representatives of topological superconductors. The chiral $p$-wave superconductivity with the pairing function $p_{\pm} = p_{\pm} \pm i p_{\parallel}$ is a possible pairing state, when the $p$-wave pairing interaction works instead of the conventional $s$-wave pairing. The pairing function $p_{\pm}$ breaks time-reversal symmetry, inducing spontaneous magnetic fields observed by $\mu$SR experiments. We also expect that Majorana states are accommodated at vortex and surfaces in chiral $p$-wave superconductors. This type of pairing is realized in the A phase of superfluid $^3$He, and is a candidate for the superconducting phase of Sr$_2$RuO$_4$. However, there remain mysteries for the pairing symmetry of Sr$_2$RuO$_4$, since we have not observed some typical phenomena expected in chiral $p$-wave superconductors. For example, when a magnetic field $\mathbf{B}$ is applied in the orientation $\mathbf{B} \parallel ab$, theoretically we expect the transition from the chiral $p_{\pm}$-wave state to the nonchiral $p_{d}$-wave state (when $\mathbf{B} \parallel y$) at high fields. That is, at low fields, the free energy of the chiral $p_{\pm}$-wave state is lower than that of nonchiral $p_d$- or $p_{d^*}$-wave states, because the latter nonchiral states have vertical line nodes. On the other hand, the nonchiral state is realized at high fields, because the upper critical field $H_{c2}$ of the $p_d$ state is higher than that of the chiral $p_{\pm}$-wave state when $\mathbf{B} \parallel y$. While the chiral-nonchiral transition was suggested by experiments of the magnetization curve, this transition was not observed in other experimental methods. There were discussions in that the double transition near $H_{c2}$ corresponds to the chiral-nonchiral transition.

On the other hand, in superconductors with uniaxial anisotropy, transverse magnetic fields appear in the vortex state when the field orientation is tilted from the $ab$ plane. This transverse field is detected by the spinflip scattering of the small angle neutron scattering (SANS) in the vortex states, as demonstrated in YBa$_2$Cu$_3$O$_{7-\delta}$. Recently, the spin flip SANS by the transverse field was reported in Sr$_2$RuO$_4$. Therefore, the quantitative theoretical estimate of the transverse field is expected. It is also important to find new phenomena by the contribution of Cooper pair’s angular momentum $L_{z}/\hbar = \pm 1$ of the $p_{\pm}$-wave pairing.

The purpose of this study is to establish quantitative theoretical estimations of the vortex structure in chiral $p$-wave superconductors when a magnetic field is applied exactly $\mathbf{B} \parallel ab$, and when the field orientation is slightly tilted from the $ab$ plane. On the basis of Eilenberger theory by which we can quantitatively calculate the spatial structure and the physical quantities of the vortex states, we will clarify behaviors of the chiral-nonchiral transition and the transverse field structure as a function of a magnetic field $\mathbf{B}$.

II. FORMULATION BY EILENBERGER THEORY

As a model of the Fermi surface, we use a quasi-two dimensional Fermi surface with a rippled cylinder shape. The Fermi velocity is assumed to be $v = (v_x, v_y, v_z) \propto (\cos \phi, \sin \phi, \tilde{v}_z \sin \phi)$ at $\mathbf{p} = (p_a, p_b, p_c) \propto (p_x \cos \phi, p_y \sin \phi, p_z)$ on the Fermi surface. We consider a case $\tilde{v}_z = 1/60$, producing large anisotropy ratio of coherence lengths, $\gamma \equiv \xi_c/\xi_b \sim (v_x^2)^{1/2}/(v_y^2)^{1/2} \sim 1/60$, where $(\cdot \cdot \cdot)_p$ indicates an average over the Fermi surface. The magnetic field is tilted within 1° from the $ab$ plane. Since we set the $z$ axis to the vortex line direction, the coordinate $(x, y, z)$ for the vortex structure is related to the crystal coordinate $(a, b, c)$ as $(x, y, z) = (a \cos \theta + c \sin \theta, c \cos \theta - b \sin \theta)$ with $\theta = 90^\circ \sim 89^\circ$.

In a chiral $p$-wave superconductor, the pair potential takes the form,

$$\Delta (\mathbf{p}, \mathbf{r}) = \Delta_+(\mathbf{r})\phi_+(\mathbf{p}) + \Delta_- (\mathbf{r})\phi_- (\mathbf{p})$$

with the pairing functions $\phi_{\pm} (\mathbf{p}) = (p_a \pm ip_b)/p_W = e^{\pm \phi}$. $\Delta_+(\mathbf{r})$ describes the vortex structure as a function of $\mathbf{r}$ (the center of mass coordinate of the pair). In our study, $\Delta_- (\mathbf{r})$ is a main component and $\Delta_+(\mathbf{r})$ is a passive
component induced around a vortex.\(^{15,16}\) At a zero field, \(\Delta_+(\mathbf{r}) = 0\). When we consider the \(p_x\) and \(p_y\) orbital components, the pair potential is decomposed as \(\Delta(\mathbf{p}, \mathbf{r}) = \Delta_x(\mathbf{r})\phi_x(\mathbf{p}) + \Delta_y(\mathbf{r})\phi_y(\mathbf{p})\) with \(\phi_x(\mathbf{p}) = \sqrt{2}p_x = \sqrt{2}\cos \phi\) and \(\phi_y(\mathbf{p}) = \sqrt{2}p_y = \sqrt{2}\sin \phi\).

Using the anisotropic ratio \(\Gamma_{\theta} \equiv \xi_y/\xi_x \sim \langle v^2 \rangle^{1/2}/\langle v_x^2 \rangle^{1/2} \sim (\cos^2 \theta + \gamma^{-2} \sin^2 \theta)^{-1/2}\), we set the unit vectors of the vortex lattice as \(\mathbf{u}_1 = c(\alpha/2, -\sqrt{3}/2)\) and \(\mathbf{u}_2 = c(\alpha/2, \sqrt{3}/2)\) with \(c^2 = 2\phi_0/(\sqrt{3}\alpha B)\) and \(\alpha = 3\Gamma_{\theta}/2\) as shown in Fig. 1(a). \(\phi_0\) is the flux quantum, and \(B\) is the flux density. As shown in Fig. 1(b), the unit vectors in the reciproc space are given by \(\mathbf{q}_1 = (2\pi/c)(1/\alpha, -1/\sqrt{3})\) and \(\mathbf{q}_2 = (2\pi/c)(1/\alpha, 1/\sqrt{3})\), where spots of the SANS appear.

Quasiclassical Green’s functions \(f(\omega_n, \mathbf{p}, \mathbf{r}), f^\dagger(\omega_n, \mathbf{p}, \mathbf{r}), g(\omega_n, \mathbf{p}, \mathbf{r})\) in the vortex lattice states are obtained by solving the Riccati equation, which is derived from the Eilenberger equation

\[
\begin{align*}
\{\omega_n + \tilde{\mathbf{v}} \cdot (\nabla + i\mathbf{A})\} f &= \Delta g, \\
\{\omega_n - \tilde{\mathbf{v}} \cdot (\nabla - i\mathbf{A})\} f^\dagger &= \Delta^* g.
\end{align*}
\]

in the clean limit, with a normalization condition \(g = (1 - f f^\dagger)^{1/2}\) and the Matsubara frequency \(\omega_n\).\(^{15-18}\) That is, we have scaled the length, temperature, magnetic field, and energies in units of \(\xi_0, T_c, B_0\), and \(\pi\lambda/T_c\) respectively, where \(\xi_0 = \hbar v_F/2\pi\lambda T_c\) \(B_0 = \phi_0/2\pi\xi_0^2\). The vector potential \(\mathbf{A} = \frac{1}{2}\mathbf{B} \times \mathbf{r} + \mathbf{a}(\mathbf{r})\) is related to the internal field as \(\mathbf{B}(\mathbf{r}) = \nabla \times \mathbf{A} = (B_x(\mathbf{r}), B_y(\mathbf{r}), B_z(\mathbf{r}))\) with \(\mathbf{B} = (0, 0, B)\), \(B_x(\mathbf{r}) = B + b_x(\mathbf{r})\) and \((B_x, B_y, b_z) = \nabla \times \mathbf{a}\).

The spatial averages of \(B_x, B_y\), and \(b_z\) are zero.\(^{15,16}\)

We calculate \(\Delta(\mathbf{p}, \mathbf{r})\) by the gap equation

\[
\Delta(\mathbf{p}, \mathbf{r}) = \lambda_0 2T \sum_{\omega_n > 0} \langle \phi^*_n(\mathbf{p}) f \rangle_p,
\]

where \(\lambda_0 = N_0 g_0\) is the dimensionless \(p\)-wave pairing interaction in the low-energy band \(\omega_n \leq \omega_c\), defined by the cutoff energy \(\omega_c\) as \(1/\lambda_0 = \ln T + 2T \sum_{\omega_n > 0} \omega_n^{1/2}\). We carry out calculations using the cutoff \(\omega_c = 20k_BT_c\). The current equation to obtain \(\mathbf{A}\) is given by

\[
\mathbf{j}(\mathbf{r}) = \nabla \times (\nabla \times \mathbf{A}) = -\kappa^{-2} 2T \sum_{\omega_n > 0} \langle \mathbf{v} \mathbf{I} \mathbf{m} g \rangle_p.
\]

The Ginzburg-Landau (GL) parameter \(\kappa\) is the ratio of the penetration depth to coherence length for \(\mathbf{B} \parallel c\), and set to be \(\kappa = 2.7\) appropriate to \(\text{Sr}_2\text{RuO}_4\).\(^\dagger\) The case of effective GL parameter \(\kappa_0 \sim \kappa T_\theta\) for a field orientation \(\theta\) is reproduced by the anisotropy of \(\tilde{\mathbf{v}}\) in Eq. (3). Iterating Eqs. (2)\(^{-}\) at \(T = 0.5T_c\), we obtain self-consistent solutions of \(\Delta_{\pm}(\mathbf{r})\), \(\mathbf{A}(\mathbf{r})\), and quasiclassical Green’s functions.

### III. EXACTLY PARALLEL FIELD TO THE BASAL PLANE

First, we study the vortex states for exactly \(\mathbf{B} \parallel ab\) (\(\theta = 90^\circ\)). In Fig. 2 we show the calculated spatial structures within a unit cell of the vortex lattice at low and high fields. The main component \(\Delta_{-}(\mathbf{r})\) has a winding 1 of the phase at the vortex center, where the amplitude \(|\Delta_{-}(\mathbf{r})|\) in Fig. 2(a) is suppressed. At low fields, the vortex core is localized at the center. At high fields the vortex core contribution becomes important in the properties of the vortex states, since the inter-vortex distances become shorter with increasing fields. As a property of chiral \(p\)-wave superconductors, the opposite chiral component \(\Delta_{+}(\mathbf{r})\) also appears where the main chiral component \(\Delta_{-}(\mathbf{r})\) has spatial modulations around vortex cores.\(^{15,19}\) The amplitude of the induced component \(\Delta_{+}(\mathbf{r})\) is presented in Fig. 2(b). It appears locally around the vortex core at a low field \(B\) = 2. With increasing fields, since the inter-vortex distances become shorter, \(\Delta_{+}(\mathbf{r})\) of neighbor vortex cores overlap with each other, as shown in panels for \(B = 8\) and 16. With further increasing \(\mathbf{B}\), the amplitude of \(\Delta_{+}(\mathbf{r})\) is reduced to \(\Delta_{+}(\mathbf{r}) = |\Delta_{-}(\mathbf{r})|\), as shown in a panel for \(B = 20\) in Fig. 2(b). This indicates disappearance of \(\Delta_{+}(\mathbf{r})\) by the chiral-nonchiral transition from the chiral \(p\)-wave state to the nonchiral \(p_y\)-wave state.

The \(z\)-component of the internal field, \(B_z(\mathbf{r})\), has a conventional spatial structure of the vortex lattice also for \(\mathbf{B} \parallel ab\) in chiral \(p\)-wave superconductors, if the length is re-scaled by the effective coherence length in each direction. As shown in Fig. 2(c), \(B_z(\mathbf{r})\) has a peak at a vortex center, and decreases as a function of radius from the center. We note that the transverse components \(B_x(\mathbf{r})\) and \(B_y(\mathbf{r})\) appear even when exactly \(\mathbf{B} \parallel ab\) in the chiral \(p_-\) state at low fields. This is unconventional behavior due to the contribution of the internal angular momentum \(L_z\) of the chiral pairing function. The transverse components vanish in nonchiral \(p_y\) states at high fields.
To see the behaviors of the chiral-nonchiral transition, we plot the amplitudes of each component of the pair potential as a function of $\bar{B}$ in Fig. 2(a). With increasing $\bar{B}$, the $p_-$ wave component decreases and the $p_+$ wave component increases. After the chiral-nonchiral transition at $B > B^* \sim 18$, the amplitudes of $p_+$ and $p_-$ are the same. If we see the pair potential in the decomposition of $p_x$ and $p_y$, with increasing $\bar{B}$, the $p_y$ component decreases toward zero at $H_{c2}$ and the $p_x$ component decreases toward zero at $B^*$. $\Delta_x = 0$ at $B > B^*$ by the chiral-nonchiral transition. In Fig. 2(a), we also show the case of conventional $s$-wave pairing. Compared with the $s$-wave case, $H_{c2}$ is enhanced in the $p_y$-wave state. This comes from the fact that $H_{c2}$ is enhanced when the pairing function has a horizontal line node relative to the field direction.

To discuss the $\bar{B}$-dependence of the internal field distribution, we consider flux line lattice (FLL) form factors $F(q_{h,k}) = (F_x(h,k), F_y(h,k), F_z(h,k))$ calculated as Fourier transformation of the internal field distribution, $\mathbf{B}(r) = \sum_{h,k} F(q_{h,k}) \exp(iq_{h,k} \cdot r)$ with the wave vectors $q_{h,k} = h\mathbf{q}_1 + k\mathbf{q}_2$. $h$ and $k$ are integers. The $z$-component $|F_z(h,k)|^2$ from $B_z(r)$ gives the intensity of spots in the conventional non-spinflip SANS experiments. The transverse components, $|F_x(h,k)|^2 = |F_x(h,k)|^2 + |F_y(h,k)|^2$, is accessible by the spin-flip SANS experiments. In Fig. 3(b), we see exponential decays of $|F_z(h,k)|^2$ as a function of $\bar{B}$, as in the conventional behavior of the vortex states, since we do not take care of the Pauli-paramagnetic effect. The transverse components $|F_x(1,0)|^2$ and $|F_y(1,1)|^2$ appear only in the chiral states at $\bar{B} < B^*$. From the stripe pattern in the spatial structure of $B_y(r)$ as in Fig. 2(c), the main spot of $|F_y(1,1)|^2$ is at $(h,k) = (1,0)$. The intensity of $|F_x(1,0)|^2$ is much smaller than $|F_y(1,1)|^2$.
IV. FIELD ORIENTATION TILTED FROM THE BASAL PLANE

Next, we discuss the vortex states when the magnetic field is slightly tilted from the $ab$ plane as $89^\circ \leq \theta < 90^\circ$. The vortex states in the $p_-$-wave domain, where $\Delta_-(\mathbf{r})$ is the main component, has lower free energy than that in the $p_+$-wave domain where $\Delta_+(\mathbf{r})$ is the main component. This is because the field orientation lifts up the degeneracy of the $p_-$ and $p_+$-wave domains.\cite{15, 16}

Therefore, we study the stable $p_-$-wave domain case here.

When $\theta = 89^\circ$, around the vortex core of the main component $\Delta_-(\mathbf{r})$, the opposite chiral component $\Delta_+(\mathbf{r})$ is also induced as presented in Fig. 2(a). There, the spatial pattern of $|\Delta_-(\mathbf{r})|$ at $\bar{B} = 2$ is similar to that of $\mathbf{B} \parallel c$ case as rather than that of $\mathbf{B} \parallel ab$ in Fig. 2(b). $|\Delta_+(\mathbf{r})|$ at a high field $\bar{B} = 20$ keeps similar spatial structure to that of $\mathbf{B} \parallel ab$ case with $\bar{B} = 8$ in Fig. 2(b). Thus $|\Delta_+(\mathbf{r})| \neq |\Delta_-(\mathbf{r})|$ even at high fields, indicating that the nonchiral state with $\Delta_+(\mathbf{r}) = 0$ does not realize. To see the disappearance of the chiral-nonchiral transition, we study the $\bar{B}$-dependence of each component of the pair potential for $\theta = 89.9^\circ$, $89.5^\circ$, and $89.0^\circ$. As seen from the curve for $89.9^\circ$ in Fig. 2(a), even if the field orientation is tilted by $0.1^\circ$, the chiral-nonchiral transition changes to a crossover behavior. At high fields, small differences between the $p_-$ and $p_+$ components still exist. Thus, in Fig. 2(b), a small amplitude of $p_+$-wave component survives up to $H_{c2}$. Further tilting the field orientation to $89.5^\circ$ and $89.0^\circ$, the crossover behaviors are smeared, and we can not see the remnant of the chiral-nonchiral transition anymore. The amplitude of the $p_+$-wave component monotonically decreases toward $H_{c2}$.

The chiral-nonchiral transition is reflected by the magnetization curve, $M = \bar{B} - H$ as a function of $\bar{B}$. From the selfconsistent solutions we obtain the relation of $\bar{B}$ and the external field $H$ as

$$H = \bar{B} + \left( \langle B_z(\mathbf{r}) - \bar{B} \rangle \right)_r / \bar{B}$$
been observed yet. This may be because the experimental situation of exactly $\vec{B} \parallel ab$ is difficult to be realized. Our study shows that the chiral-nonchiral transition vanishes by tilting the field orientation within $1^\circ$.

In Sr$_2$RuO$_4$, when $\vec{B} \parallel ab$, $H_{c2}$ is suppressed and changes to the first order phase transition. Our simple formulation in this work does not include the mechanism for the suppression of $H_{c2}$, such as a Pauli-paramagnetic-like effect. The study for this $H_{c2}$ behavior belongs to future works.

Both in the $\vec{B}$-dependence of $|F_x(1,0)|^2$ in Fig. 7(a) and $|F_z(1,1)|^2$ in Fig. 7(b), with decreasing $\theta$ from $90^\circ$, $|F_z|^2$ becomes larger at low fields, reflecting the decrease of the effective GL parameter $\kappa_\theta$. Roughly $|F_z| \propto (\kappa_\theta)^{-1}$ from Eq. (4). On the other hand, $|F_x|^2$ becomes smaller at high fields, because $H_{c2}$ decreases by the decrease of $\theta$. As shown in Figs. 7(b)-(c), $B_x(r)$ and $B_y(r)$ have similar spatial structures until high fields to those of $B = 2$ and $\theta = 90^\circ$ in Figs. 2(d)-(e). However, the amplitudes of $B_x$ and $B_y$ at $\theta = 89^\circ$ are much larger than those at $\theta = 90^\circ$. Thus, $|F_x(1,0)|^2$ in Fig. 7(c) and $|F_y(1,1)|^2$ in Fig. 7(d) increase rapidly with decreasing $\theta$ from $90^\circ$. $|F_x(1,1)|^2$ and $|F_y(1,0)|^2$ are less than $10^{-12}$.

When these form factors are compared with each other, the intensity of the spinflip SANS at $q_{1,1}$ from $F_{y}(1,1)$ is much larger than that of the non-spinflip SANS intensity of $|F_{x}(1,1)|^2$ and $|F_{z}(1,0)|^2$. On the other hand, very small intensity of the spinflip SANS at $q_{1,0}$ from $|F_{z}(1,0)|^2$ is difficult to be observed. These correspond to the SANS experimental results on Sr$_2$RuO$_4$, where the spinflip SANS spot was observed only at $q_{1,1} \parallel \vec{B}$. Within the experimental resolution, the spin-flip SANS spot at $q_{1,0}$ and the non-spinflip SANS spots have not been observed yet. We note that similar behaviors of the transverse fields appear also in the nonchiral state including $s$-wave pairing, if $\theta \neq 90^\circ$. Thus, for $\theta \neq 90^\circ$, it is not easy that unique effects due to the chiral state are extracted from qualitative behaviors of the transverse fields.

![FIG. 7: (Color online) $B$-dependence of the FLL form factors in the vortex states, when the magnetic field is slightly tilted from the $ab$ plane, i.e., $\theta = 89.9^\circ$, 89.5$^\circ$, and 89.0$^\circ$. (a) $|F_x(1,0)|^2$. (b) $|F_z(1,1)|^2$. (c) $|F_y(1,0)|^2$. (d) $|F_y(1,1)|^2$. In (a)-(d), the vertical axis is log-scale.](image)

\[ \sum_{\omega_n > 0} \langle \langle \text{Re} \left[ \frac{\Delta^\dagger + \Delta^*}{2(g+1)} \right] \rangle \rangle_p \tau, \]  

which is derived by Doria-Gubernatis-Rainer scaling. The average over $\tau$ indicates a spatial average. In the magnetization curve in Fig. 4(a), we see a change of the slope at $B^*$ for exactly $\vec{B} \parallel ab$ ($\theta = 90^\circ$). This is clearly seen as a step at $B^*$ in the plot of the derivative $dM/dB$ in Fig. 4(b). However this behavior of second-order phase transition is smeared by tilting the field orientation within $1^\circ$. The step in $dM/dB$ was suggested in the experimental observation of the magnetization curve. However in other experimental methods such as specific heat and thermal conductivity, the chiral-nonchiral transition has not been observed yet. This may be because the experimental situation of exactly $\vec{B} \parallel ab$ is difficult to be realized. Our study shows that the chiral-nonchiral transition vanishes by tilting the field orientation within $1^\circ$.

**V. SUMMARY**

We studied the vortex states for nearly $\vec{B} \parallel ab$ in chiral $p$-wave superconductors on the basis of Eilenberger theory. The chiral-nonchiral transition at exactly $\vec{B} \parallel ab$ vanishes by tilting the magnetic field within $1^\circ$. We quantitatively estimated the FLL form factors including transverse fields, and showed that the spin-flip SANS intensity by the transverse fields has large intensity at $(1,1)$-spot. The transverse fields appear even when exactly $\vec{B} \parallel ab$, as a unique effect of the chiral states. These theoretical results indicate the importance of careful studies about the vortex states for nearly $\vec{B} \parallel ab$, to detect some natures of chiral $p$-wave superconductors or Sr$_2$RuO$_4$. 
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