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Magnetoelectric effects in heavy-fermion superconductors without inversion symmetry

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We investigate the effects of strong electron correlation on magnetoelectric transport phenomena in noncentrosymmetric superconductors with particular emphasis on its application to the recently discovered heavy-fermion superconductor CePt$_3$Si. Taking into account electron correlation effects in a formally exact way, we obtain the expression of the magnetoelectric coefficient for the Zeeman-field-induced paramagnetic supercurrent, the existence of which was predicted more than a decade ago. It is found that in contrast to the usual Meissner current, which is much reduced by the mass renormalization factor in the heavy-fermion state, the paramagnetic supercurrent is not affected by the Fermi liquid effect. This result implies that the experimental observation of the magnetoelectric effect is more feasible in heavy-fermion systems than that in conventional metals with moderate effective mass.

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I. INTRODUCTION

It has been discussed for decades that metallic systems with noncentrosymmetric crystal structure may exhibit nontrivial magnetoelectric effects. The existence of an asymmetric potential gradient $\nabla V$ due to the noncentrosymmetric structure gives rise to the spin-orbit interaction $(\vec{p} \times \nabla V) \cdot \vec{a}$, which breaks the individual inversion and spin rotation symmetry. As a result, the charge or energy current operator may couple to the spin-density operator. In this context, current flows induced by an applied magnetic field and current-flow-driven magnetization have been investigated extensively both in normal metals and superconductors. In particular, Edelstein predicted the remarkable magnetoelectric effect in superconducting states; i.e., in noncentrosymmetric superconductors, the Zeeman field induces a supercurrent, and conversely, the supercurrent flow induces a magnetization. Later, the former effect is elegantly reformulated by Yip in terms of the “van Vleck” contribution which stems from the inversion-symmetry-breaking spin-orbit interaction. Since a static magnetic field cannot induce dissipative current flows, the Zeeman-energy-induced current should vanish in the normal state. However, in the superconducting state, the existence of the paramagnetic supercurrent is not forbidden in the absence of the inversion symmetry. The recent discovery of superconducting materials without inversion symmetry such as CePt$_3$Si, UIr, and Cd$_2$Re$_2$O$_7$ stimulates the renewed interest in this issue. Under an applied magnetic field, the Meissner diamagnetic supercurrent in addition to the Zeeman-field-induced paramagnetic supercurrent should exist. Thus, it is important for the experimental observation of this effect to discriminate between these two supercurrents.

In the present paper, we would like to investigate the Fermi liquid corrections to this Edelstein magnetoelectric effect, which may be important for the application to heavy-fermion superconductors such as CePt$_3$Si and UIr. We obtain the formula for the magnetoelectric effect coefficient taking into account Fermi liquid corrections exactly. The most important finding is that the Zeeman-energy-induced paramagnetic supercurrent is not at all affected by electron correlation effects provided that ferromagnetic spin fluctuation is not developed, in contrast with the diamagnetic Meissner current of which the magnitude is much reduced by the mass renormalization effect. This result implies that the experimental detection of the paramagnetic supercurrent in heavy-fermion superconductors may be more feasible than that in weakly correlated metals.

The organization of this paper is as follows. In Sec. II, we present the basic formulation of the Fermi liquid theory for a model system without inversion symmetry. We would like to make a brief comment on the superconducting state realized in CePt$_3$Si in Sec. III. In Sec. IV, the exact formula of the magnetoelectric coefficient is obtained. In Sec. V, the implication for the experimental observation of this effect is discussed. Summary and discussion are given in the last section.

II. MODEL AND ANALYSIS BASED ON THE FERMI LIQUID THEORY

As a simplest model which realizes the broken inversion symmetry, we consider an interacting electron system with the Rashba spin-orbit interaction. The Hamiltonian is given by

$$\mathcal{H} = \sum_{p,\sigma} \varepsilon_p c_p^{\dagger} c_p + \alpha \sum_{p,\sigma,\sigma'} (\vec{p} \times \vec{n}) \cdot \vec{a}_{\sigma\sigma'} c_p^{\dagger} c_{p'p} + U \sum_i n_i n_{i+1}$$

where $c_p^{\dagger}$ ($c_p$) is the creation (annihilation) operator for an electron with spin $\sigma$ and momentum $p$. The number density operator at the site $i$, $n_i = c_i^{\dagger} c_i$. The second term of Eq. (1) is the Rashba spin-orbit interaction which incorporates the broken inversion symmetry. Here, the unit vector parallel to the asymmetric potential gradient is given by $\vec{n} = (0, 0, 1)$. This system is considered to be a model of CePt$_3$Si, with which we are mainly concerned in this paper. The $f$ electron of CePt$_3$Si is in the $\Gamma_7$ Kramers doublet state. Expanding the $\Gamma_7$ doublet in terms of the $s_z=\pm 1/2$ and $-1/2$ basis, we found that the Rashba spin-orbit interaction term expressed in terms of the $\Gamma_7$ basis has the matrix structure given above.
up to a constant factor which can be absorbed into the redefinition of the coupling constant $\alpha$. Thus, in the case of CePt$_3$Si, the spin index $\sigma$ in Eq. (1) represents the $\Gamma_7$ Kramers doublet. According to the LDA band calculations, the electronic structure near the Fermi level in CePt$_3$Si is predominated by the $f$-electron component. Since the leading interaction between $f$ electrons is an on-site Coulomb type, we assume the interaction term as given by Eq. (1) for simplicity.

In the following, we do not specify the pairing mechanism of superconductivity, but assume that the superconducting state is realized by an effective pairing interaction with an angular momentum $l \geq 1$ which may stem from the on-site Coulomb interaction in Eq. (1) or may have any other origin not included in the Hamiltonian (1). Then, we can analyze electron correlation effects on this superconducting state in a formally exact way by using the superconducting Fermi liquid theory.\cite{14,15,16}

In the conventional Nambu representation,\cite{17} the inverse of the single-particle Green’s function is defined as

$$\hat{G}^{-1}(p) = \left( \begin{array}{cc} i\varepsilon_n - \hat{H}(p) & -\hat{\Delta}(p) \\ -\hat{\Delta}^\dagger(p) & i\varepsilon_n + \hat{H}^\dagger(-p) \end{array} \right),$$

where $p=(\hat{p},i\varepsilon)$, and

$$\hat{H}(p) = \hat{H}_0(p) + \hat{\Sigma}(p),$$

$$\hat{H}_0 = \varepsilon_p - \mu + \alpha(\hat{p} \times \hat{n}) \cdot \hat{\sigma} - \mu_0 \tau \hat{H}_s,$$

with $\mu$ the chemical potential. Here, to discuss the magnetoelectric effect, we take into account the Zeeman magnetic field in the $x$ direction, $H_s$. For simplicity, we assume that the $g$ value is equal to 2. The self-energy matrix $\hat{\Sigma}$ consists of both diagonal and off-diagonal components,

$$\hat{\Sigma} = \left( \begin{array}{cc} \Sigma_{\uparrow\uparrow}(p) & \Sigma_{\uparrow\downarrow}(p) \\ \Sigma_{\downarrow\uparrow}(p) & \Sigma_{\downarrow\downarrow}(p) \end{array} \right).$$

We can easily see from the symmetry argument that under the applied in-plane magnetic field, $\Sigma_{\uparrow\uparrow}(p) = \Sigma_{\downarrow\downarrow}(p) = \Sigma(p)$ and $\Sigma_{\uparrow\downarrow}(\hat{p},i\varepsilon) = \Sigma_{\downarrow\uparrow}(\hat{p},-i\varepsilon)$. The superconducting gap function is $\{\Delta(p)\}_{\alpha\beta} = \Delta_{\alpha\beta}(p)$ ($\alpha, \beta = \uparrow, \downarrow$). The self-energy correction terms $i\varepsilon_n - \hat{H}(p)$ and $i\varepsilon_n + \hat{H}^\dagger(-p)$ in $\hat{G}^{-1}(p)$ are diagonalized by the unitary transformation $\hat{A}(p)\hat{G}^{-1}(p)\hat{A}^\dagger(p)$ with

$$\hat{A}(p) = \left( \begin{array}{cc} \hat{U}(\hat{p}) & 0 \\ 0 & \hat{U}^\dagger(-\hat{p}) \end{array} \right),$$

$$\hat{U}(\hat{p}) = \frac{1}{\sqrt{2}} \left( \begin{array}{cc} 1 & i\hat{I}_z \\ i\hat{I}_y & 1 \end{array} \right),$$

where $\hat{I}_z = \hat{I}_z \pm i\hat{I}_y$, and $\hat{I}_x$ and $\hat{I}_y$ are, respectively, the $x$ and $y$ components of the unit vector $\hat{t}_p = \frac{\hat{p}}{\vert \hat{p} \vert}$ with

$$\hat{t}_p(ie) = \left( \frac{p_x + \frac{1}{\alpha} \text{Im} \Sigma_{\uparrow\uparrow}(p) + \frac{1}{\alpha} \text{Re} \Sigma_{\uparrow\uparrow} - \frac{\mu_0 H_s}{\alpha}}{\mu + \alpha(\hat{p} \times \hat{n}) \cdot \hat{\sigma} - \mu_0 \tau \hat{H}_s}, 0 \right).$$

As seen from Eqs. (3)–(5), the main effect of the off-diagonal self-energy $\Sigma_{\uparrow\downarrow}$ is to renormalize the Rashba interaction term, replacing the momentum $\hat{p}$ in the Rashba term with the vector $\hat{t}_p$. Since the on-site Coulomb interaction does not change the symmetry of the system, the off-diagonal self-energy should satisfy the following condition in the absence of magnetic fields:

$$\text{Re} \Sigma_{\uparrow\downarrow}(\hat{p},\mu) = - \text{Re} \Sigma_{\uparrow\downarrow}(\hat{p},\mu),$$

$$\text{Im} \Sigma_{\uparrow\downarrow}(\hat{p},\mu) = - \text{Im} \Sigma_{\uparrow\downarrow}(\hat{p},\mu),$$

and $\text{Re} \Sigma_{\uparrow\downarrow}(\text{Im} \Sigma_{\uparrow\downarrow})$ is an even function of $p_x(p_y)$.

In the normal state, the single-particle excitation energy $\varepsilon_{p\tau}$ for the quasiparticle with the helicity $\tau = \pm 1$ is given by the solution of the equation $z - \hat{H}(\hat{p},z) = 0$, which, in the diagonalized representation,

$$\varepsilon_{p\tau} + \mu - \tau a_p[\varepsilon_{p\tau} - \Re \Sigma(\hat{p},\varepsilon_{p\tau}^*)] = 0.$$

The gap functions $\Delta(p)$ and $\Delta^\dagger(p)$ are diagonalized simultaneously, and the nonzero off-diagonal components of $\Delta(p)$ which correspond to the Cooper pairing between the different Fermi surfaces induce pair-breaking effects, resulting in the decrease of the transition temperature $T_c$. Thus, the highest transition temperature is achieved by the gap function given by Eq. (12).\cite{19} The realization of the gap function (12) in the case with no inversion center is also elucidated by the group theoretical argument.\cite{20}

Taking the inverse of Eq. (2), we have

$$\hat{G}(p) = \left( \begin{array}{cc} \hat{G}(p) & \hat{F}(p) \\ \hat{F}^\dagger(p) & -\hat{G}^\dagger(-p) \end{array} \right),$$

where

$$\hat{G}(p) = \sum_{\tau=\pm 1} \frac{1 + \tau \hat{I}_z \cdot \hat{\sigma}}{2} \hat{G}_\tau(p),$$

$$\hat{F}(p) = \sum_{\tau=\pm 1} \frac{1 + \tau \hat{I}_z \cdot \hat{\sigma}}{2} i\sigma_y F_\tau(p),$$

and

$$\hat{G}_\tau(p) = \frac{\varepsilon_{p\tau}(ie + e_{p\tau}^*)}{(ie + i\gamma \text{sgn} \varepsilon)^2 - t_{p\tau}^2},$$

024515-2
The symmetries of \( \Gamma_s(p,p') \) and \( \Gamma_f(p,p') \) in the momentum space are characterized by the same irreducible representation, and as a result, \( \Delta_f(p) \) and \( \Delta_f(p) \) in Eq. (12) have the same symmetry in the momentum space. The realized superconductor state is the mixture of the spin singlet and triplet states. In this case, the possible pairing state is \( s+p \) or \( d+f \) or \( g+h \), and so forth.

III. A COMMENT ON THE SUPERCONDUCTING STATE REALIZED IN CePt_3Si: THE POSSIBILITY OF AN UNUSUAL COHERENCE EFFECT

Here, we would like to make a brief remark about the pairing state realized in CePt_3Si. The NMR measurement carried out by Yogi et al. shows the existence of the coherence peak of \( 1/T_1T \) just below \( T_c \), indicating the full-gap state without nodes. On the other hand, the recent experiment on the thermal transport done by Izawa et al. supports the existence of line nodes of the superconducting gap. A possible resolution of this discrepancy is that the line nodes of the superconducting gap are generated accidentally at the magnetic zone boundary which emerges as a result of the antiferromagnetic phase transition at \( T_N = 2.2 \) K, and crosses the Fermi surface. For such accidental nodes without the sign change of the superconducting gap function, the coherence factor of \( 1/T_1T \) does not vanish, resulting in the enhancement of the coherence peak just below \( T_c \). In this case, a plausible candidate for the pairing state is the \( s+p \) wave state. An important point which we would like to stress here is that even when the superconducting state is dominated by the \( p \) wave pairing, i.e., \( \Delta_f(p) \ll \Delta_s(p) \), the coherence factor which enters into \( 1/T_1T \) does not vanish. This contrasts with the case of the usual \( p \) wave state realized in centrosymmetric superconductors, where the coherence factor of \( 1/T_1T \) disappears. This is understood as follows. For simplicity, we ignore electron correlation effects. Then, in noncentrosymmetric superconductors, the nuclear relaxation rate is given by

\[
\frac{1}{T_1T} \approx \lim_{\omega \rightarrow 0} \frac{1}{\omega} \text{Im} \left( T \sum_{\epsilon_m p-p'} \left\{ \frac{1}{2} \left[ \frac{\Delta_f(p,\epsilon_m)}{\Delta_s(p,\epsilon_m)} - \Delta_s(p,\epsilon_m) - \Delta_f(p,\epsilon_m) \right] \right\} \right) \frac{1}{2T} \frac{1}{\cosh^2 \frac{\epsilon}{2T}} \left[ \frac{1}{2} \left( N_s(e)^2 + N_s(e)^2 \right) \right],
\]

where \( N_s(e) \) and \( N_s(e) \) defined by the retarded Green’s functions as

\[
N_s(e) = - \sum_p \sum_{\tau=\pm} \text{Im} \ G_s^\tau(p,e),
\]

\[
N_s(e) = - \sum_p \sum_{\tau=\pm} \text{Im} \ F_s^\tau(p,e).
\]

The expression of \( 1/T_1T \) (21) does not rely on the phase factor \( \vec{\tau}_p \cdot \vec{n} \) of the triplet component of the gap function (12). The second term on the right-hand side of Eq. (21) gives the nonzero contribution from the coherence factor, as in the case of conventional \( s \)-wave superconductors. This property enhances the coherence peak of \( 1/T_1T \) prominently. It may be important to take into account this unusual coherence effect for clarification of the origin of the notable coherence peak of \( 1/T_1T \) in CePt_3Si. It is also intriguing to explore the unusual coherence effect on other response functions, such as the ultrasonic attenuation. We would like to address this issue elsewhere.
IV. MANY-BODY EFFECTS ON THE MAGNETOELECTRIC TRANSPORT IN THE SUPERCONDUCTING STATE

In this section, we consider electron correlation effects on the magnetoelectric transport in the superconducting state first found by Edelstein and later discussed in detail by Yip i.e., the emergence of the paramagnetic supercurrent induced by the Zeeman magnetic field in the direction normal to the \( \hat{\mathbf{n}} \) vector.\(^3\)

We consider the charge current flowing in the \( y \) direction, which is defined as

\[
J_y = T \sum_{\nu,\rho} \frac{1}{2} \text{Tr}[\hat{V}_{\nu\rho}(\rho)G(p)]
\]

with

\[
\hat{V}_{\nu\rho} = \begin{pmatrix} \hat{v}_{\nu\rho}(\rho) & 0 \\ 0 & -\hat{v}_{\nu\rho}(\rho) \end{pmatrix}
\]

and \( \hat{v}_{\nu\rho} = -\partial_{\nu\rho} - e\vec{\sigma} \).

The transport coefficient which characterizes Edelstein’s magnetoelectric effect is given by

\[
\mathcal{K}_{yx} = \left. \frac{\partial J_y}{\partial H_x} \right|_{H_y=0} = - T \sum_{\nu,\rho} \frac{1}{2} \text{Tr} \left[ \hat{V}_{\nu\rho}(\rho) \frac{\partial G^{-1}(p)}{\partial H_x} G(p) \right] \bigg|_{H_y=0}.
\]

Following Ref. 3, to simplify the expression of \( \mathcal{K}_{yx} \) we use Ward’s identity for the current vertex,

\[
\frac{\partial \hat{G}(p)}{\partial p_y} = \hat{G}(p) \hat{V}_y(p) \hat{G}(p) + \hat{F}(p) \hat{V}_y(-p) \hat{F}(p, -ie) + \hat{R}(p),
\]

where \( \hat{V}_y(p) = \hat{v}_{\nu\rho} + \frac{\partial \hat{\Sigma}(p)}{\partial p_y} \), and

\[
\hat{R}(p) = \begin{pmatrix} 0 & r(p) \\ r^*(p, -ie) & 0 \end{pmatrix},
\]

\[
r(p) = \frac{G_+ - G_-}{2 \alpha |\tilde{p}|} - G_+ E_F - \frac{1}{2} \alpha |\tilde{p}| \Lambda_{+uu}^v(p).
\]

\[
\Lambda_{+uu}^v(p) = \hat{t}_x \left( \alpha + \frac{\partial \text{Re} \Sigma_{11}}{\partial p_y} \right) - \hat{t}_y \frac{\partial \text{Im} \Sigma_{11}}{\partial p_y}.
\]

Then, from Eqs. (13), (26), and (27), we obtain

\[
\mathcal{K}_{yx} = - 2T \sum_{\nu,\rho} \text{Tr} \left[ \hat{F}(p) \hat{v}_{\nu\rho}(-p) \hat{F}(p) \left( \mu_B \sigma_\nu - \frac{\partial \hat{\Sigma}}{\partial H_x} \right) \right]
\]

\[
- 2 \mu_B T \sum_{\nu,\rho} \text{Tr} \left[ \frac{\partial \hat{\Sigma}}{\partial H_x} \right]
\]

\[
- T \sum_{\nu,\rho} \text{Tr} \left[ \frac{\partial \hat{G}}{\partial H_x} \right],
\]

where the three-point vertex function is

\[
\Lambda_{+uu}^v(p) = \hat{t}_x \left( 1 - \frac{1}{\mu_B} \frac{\partial \text{Re} \Sigma_{11}}{\partial H_x} + \hat{t}_y \frac{\partial \text{Im} \Sigma_{11}}{\partial H_x} \right).
\]

The last two terms of Eq. (31), which emerge as a result of Fermi liquid corrections, can be rewritten by using Luttinger-Ward’s identity generalized to the superconducting state. Luttinger-Ward’s identity in the normal state reads\(^{21,22}\)

\[
T \sum_{\nu,\rho} \text{Tr} \left[ \frac{\partial \hat{G}}{\partial \rho} \right] = 0.
\]

This relation is obtained by differentiating all closed linked diagrams with respect to \( \rho \).\(^{21,22}\) In the superconducting state, a similar analysis leads to

\[
T \sum_{\nu,\rho} \text{Tr} \left[ \frac{\partial \hat{G}}{\partial \rho} \right] + T \sum_{\nu,\rho} \text{Tr} \left[ \frac{\partial \hat{F}}{\partial \rho} \right] + T \sum_{\nu,\rho} \text{Tr} \left[ \frac{\partial \hat{\Sigma}}{\partial \rho} \right] = 0.
\]

Differentiating Eq. (34) with respect to \( H_x \), and integrating over \( \rho \) by parts, we have

\[
T \sum_{\nu,\rho} \text{Tr} \left[ \frac{\partial \hat{G}}{\partial H_x} \right] - T \sum_{\nu,\rho} \text{Tr} \left[ \frac{\partial \hat{F}}{\partial H_x} \right]
\]

\[
= - T \sum_{\nu,\rho} \text{Tr} \left[ \frac{\partial \hat{\Sigma}}{\partial H_x} \right] + T \sum_{\nu,\rho} \text{Tr} \left[ \frac{\partial \hat{\Sigma}}{\partial H_x} \right]
\]

\[
- T \sum_{\nu,\rho} \text{Tr} \left[ \frac{\partial \hat{\Sigma}}{\partial H_x} \right] + T \sum_{\nu,\rho} \text{Tr} \left[ \frac{\partial \hat{\Sigma}}{\partial H_x} \right].
\]

We see from Eqs. (37) and (35) that \( \mathcal{K}_{yx} \) vanishes exactly in the normal state, as is consistent with the thermodynamic argument that a static magnetic field cannot induce nonequilibrium current flows.\(^3\)\(^5\) The right-hand side of Eq. (35) consists of the terms which have the form \( \Sigma_{n,p} G_x F_x \). The ratio of the contributions from these terms to those from other terms of Eq. (31) is of order \( \Delta/E_F \), and thus we can neglect the last two terms of Eq. (31) approximately. Then, the magnetoelectric coefficient is expressed as
where

\[ \Lambda^{\text{fi}}(p) = i \left( 1 - \frac{1}{\mu_B} \frac{\partial \text{Re} \Sigma_{\downarrow\downarrow}}{\partial H_x} \right) - \frac{i}{\mu_B} \frac{\partial \text{Im} \Sigma_{\uparrow\downarrow}}{\partial H_x} - \frac{i}{\mu_B} \frac{\partial \text{Re} \Sigma}{\partial H_x}. \]  

(37)

It is noted that the vertex corrections due to electron correlation, \( \Lambda^{\text{fi}}(p) \), which appear in the above expression (36) is nothing but the vertex corrections to the uniform spin susceptibility,

\[ \chi_{xx} = \frac{\mu_B^2}{2} \sum_p \frac{\bar{E}_{p\uparrow} \bar{E}_{p\downarrow}}{2E_{p\uparrow}} - \frac{\bar{E}_{p\uparrow} \bar{E}_{p\downarrow}}{2E_{p\downarrow}} \frac{i}{\alpha(p)} \Lambda^{\text{fi}}_x(p). \]  

(38)

Equation (38) is easily obtained by differentiating the \( x \) component of the total magnetization \( S^i = \mu_B \Sigma_{n,p} \text{Tr} \sigma_i \tilde{G}(p) \) with respect to \( H_x \). Note that the first term on the right-hand side of Eq. (38) is the Pauli paramagnetic contribution and the second one is the “van Vleck” term which arises from excitations between spin-orbit split two bands. Generally, in heavy-fermion systems, the magnitude of the uniform spin susceptibility is enhanced by the vertex corrections \( \Lambda^{\text{fi}}(p) \). In typical heavy-fermion systems including CePt_3Si, the Wilson ratio \( R_{\text{W}} = T \chi(C/T_0/C_0) \sim 2 \), which implies that the vertex corrections \( \Lambda^{\text{fi}} \) is approximately of order the mass enhancement factor \( 1/\bar{v}_p^2 \). Therefore, in Eq. (36), effects of the vertex corrections and the mass renormalization factors \( \bar{v}_p \) cancel with each other. This cancellation holds as long as there is no strong ferromagnetic spin fluctuation which increases notably the magnitudes of the vertex corrections \( \Lambda^{\text{fi}} \). Another important feature of Eq. (36) is the absence of the backflow term of the charge current, which usually exists in the nonequilibrium current flow. (See the discussion on the usual Meissner current in the next section.) This is related to the fact that the current induced by a static magnetic field is a dissipationless equilibrium flow. As a result, the Fermi liquid corrections do not exist in this magnetoelectric coefficient for heavy-fermion superconductors, provided that there is no ferromagnetic fluctuation. This is one of the main results of this paper. In terms of the Kubo formula, the absence of electron correlation effects for \( K_{yy} \) is understood as follows. \( K_{yy} \) is given by the correlation function of the current and spin-density operators. The spin-density vertex is renormalized by electron correlation in the opposite way to the current vertex, resulting in the cancellation of the mass renormalization factors. The important implication of this result is that the Zeeman-field-induced paramagnetic supercurrent is not suppressed by electron correlation effects in contrast with the usual diamagnetic supercurrent of which the magnitude is much reduced by the large mass enhancement in heavy fermion systems. This property may make the experimental observation of the magnetoelectric effect easier, as discussed in the next section.

V. IMPLICATIONS FOR EXPERIMENTAL OBSERVATIONS

On the basis of the formula (36), we would like to discuss how the Zeeman-field-induced paramagnetic supercurrent is experimentally observed.

When the magnetic field is applied in the \( x \) direction, the London equation is modified to

\[ \tilde{J}_\| = - \frac{e}{4 \pi \lambda^2} \vec{A} + K_{yy}(\vec{n} \times \vec{H}), \]  

(39)

where \( \vec{A} \) is the vector potential. Since the applied magnetic field always induces both the diamagnetic and the paramagnetic supercurrent in the system with the Rashba spin-orbit interaction, it is important for the experimental observation of this effect to discriminate between these two supercurrents. If one measures currents simply attaching leads to the sample and applying an in-plane magnetic field, the Zeeman-field-induced paramagnetic supercurrent cannot flow in the sample because it generates the Joule heat in the normal metal leads. In this case, the magnetoelectric effect is canceled with the nonzero phase gradient of the superconducting order parameter in the equilibrium state.

Here, to highlight the observation of the paramagnetic supercurrent, we consider an experimental setup composed of two superconducting samples joined together as depicted in Fig. 2. The \( \vec{n} \) vectors of the two superconducting samples A and B (depicted by the gray arrows) are aligned in the directions (0,0,1) and (0,0,−1), respectively. The in-plane magnetic field \( \vec{H} \) is applied in the \( x \) direction. The paramagnetic supercurrent circulates in the system as depicted by the black thin arrows.
where the quasiparticle velocity is the lowest order. The result at the semiclassical approximation based upon the Doppler klapp scattering is expected to be strong, measurements for CePt$_3$Si supports the existence of the line and upper critical field. In fact, the recent thermal transport density of states under the applied in-plane magnetic field for served through the V olovik effect on the single-particle density. In contrast to the diamagnetic Meissner effect, the toelectric effect.

Here $\frac{\mathbf{v}^*}{H_{c2}}$ is the renormalized Fermi velocity, $n_s$ is the superfluid density, and $p_F$ is the Fermi momentum. It should be notified that $n_s$ is renormalized by the backflow effect of $J_{pr}$, and is not equal to the carrier density even at zero temperature. In particular, for heavy fermion systems in which umklapp scattering is expected to be strong, $n_s$ is smaller than the carrier density. In contrast to the diamagnetic Meissner current, the Zeeman-field-induced paramagnetic supercurrent is not influenced by the many-body effects described above, as discussed in the previous section. This difference of electron correlation effects between the diamagnetic and paramagnetic supercurrents can be utilized to detect the magnetoelectric effect.

If the superconducting gap has a nodal structure as is often realized in some heavy-fermion superconductors, the existence of the paramagnetic supercurrent is indirectly observed through the Volovik effect on the single-particle density of states under the applied in-plane magnetic field for $H_{c1} < H < H_{c2}$. Here $H_{c1}$ and $H_{c2}$ are, respectively, the lower and upper critical field. In fact, the recent thermal transport measurements for CePt$_3$Si supports the existence of the line node in the superconducting state of this system. Applying the semiclassical approximation based upon the Doppler shift effect, and assuming a spherical Fermi surface, we calculate the local density of states from the modified London equation (39). In the calculation of $K_{yx}$, we use the fact that the vertex correction $\Lambda_{\alpha+i\eta}$ is appropriately approximated as $\sim -e^{-\Delta^2}$ in typical heavy-fermion systems as discussed in Sec. IV, and we expand Eq. (36) in terms of $\alpha p_F/E_F$ up to the lowest order. The result at $T=0$ is

$$
\begin{align*}
\delta D_{loc}(0) & \sim \left[ \frac{H}{H_{c2}} + \frac{H e \mu_B m^* \phi_0}{\xi c^*} \right] 10^{-24} \cdot 0.48 \times 10^{-24} \frac{H}{H_{c2}}.
\end{align*}
$$

Here $\Phi_0=\hbar c/(2e)$, and $n_0$ is the density of electrons. $E_F$ is the unrenormalized Fermi energy. The first term on the right-hand side of Eq. (42) is due to the usual Volovik effect, and the second term linearly proportional to the applied magnetic field stems from the Zeeman-energy-induced paramagnetic supercurrent. Thus, the magnetic-field dependence distinguishes between the paramagnetic and diamagnetic currents. The above behavior of the local density of states may be observed by the measurement of the specific-heat coefficient or the thermal conductivity in sufficiently low field regions. The above expression of $\delta D_{loc}(0)$, it is seen that the conventional diamagnetic contribution is suppressed by the mass renormalization factor $z^2_{pr}$ which appears through $v_p^*$, whereas the magnetoelectric contribution is not affected by this correlation effect. It is also noted that the carrier density which enters into the paramagnetic term is not $n_s$ but equal to the electron density $n_0$. This is due to the absence of the backflow term in the Zeeman-energy-induced supercurrent. As mentioned before, $n_s$ is affected by the backflow term. For simplicity, we assume that $n_s = n_0$ for a while.

In the case of CePt$_3$Si, according to the measurement of $H_{c2}$, the coherence length $\xi \sim 8.1 \times 10^{-7}$ cm. It is a bit difficult to estimate the renormalized Fermi velocity from experimental measurements. Bauer et al. obtained $v_p^* \sim 5.29 \times 10^5$ cm/s from the data of $dH_{c2}/dT$ and the specific-heat coefficient, assuming a spherical Fermi surface. This value of $v_p^*$ is almost of the same order as that obtained by combining the unrenormalized Fermi velocity computed from the LDA method and the mass enhancement factor $z^{-1} \sim 100$ estimated from the specific-heat measurement. According to the LDA band calculations, the spin-orbit splitting is not so small compared to the Fermi energy, and may be approximated as $\alpha p_F/E_F \sim 0.1$. Then, for CePt$_3$Si, we have

$$
\begin{align*}
\delta D_{loc}(0) & \sim 1.0 \times 10^{-24} \sqrt{\frac{H}{H_{c2}}} \pm 0.48 \times 10^{-24} \frac{H}{H_{c2}}.
\end{align*}
$$

It is remarkable that the contribution from the paramagnetic supercurrent [the second term of Eq. (43)] is comparable to that from the Meissner supercurrent [the first term of Eq. (43)]. It should be stressed that the feasibility of the experimental observation of Edelstein’s magnetoelectric effect is due to the large mass enhancement in the heavy-fermion system, which suppresses strongly the Meissner supercurrent, but in contrast, does not affect the Zeeman-field-induced paramagnetic supercurrent. Moreover, if the superfluid density $n_s$ is reduced by the backflow effect, the Meissner term of Eq. (42) is more suppressed compared with the paramagnetic term, and thus the observation of the magnetoelectric effect may become easier.

VI. SUMMARY AND DISCUSSION

We have investigated electron correlation effects on the magnetoelectric transport phenomena in superconductors
without inversion symmetry. It is found that, in contrast to the Meissner diamagnetic supercurrent which is much reduced by the mass enhancement factor in the absence of translational symmetry, the Zeeman-field-induced paramagnetic supercurrent is not affected by the strong electron correlation provided that ferromagnetic fluctuation is not developed. Because of this remarkable property, the experimental detection of the magnetoelectric effect may be more feasible in heavy-fermion superconductors without inversion symmetry such as CePt$_3$Si, where the enormous mass enhancement suppresses the magnitude of the Meissner supercurrent, than in conventional metals with moderate effective electron mass. We have proposed the experimental setup for the observation of the magnetoelectric effect in CePt$_3$Si which utilizes the Volovik effect. It has been also pointed out that in noncentrosymmetric $p$ wave superconductors, the coherence effect on the nuclear relaxation rate $1/T_1T$ is similar to that of conventional $s$-wave superconductors.

Finally, we would like to comment on the implication of our results for UIr, which is the recently discovered ferromagnetic superconductor without inversion symmetry. UIr exhibits superconductivity under high pressure in the vicinity of the phase boundary between ferromagnetic and nonmagnetic states. The resistivity of this system increases remarkably as the applied pressure approaches the critical value at which the ferromagnetism disappears, indicating the existence of ferromagnetic critical fluctuation. In this case, the magnetoelectric coefficient $\varepsilon(T)$ may be enhanced by the three-point vertex functions $\Lambda^{x\pm\pm}_T$ of which the magnitudes are much increased by ferromagnetic fluctuation, provided that the spin easy axis is taken as the $x$ axis. Since the crystal structure of UIr is monoclinic, and does not possess any mirror planes, the Rashba spin-orbit interaction with the $\vec{\alpha}$ vector perpendicular to the spin easy axis should always exist. Thus, the magnetoelectric effect strongly enhanced by ferromagnetic fluctuation may be observed in UIr under an applied magnetic field parallel to the spin easy axis.

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24. If the attached leads are made of superconductors with the transition temperature higher than that of the sample and the same pairing symmetry as that of $\Delta_0(p)$ and $\Delta_1(p)$, the Zeeman-field-induced supercurrent can flow in the system. Note that the tunneling amplitude between the leads and the noncentrosymmetric superconductor does not depend on the symmetric properties of the $\vec{d}$ vector of the gap function (12), $\vec{d}_S \times \vec{n}$, as pointed out by I. A. Sergienko and S. H. Curnoe, in Ref. 18. The observation of the magnetoelectric effect may be possible also in this case, as discussed in Sec. V.