Orbital magnetic dynamics in chiral $p$-wave superconductors

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We present a theory of orbital magnetic dynamics for a chiral $p$-wave superconductor with broken time-reversal symmetry. In contrast to the common Landau-Lifshitz theory for spin ferromagnets, the case of orbital magnetism cannot be described in terms of local magnetization density. Hence it is impossible to define unambiguously the spontaneous magnetic moment: the latter would depend on conditions of its experimental investigation. As an example of this we consider orbital magnetization waves and the domain structure energy.

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Superconductors with unconventional pairing mechanisms have attracted a lot of attention during the past decade. A prominent feature characteristic of unconventional superconductivity is the possibility of states with broken time-reversal symmetry (TRS), which is expressed in the presence of magnetic structure in these materials. Broken TRS has been detected in several superconductors, including Sr$_2$RuO$_4$, ZrZn$_2$, and UGe$_2$, which brought renewed interest to the old-standing problem of coexistence of superconductivity and magnetism.

The layered perovskite Sr$_2$RuO$_4$ is one of the best known examples of unconventional superconductors with broken TRS. Experimental observations of a temperature-independent Knight shift for $H \perp \hat{z}$ (Ref. 3) and an increased muon spin relaxation below $T_c$ (Ref. 1) support the theoretically proposed spin-triplet $p$-wave order parameter

\[ \Delta_{\alpha\beta}(k) = i d^\dagger(k)(\sigma^\alpha\sigma^\beta)_{\alpha\beta}, \]

with $d^\dagger(k) = D_0\hat{z}(k_x + ik_y)$. Such an order parameter has a non-zero orbital moment

\[ l = -i(d^\dagger(k)|k| \nabla_k|d^\dagger(k)|/|d^\dagger(k)|d^\dagger_j|d^\dagger_j), \]

which should lead, in principle, to spontaneous magnetization like in usual ferromagnets. However, coexisting superconductivity screens out this magnetization, making its experimental detection rather difficult. As a possible way to overcome this difficulty, it was proposed to perform microwave response measurements, where excitation of spin waves would provide an immediate signature for the presence of magnetic order.

Previous works which considered magnetic dynamics in superconductors with coexisting magnetism (SCFM), assumed a phenomenological model in which the magnetization was independent from the superconducting order parameter, and its dynamics was described by the Landau-Lifshitz equation (LL dynamics). This model, however, is not so obvious for materials such as Sr$_2$RuO$_4$, where TRS is broken at the superconducting transition, and the magnetic properties are expected to stem from the orbital part of the multicomponent superconducting order parameter. Therefore a proper treatment of magnetic dynamics in such materials should derive it from the dynamics of the superconducting order parameter. An analogous derivation has been done for the A phase of superfluid $^3$He ($^3$He-A); then a modification of the theory taking into account the charge of Cooper pairs would yield magnetic orbital dynamics for an isotropic $p$-wave superconducting electron liquid. However, such a derivation cannot be applied directly to the case of a superconducting metal because of the crystal-field anisotropy.

In this paper, we address this problem and derive an effective orbital magnetic dynamics for a $p$-wave superconductor in a strong crystal-field potential. The resulting dynamics is not equivalent to the phenomenological LL SCFM magnetization dynamics, except for some simple cases. Moreover, in general, it cannot be described in terms of local magnetization at all! Instead, we show that it can be described with the help of a unit vector $l$, corresponding to the angular momentum of Cooper pairs. A similar situation happens in $^3$He A phase, where the presence of an anomalous term in the current does not allow us to express the dynamical equations in terms of local angular momentum density.

We start by considering a Ginzburg-Landau (GL) free energy functional for a $p$-wave superconductor in a strong crystal field possessing a tetragonal symmetry $D_{4h}$ (which is the symmetry of Sr$_2$RuO$_4$). For simplicity, we will assume that the spin does not participate in the dynamics. This assumption is justified, for example, in the case when the direction of $d(k)$ is fixed along the crystal axis $\hat{z}$ by strong spin-orbit coupling. The order parameter can be decomposed into five odd-parity irreducible representations of $D_{4h}$ symmetry. However, limiting the order parameter to a $p$-wave form and requiring $d \parallel \hat{z}$ leaves only two of them: a two-dimensional $\Gamma_5^r = \{\hat{z}k_x, \hat{z}k_y\}$, and a one-dimensional $\Gamma_1^r = \{\hat{k}k_z\}$. Having in mind the order parameter $\sim k_x + ik_y$ for Sr$_2$RuO$_4$, we assume that the $\Gamma_5^r$ representation is most favorable, having the highest transition temperature. It will be described by the GL order parameter $\vec{n} = (\eta_x, \eta_y)$. However, this representation alone is not enough to describe the magnetic properties, since it does not allow any magnetic moment perpendicular to $\hat{z}$, and hence we will need an admixture
of the higher-energy representation $\Gamma^-_1$. The free energy then has the form
\[ f = f_5 + f_1 + f_{1-5} + b^2/8\pi, \] (2)

where $b = \nabla \times A$ is the magnetic field. The form of the GL functional corresponding to representations $\Gamma^-_5$ and $\Gamma^-_1$ is well-known:
\[ f_{5\text{hom}} = P_1(T)|\bar{\eta}|^2 + \beta_1|\eta|^4 + \beta_2(\eta_x^* \eta_y - \eta_y^* \eta_x)^2 \\
+ \beta_3|\eta_z|^2|\eta_y|^2, \] (3)

for the homogeneous part of the functional corresponding to the $\Gamma^-_5$ representation,
\[ f_{5\text{grad}} = K_1[D_{x,y} \cdot \bar{\eta}]^2 + K_2[D_x \eta_y]^2 + |D_y \eta_x|^2 \\
+ K_3(D_x \eta_z)^2(D_y \eta_y) + \text{c.c.} + K_4(D_x \eta_y)^2D_y \eta_z + \text{c.c.} \\
+ K_5[D_x \eta_z]^2 + D_z \eta_y|^2, \] (4)

where $D_i \equiv \delta_i - i(e^*/c)\mathbf{A}_i$ is the gradient mixing term. For the $\Gamma^-_1$ functional, and
\[ f_1 = P_2(T)|\psi|^2 + K_6[D_x \psi]^2 + |D_y \psi|^2 \\
+ K_7[D_z \psi]^2 \] (5)

for the functional corresponding to the $\Gamma^-_1$ representation. For $P_1(T) < 0$ and $P_2(T) > 0$, the ground state is determined by $f_{5\text{hom}}$, which gives $\bar{\eta} = \eta(1, \pm i)$, provided that $\beta_2 > 0$, $4\beta_2 > \beta_3$, and $4(\beta_1 - \beta_2) + 3\beta_3 > 0$. This reproduces the order parameter for Sr$_2$RuO$_4$, and it is also analogous to the Anderson-Brinkman-Morel (ABM) state in superfluid $^3$He with the 1 vector parallel to the $\zeta$ axis. Assuming the couplings of $f_{5\text{hom}}$ to be dominant (consistent with the strong anisotropy and the tendency to TRS breaking), the order parameter $\bar{\eta}$ is frozen in the ground state
\[ \bar{\eta} = \eta(1, i) \quad \text{and} \quad |\eta|^2 \equiv \eta_0^2 = -P_1(T)/(4(\beta_1 - \beta_2) + 3\beta_3), \] (6)

while for the second type four terms are possible:
\[ |\bar{\eta}|^2[D_{x,y}^* D_x \psi D_y \psi + \text{c.c.} \\
+ (\eta_x^* \eta_y - \eta_y^* \eta_x)(D_x \psi)^* D_y \psi - (D_y \psi)^* D_x \psi) \\
+ (\eta_z^2 - \eta_y^2)(D_x \psi)^2 - (D_y \psi)^2) + \text{c.c.} \] (9)

Collecting all terms and using Eq. (6), we obtain for the free energy:
\[ f = (K_1 + K_2)|D_{x,y}\eta|^2 + (K_3 - K_4)i(|D_x \psi|D_y \eta - \text{c.c.}) \\
+ 2K_3|D_x \eta|^2 + Q_1(|D_x \psi|^2(D_x + iD_y) \eta + \text{c.c.}) + Q_2(|D_y \psi|^2(D_y + iD_x) \psi + \text{c.c.}) + Q_3|D_x \psi|^2|D_y \psi|^2 + \text{c.c.} \\
+ Q_4(|\eta|^2)^2(D_x \psi)^2 - (D_y \psi)^2) + \text{c.c.} + f_1 + P_3|\bar{\eta}|^2 + b^2/8\pi. \] (10)

As we have already mentioned, the state $\psi = 0, \bar{\eta} = \eta(1, i) \equiv \eta(1, i, -\delta \alpha \bar{e}^i \varphi)$, where $\delta \alpha$ is a small rotation angle, and $\varphi$ is the angle between the axis of rotation $\mathbf{1}$ and the $\zeta$ axis. From this we identify $\psi/\eta = -\delta \alpha \exp[i\varphi]$. On the other hand, $l_z \delta \alpha \exp[i\varphi] = (l_z + il_y)$, so we make a replacement $\psi = -\eta(l_z + il_y)$. Thus the presence of a small admixture of the order parameter $\psi$ is analogous to tilting the vector $\mathbf{1}$ away from the $\zeta$ direction. We stress that this is only correct as long as $\psi$ is small. Large deflections of $\mathbf{1}$ require a breaking of the state given by Eq. (10), and then the configuration space $(\bar{\eta}, \psi)$ might include also states in which $\mathbf{1}$ is not defined at all (TRS-conserving states). This is a consequence of strong crystal-field anisotropy (for a weak anisotropy, like in $^3$He, this would not be a problem, since there states with defined $\mathbf{1}$ - namely, the ABM states are well separated by the energy from other states).

In addition to the tilting of the vector $\mathbf{1}$, excitation of the orbital mode results also in superconducting flow with the velocity $\mathbf{v} = \nabla \phi/m^* - (e^*/m^*c)\mathbf{A}$. We would like to mention in passing that since only small deviations of $\mathbf{1}$ from $\zeta$ are considered, no problem of nonzero curl of the phase gradient (Mermin-Ho relation) arises here. Expressing Eq. (10) in terms of the new variables $l_\perp = (l_x, l_y)$ and $\mathbf{v}$ (in the harmonic approximation), we obtain for the free energy
\[ f = \frac{1}{2}|D_{x,y}|v_{xy}|^2 + \frac{1}{2}l_z|v_{z}|^2 + B_1|\nabla \cdot l_\perp|^2 \\
+ B_2|\nabla \times l_\perp|^2 + B_3|\partial_t l_\perp|^2 + C \mathbf{v} \cdot \nabla \times l_\perp \\
- C \nabla \cdot l_\perp + Z \partial_t l_\perp + \frac{\alpha}{2}|l_\perp|^2 + b^2/8\pi, \] (12)
where all the coefficients may be expressed via those of
Eq. 10: for example, \(\alpha = 2[P_2(T)\eta_0^2 + P_3\eta^4]\), \(C = Q_1\eta_0^2 m^*\), and \(C_{an} = (Q_1 + Q_2)\eta_0^2 m^*\). Without anisotropy
terms \(\propto \alpha \) and \(\propto Z\), this functional is analogous to the
energy of the \(^3\)He-A phase, \[12\], linearized with respect
to small deviations from the \(\hat{z}\) axis. The free energy de-
termines the (charge) current density:
\[
j = \frac{e^*}{m^*} \frac{\partial f}{\partial \nu} = j_r + j_m ,
\]
where
\[
j_r = \frac{e^*}{m^*} (\rho_{xy} v_{xy} + \rho_z v_z)
\]
is the transport current, and
\[
j_m = \frac{e^*}{m^*} (C \nabla \times \mathbf{l}_\perp - C_{an} \nabla \times \mathbf{m})
\]
is the magnetization current [here \(\nabla \times \mathbf{m} = (\partial_x, \partial_y, 0)\)].

Comparing with the LL SCFM model, \[9, 10\] one can see that it differs from the latter not only by anisotropy in
the London penetration depth \(\sim 1/\rho_{l0}\) and the in-plane
term \(\sim Z\), but, most importantly, by the presence of an
anomalous term \(\sim C_{an}\) in the magnetization current. This
term renders it impossible to have a consistent de-
nition of magnetization for the system (despite the pres-
ence of magnetization currents!), which makes a crucial
difference between our \(p\)-wave superconductor and LL
SCFM’s. Indeed, this current cannot be expressed as a
curl of magnetization vector \(c \nabla \times \mathbf{m}\), and therefore the
energy of interaction \((m^*/e^*) \mathbf{v} \cdot \mathbf{j}_m\) cannot be cast in the
Zeeman form \(-\mathbf{b} \cdot \mathbf{m}\). We emphasize here that the anoma-
lus term appears not due to the anisotropy, but, rather,
due to the orbital nature of magnetism in our system.

The magnetic dynamics is generally described by Ampere’s
law
\[
\frac{4\pi}{c} (j_r + j_m) = \nabla \times \mathbf{b} ,
\]
by the London equation
\[
\nabla \times \mathbf{v} = -\mathbf{\xi} \mathbf{b} ,
\]
where \(\mathbf{\xi} \equiv (e^*/m^*c)\), and by an equation of motion for \(I\).
This motion is described by the equation
\[
\frac{\partial \mathbf{l}_\perp}{\partial t} = -g \hat{z} \times \frac{\partial f}{\partial \mathbf{l}_\perp} = -g \hat{z} \times \left( \frac{\partial f}{\partial \mathbf{l}_\perp} - \partial_i \frac{\partial f}{\partial \partial_i \mathbf{l}_\perp} \right) ,
\]
which is a precession equation for a unit axial vector,
with a generalized torque given by the expression in the
brackets. This equation can be obtained from thermody-
namic conservation requirements by a general procedure
used for the derivation of hydrodynamic equations. \[20\]
Its generalization for the orbital dynamics of \(^3\)He-A has
been considered in Ref. \[21\] [Eq. \[18\]] is obtained by assum-
ing small deviations from the equilibrium and zero
normal-fluid velocity. Although this equation formally
looks the same as the LL equation for spin ferromagnets,
there is an essential difference: the dynamical constant \(g\)
here is an unknown phenomenological parameter, while in
the LL equation it is given by \(g = \gamma/M_0\), where \(\gamma\)
is the gyromagnetic ratio, and \(M_0\) is the magnetization.
The value of \(g\) was disputed for \(^3\)He-A, \[14, 17\] but its
determination is a prerogative of the microscopic theory.

Explicitly the equation of motion for the \(I\) vector reads
\[
\dot{\mathbf{l}}_\perp = -g \hat{z} \times \left[ \mathbf{\alpha} \mathbf{1}_\perp \right. \left. - 2B_1 \nabla (\nabla \cdot \mathbf{1}_\perp) - 2B_2 \hat{z} \times \nabla (\nabla \times \mathbf{1}_\perp) \right.
\]
\[
- 2B_3 \partial_i^2 \mathbf{1}_\perp - C \mathbf{\xi} \mathbf{b} - C_{an} \nabla \times \mathbf{v}_z - Z \partial_i \partial_i \mathbf{1}_\perp (1 - \delta_{ij}).
\]
Equations \[13-17\] and \[19\] constitute a full system of
equations describing the magnetic dynamics. As was ex-
plained above, the magnetic order parameter \(I\) cannot be
identified as a local magnetization, despite being analo-
gous to it, and hence the dynamics is, in general, more
complicated than the LL SCFM dynamics. In special
cases, however, namely in situations where all variations
are either parallel or perpendicular to the \(\hat{z}\) axis, the
dynamics can be described in terms of local magnetiza-
tion, so an effective LL SCFM description is valid (with
a generalization of using tensorial stiffness parameters
and the London penetration depth, as well as in-plane
anisotropy). Still, the value of the effective magnetiza-
tion determined this way would not be unique, but rather
different in each case. For example, for magnetization
waves, propagating in the \(\hat{z}\) direction (perpendicular
geometry), the free energy is
\[
f = \frac{\rho_{xy}}{2} |v_{x,y}|^2 - C \mathbf{\xi} \mathbf{b} \cdot \mathbf{l}_\perp + B_1 |\partial_z \mathbf{l}_\perp|^2 + \frac{\alpha}{2} |\mathbf{l}_\perp|^2 + \frac{b^2}{8\pi} \frac{C^2 \xi^2 q^2}{q^2 + 4\pi \rho_{xy} \xi^2}. \]
leading for plane waves \(\propto e^{iQ \cdot r - i\omega t}\) to a dispersion
\[
\frac{\omega}{g} = \pm \left( \alpha + 2B_3 q^2 - 4\pi C^2 \xi^2 q^2 \right)^{1/2}
\]
This has a form of a LL SCFM spectrum that was con-
sidered in Ref. \[10\] with an equilibrium magnetization
\(M_0 = \xi C\). Hence the results for the microwave response
of a LL SCFM given in that paper are directly applicable
to the case of an unconventional superconductor consid-
ered here, provided that an effective equilibrium magne-
tization is chosen as \(M_0 = \xi C\). On the other hand, for
waves propagating in the \(\hat{\mathbf{x}} - \hat{\mathbf{y}}\) plane, the spectrum is
\[
\frac{\omega^2}{g^2} = - \left( \frac{Q_3 \sin 3\phi}{4} \right)^2 + \left( \alpha + 2B_1 + \frac{Z \sin^2 2\phi}{8\pi} q^2 \right)^2
\]
\[
\times \left( \alpha + 2B_2 - \frac{Z \sin^2 2\phi}{8\pi} q^2 \right)^2 - 4\pi \left( C - C_{an}\right)^2 \xi^2 q^2 \]
\[
\left( q^2 + 4\pi \rho_{xy} \xi^2 \right)^2 \]
where \(\phi\) is the angle the wave vector \(q\) makes with the
\(\hat{\mathbf{z}}\) axis. This expression corresponds to the LL SCFM
spectrum in the parallel geometry with a different
value of equilibrium magnetization \(M_0 = \xi (C - C_{an})\) (in
that analysis no $\hat{x} - \hat{y}$ plane anisotropy was introduced, so the $\phi$ dependence was irrelevant there). We stress again, however, that these results cannot be interpreted in terms of a tensorial $M_0$. For a general propagation direction the dynamics cannot be interpreted in terms of local magnetization at all.

Another issue in which important differences arise between our unconventional superconductor and a LL SCFM, is the issue of domain walls. A basic feature of ferromagnets, both insulating and superconducting, which determines the field distribution inside domains, is the discontinuity of the magnetic induction $\Delta b_z$ across the wall being given by $8\pi M_0$, with $M_0$ the equilibrium magnetization inside the domains. Moreover, this discontinuity may be used as an experimental definition of domain magnetization. One might ask, whether in our case $\Delta b_z/8\pi$ has any meaning of magnetization and whether it is related in any sense to dynamic response properties discussed above. The answer to the second question is, in general, negative. Indeed, in order to determine $\Delta b_z$, one has to examine what happens inside the domain wall, which involves a strong deviation from the ground state, Eq. (2). For that purpose, the free energy Eq. (12) is inapplicable, and one has to go back to the general expression, Eq. (2). With the energy $f_{5\, \text{hom}}$ being dominant, a domain-wall solution may be restricted to the $(\eta_x, \eta_y)$ part of the order parameter. For specific parameter values, domain-wall solutions have been obtained, with the discontinuity being determined by the parameters of $f_{5\, \text{grad}}$. For example, when $\beta_1, \beta_3 \gg \beta_2$, and $K_3 = K_4$, the discontinuity is

$$\Delta b_z = \frac{4\pi}{c} \frac{P_1(T)\epsilon^*}{\beta_3} \frac{K_3 - K_2}{K_1 + K_2}. \quad (23)$$

Thus, $\Delta b_z$ has nothing to do either with the parameters $C$ and $C_{an}$, which determine the magnetization current for small deviations, or the dynamical parameter $g$. Only in the limit of very weak crystal anisotropy, when the vector $l$ is well defined even inside the domain wall as shown in Fig. 1 the discontinuity $\Delta b_z$ is determined by a combination of $C$ and $C_{an}$. Indeed, in this case the magnetization current inside a domain wall is given by the isotropic form

$$j_m = \frac{\epsilon^*}{m^*} [C \nabla \times l - C_{an}(l \cdot \nabla \times l)], \quad (24)$$

and then the field discontinuity for a Bloch wall (where $l$ rotates in the wall plane) is given by $\Delta b_z = 8\pi \xi (C - C_{an})$, corresponding to an effective magnetization $\xi (C - C_{an})$. On the other hand, for a Néel wall (where $l$ rotates in the plane perpendicular to the wall) the discontinuity becomes $\Delta b_z = 8\pi \xi C$, corresponding to an effective magnetization $\xi C$.

As to the first question, $\Delta b_z/8\pi$ does have the correspondence to magnetization in the usual ferromagnets.

To see it, let us recall that for usual ferromagnets the domain-structure energy is given by the sum of a local domain-wall energy (which may be considered an effective parameter independent of the fields) and an electromagnetic energy inside the domains (including the energy of Meissner currents and the Zeeman term $-b_z M_0$). In our case, the Zeeman term is absent, and instead the domain-wall energy contains a contribution due to the domain-wall currents. However, for thin domain walls this contribution can be transformed to an effective Zeeman form:

$$- \int_{\text{wall}} \frac{j_y A_y}{c} \, dx = - \int_{\text{domain}} \tilde{M} b_z(x) \, dx, \quad (25)$$

where $\tilde{M} = \Delta b_z/8\pi$, and the domains are in the $\hat{x}$ direction. Hence $\Delta b_z/8\pi$ may be interpreted as an effective magnetization of each domain, which allows the electromagnetic part of the domain-wall energy to be cast in the usual Zeeman form, so that the remaining part is entirely local (i.e., independent of the field distribution). This static “magnetization” determines the field in the domains, but is unrelated, in general, to the dynamical response of the material. As an application of this result, one can immediately obtain a criterion of stabil-

![FIG. 1: Domain wall between two domains with opposite directions of the $l$ vector in the case of weak anisotropy: a) Bloch wall; b) Néel wall. Note that when the anisotropy is strong $l$ is not defined inside the wall.](image-url)
ity against the formation of domains for our material. For this purpose we can use the result of Krey’s analysis \[24\] for a LL SCFM, that the uniformly magnetized configuration becomes unstable when \(W < 4\pi\lambda M_0^2\), where \(W\) is the local domain-wall energy (calculated excluding electromagnetic effects), and \(\lambda\) the London penetration depth. The same result is valid for our case if \(\tilde{M}\) defined above is substituted for \(M_0\), and the in-plane penetration depth is used, \(\lambda^{-2} \rightarrow 4\pi\xi^2\rho_{xy}\). Of course, \(W\) should be calculated separately, as was done in Refs. \[22\] and \[23\].

In summary, we have studied the orbital magnetic dynamics in a \(p\)-wave superconductor with strong crystal-field anisotropy. The dynamics is essentially different from the Landau-Lifshitz dynamics for a superconducting ferromagnet. The most important difference is that the directional order parameter \(l\) (orbital moment of Cooper pairs) does not lead to a definite spontaneous magnetization (magnetic moment per unit volume). While in simple cases one can introduce an effective magnetic-moment density similar to that in the Landau-Lifshitz dynamics, the value of this density varies from case to case. As examples of these cases we have considered magnetization waves along and normal to the main crystal axis, and the energy of the domain structure.

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