Dynamical Induction of \textit{s}-wave Component in \textit{d}-wave Superconductor Driven by Thermal Fluctuations

Atsuya KUMAGAI\textsuperscript{*} and Hiromichi EBISAWA\textsuperscript{**}

Graduate School of Information Sciences, Tohoku University, Sendai 980-8579

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We investigated the mutual induction effects between the \textit{d}-wave and the \textit{s}-wave components of order parameters due to superconducting fluctuation above the critical temperatures and calculated its contributions to paraconductivity and excess Hall conductivity based on the two-component stochastic TDGL equation. It is shown that the coupling of two components increases paraconductivity while it decreases excess Hall conductivity compared to the cases when each component fluctuates independently. We also found the singular behavior in the paraconductivity and the excess Hall conductivity dependence on the coupling parameter which is consistent with the natural restriction among the coefficients of gradient terms.

KEYWORDS: superconducting fluctuation, TDGL theory, two-component superconductor, paraconductivity, excess Hall conductivity

\section{Introduction}

Recently many experimental results indicate the mixed pairing symmetry of Cooper pairs in high-$T_c$ cuprates which consists of \textit{d}- and \textit{s}-wave-like components.\textsuperscript{1, 2} Theoretically, it is shown that certain classes of microscopic models lead to dominant \textit{d}-wave pairing with subdominant \textit{s}-wave pairing.\textsuperscript{3} Macroscopic properties of superconductors of these mixed symmetry can be studied by means of the Ginzburg-Landau (GL) thoery for two-component order parameter. Phenomenologically, Joynt\textsuperscript{3} first introduced the GL free energy of two-component superconductors based on the group theoretical consideration.\textsuperscript{4} Microscopic derivations of the corresponding two-component GL equations have been performed by several authors.\textsuperscript{5, 6, 7}

Due to those macroscopic arguments, it is shown that, in tetragonal lattice, the coexistence of the two components is not expected in the bulk, but the subdominant order parameter is induced by spatial variations because two components are coupled through “mixed gradient” term; this effect has been intensively investigated under the circumstances such as vortices,\textsuperscript{8} impurities.\textsuperscript{9}

\textsuperscript{*} E-mail: kumagai@cmt.is.tohoku.ac.jp
\textsuperscript{**} E-mail: ebi@cmt.is.tohoku.ac.jp
currents etc., thereby their anisotropic features being clarified. Such inductions have been also studied from more microscopic standpoints. Nonetheless there is so far no considerations of the effect beyond these static cases. There should be dynamical situations where the order parameter varies spatially so that the subdominant component is induced. We now point out this possibility, which is driven by the superconducting fluctuation above critical temperature, and give an estimation of the effect to the paraconductivity and the excess Hall conductivity. It should be stressed that, within our considerations, the origin of the induction is superconducting fluctuation itself, unlike the recent studies treating vortex states based on two-component TDGL equation, where the induction is attributed to intrinsic structures of vortices.

The conduction process which contributes to the enhancement of electrical conductivity due to superconducting fluctuation is first pointed out by Aslamasov and Larkin (AL). Besides, there is a prediction about the behavior of Hall conductivity resulting from AL process, though it is not yet clear whether it accounts for the Hall anomaly. AL process can also be described within the framework of Time-Dependent Ginzburg-Landau (TDGL) equation and several formulations have been done based on the stochastic TDGL equation.

In this work, based on the stochastic two-component TDGL equation, we show our calculation for the paraconductivity and the excess Hall conductivity and discuss the importance of the induction effect of the subdominant component through the superconducting fluctuation. Since we limit the considerations to the case of weak magnetic field, we will make the simple linear or bilinear expansion of the response current in terms of the external electric and magnetic field.

The characteristic features of the induction effect we will show is firstly that no anisotropy arises in spite of the anisotropy in GL free energy and secondly that the induction reduces the excess Hall conductivity while it enhances the paraconductivity. The results will show the singular behavior of these transport coefficients in increment of the coupling of two components through the spatial variation. This is direct reflection of the stability condition of the normal state against the superconducting fluctuation.

In Sec. we first give the formulations of fluctuation conductivity tensor based on the two-component TDGL equation as a stochastic differential equation. In Sec. we show the numerical results and see how the coupling of two components affects them. In Sec. we summarize the results and give some discussions. In this paper we set $\hbar = k_B = c = 1$.

§2. Formulation for the stochastic TDGL equation

In this section we show the formulation for the paraconductivity and excess Hall conductivity based on the stochastic two-component TDGL equation. In the following, electric and magnetic field is expressed as scalar and vector potential $\phi(r), A(r)$, respectively. The charge of a Cooper pair is $e^* = -2e < 0$. We introduce a differential operator $\Pi = -i\nabla - e^* A(r)$. 
We start with the two-component TDGL equation

\[ \gamma_d \left( \frac{\partial}{\partial t} + ie^* \phi \right) d = - \left( \frac{\Pi^2}{2m_d} + \alpha_d \right) d - \frac{(\Pi_y^2 - \Pi_x^2)}{2m_v} s + f_1 \]  
(2.1a)

\[ \gamma_s \left( \frac{\partial}{\partial t} + ie^* \phi \right) s = - \left( \frac{\Pi^2}{2m_s} + \alpha_s \right) s - \frac{(\Pi_y^2 - \Pi_x^2)}{2m_v} d + f_2 \]  
(2.1b)

\[ j(r) = \frac{e^*}{m_d} (d^* \Pi d + \text{c.c.}) + \frac{e^*}{m_s} (s^* \Pi s + \text{c.c.}) + \frac{\tilde{x} e^*}{m_v} (d^* \Pi_x s + s^* \Pi_x d + \text{c.c.}) + \frac{\tilde{y} e^*}{m_v} (d^* \Pi_y s + s^* \Pi_y d + \text{c.c.}) \]  
(2.2)

which is based on Joynt’s two-component GL free energy\(^3\) and also can be derived from such lattice models as \(t\)-\(J\) model\(^2\). \(d\) and \(s\) mean superconducting order parameters of \(d\)-wave and \(s\)-wave component, respectively. Both GL coefficients \(\alpha_d\) and \(\alpha_s\) are positive because the system is above the transition temperatures. We added the white noise source terms \(f_1, f_2\) as representations of thermodynamical fluctuations, which make the mean square expectation value of order parameter finite even above the transition temperature. It is the main characteristics of the TDGL equation that the two components are coupled via so-called “mixed gradient” terms. The main purpose of this paper is to clarify the effects of this term on the electric transport properties, and this effects are represented as the coefficient \(m^{-1}\). Because of this anisotropic term we should have set electrical field as \(E = (E_x, E_y, 0)\) for more generality, but it will be found later that this anisotropy is cancelled with that of current and no anisotropy arises in the fluctuation conductivity tensor. For this reason we set \(\phi(r) = -E_x, A(r) = Bx\tilde{y}\).

First, we rewrite the TDGL equation(2.1) into the simpler form. Keeping up to bilinear terms in external fields, Fourier transformed TDGL equation becomes

\[ i\gamma_\omega \psi(k, \omega) = - \left( \eta_k + \gamma V_k \frac{\partial}{\partial k_x} \right) \psi(k, \omega) + f(k, \omega), \]  
(2.3)

\[ V_k = e^*(E - iB\gamma^{-1}k_y W_y) \]  
(2.4)

where

\[ \psi = \begin{pmatrix} d \\ s \end{pmatrix}, \]  
(2.5)

\[ f = \begin{pmatrix} f_1 \\ f_2 \end{pmatrix}, \]  
(2.6)

\[ \eta_k = \frac{k^2}{2} W_x + \frac{k^2}{2} W_y + \alpha, \]  
(2.7)

\[ W_x = \begin{pmatrix} m_d^{-1} & -m_v^{-1} \\ -m_v^{-1} & m_s^{-1} \end{pmatrix}, \]  
(2.8)
Why = (m_d^{-1} m_v^{-1})

\alpha = \begin{pmatrix} \alpha_d & 0 \\ 0 & \alpha_s \end{pmatrix},

\gamma = \begin{pmatrix} \gamma_d & 0 \\ 0 & \gamma_s \end{pmatrix}.

Then we can formally express the solution of the TDGL equation:

\psi(k, \omega) = G(k, \omega) f(k, \omega)

(2.12)

where

G(k, \omega) \equiv \left( i\gamma \omega + \eta_k + \gamma V_k \frac{\partial}{\partial k_x} \right)^{-1}.

(2.13)

Here we assume the correlation function between white noises as

\langle f(k', \omega') f^\dagger(k, \omega) \rangle = a \delta_{kk'} \delta_{\omega \omega'},

(2.14)

where the coefficient \( a \) is 2 \times 2 matrix and should be determined from the linearized two-component GL free energy

F_k = \psi(k)^\dagger \eta_k \psi(k).

(2.15)

\psi(k) means the order parameter of wave vector \( k \) at the specific time. Corresponding distribution function is

Z_k = \int e^{-F_k/T} dd_1 dd_2 ds_1 ds_2 = \left( \frac{\pi T}{2} \right)^2 \det \eta_k,

(2.16)

where indices 1 and 2 mean real and imaginary part, respectively. This distribution function gives rise to

\langle \psi(k) \psi(k)^\dagger \rangle_0 = -T \left( \frac{\partial}{\partial \eta_{11}} \frac{\partial}{\partial \eta_{22}} \right) \ln Z_k = T \eta_k^{-1}.

(2.17)

\langle \cdots \rangle_0 means thermal average in the absence of external fields. The stochastic TDGL equation(2.1) also realizes this equilibrium values, provided we set

a = T(\gamma + \gamma^\dagger).

(2.18)

In fact, from eq.(2.14) with eq.(2.18), we obtain

\langle \psi(k) \psi(k)^\dagger \rangle_0 = \sum_\omega \langle \psi(k, \omega) \psi^\dagger(k, \omega) \rangle_0

= \sum_\omega G_0(k, \omega) \langle f(k, \omega) f^\dagger(k, \omega) \rangle G_0^\dagger(k, \omega)

= T \int_{-\infty}^{\infty} \frac{d\omega}{2\pi i} [(\omega - i\gamma^{-1} \eta_k)^{-1} \eta_k^{-1} - \eta_k^{-1}(\omega + i\eta_k \gamma^{-1})^{-1}]

= T \eta_k^{-1}

(2.19)
where $G_0(k, \omega) = (i \gamma \omega + \eta_k)^{-1}$ stands for the Green’s function of TDGL equation (2.3) in the absence of external fields.

Next we perform expansions with respect to external fields $V_k$. Resolvent $G(k, \omega)$ is expanded as follows:

$$G(k, \omega) = G_0(k, \omega) + G_1(k, \omega) + G_2(k, \omega) + \cdots.$$  \hspace{1cm} (2.20)

Note that, in the above expansion, each term acts on any function $h(k)$ as a differential operator except $G_0(k, \omega)$:

$$G_1(k, \omega)h(k) = -G_0(k, \omega)\gamma V_k \frac{\partial}{\partial k_x}[G_0(k, \omega)h(k)],$$  \hspace{1cm} (2.21)

$$G_2(k, \omega)h(k) = G_0(k, \omega)\gamma V_k \frac{\partial}{\partial k_x} \left[ G_0(k, \omega) \gamma V_k \frac{\partial}{\partial k_x} [G_0(k, \omega)h(k)] \right].$$  \hspace{1cm} (2.22)

In contrast, $G^\dagger(k, \omega)$ acts on functions at the left side.

On the other hand, the expectation value of supercurrent (2.2) is written as the thermal averaged form as

$$\langle j(r) \rangle = e^* \sum_k \sum_{k'} \sum_\omega \sum_{\omega'} \text{tr} [R(k, \omega; k', \omega')(\hat{K}(k) - e^*\hat{A})] e^{(k'-k)\cdot r}$$  \hspace{1cm} (2.23)

where

$$\hat{K}(k) = \begin{pmatrix} m_d^{-1}k & m_v^{-1}k \\ m_v^{-1}k & m_s^{-1}k \end{pmatrix}, \quad \hat{k} = (-k_x, k_y),$$  \hspace{1cm} (2.24)

$$\hat{A} = \begin{pmatrix} m_d^{-1}A & m_v^{-1}A \\ m_v^{-1}A & m_s^{-1}A \end{pmatrix}, \quad \hat{A} = (-A_x, A_y),$$  \hspace{1cm} (2.25)

and

$$R(k, \omega; k', \omega') = \langle \psi(k', \omega') \psi^\dagger(k, \omega) \rangle.$$  \hspace{1cm} (2.26)

Corresponding to the resolvent expansion (2.20), the correlation function $R(k, \omega; k', \omega')$ is expanded as follows:

$$R(k, \omega; k', \omega') = G(k', \omega')(f(k', \omega')f^\dagger(k, \omega))G^\dagger(k, \omega)$$

$$= R_0(k, \omega; k', \omega') + R_1(k, \omega; k', \omega') + R_2(k, \omega; k', \omega') + \cdots$$  \hspace{1cm} (2.27)

where

$$R_0(k, \omega; k', \omega') = (2\pi)^4 G_0(k', \omega') \alpha \delta(k - k') \delta(\omega - \omega'),$$  \hspace{1cm} (2.28)

$$R_1(k, \omega; k', \omega') = (2\pi)^4 [G_1(k', \omega') \alpha \delta(k - k') \delta(\omega - \omega')G_0^\dagger(k, \omega)$$

$$+ G_0(k', \omega') \alpha \delta(k - k') \delta(\omega - \omega')G_1^\dagger(k, \omega)],$$  \hspace{1cm} (2.29)

$$R_2(k, \omega; k', \omega') = (2\pi)^4 [G_2(k', \omega') \alpha \delta(k - k') \delta(\omega - \omega')G_0^\dagger(k, \omega)$$

$$+ G_1(k', \omega') \alpha \delta(k - k') \delta(\omega - \omega')G_1^\dagger(k, \omega)$$

$$+ G_0(k', \omega') \alpha \delta(k - k') \delta(\omega - \omega')G_2^\dagger(k, \omega)].$$  \hspace{1cm} (2.30)
Now we can calculate supercurrent from the correlation functions. The expressions for the supercurrent in each order are as follows:

\[ \langle j_0(r) \rangle = e^s \sum_k \sum_{k'} \sum_\omega \sum_{\omega'} \text{tr}[R_0(k; \omega; k', \omega') \tilde{K}(k)] e^{i(k'-k) \cdot r}, \]  

\[ \langle j_1(r) \rangle = e^s \sum_k \sum_{k'} \sum_\omega \sum_{\omega'} \text{tr}[R_1(k; \omega; k', \omega') \tilde{K}(k) - R_0(k; \omega; k', \omega') e^s \tilde{A} e^{i(k'-k) \cdot r}, \]  

\[ \langle j_2(r) \rangle = e^s \sum_k \sum_{k'} \sum_\omega \sum_{\omega'} \text{tr}[R_2(k; \omega; k', \omega') \tilde{K}(k) - R_1(k; \omega; k', \omega') e^s \tilde{A} e^{i(k'-k) \cdot r}. \]  

Here we consider a thin film with thickness \( d \), so that we calculate the summation as integral with respect to \( k_x \) and \( k_y \), and extract only the contributions from \( k_z = 0 \). Thus we have the expression for the paraconductivity from (2.32)

\[ \sigma_{xx} = \frac{e^{s}T}{d} \sum_\omega \sum_k \text{tr} \left\{ -k^2_x (G_0 W_x G_0 W_x G_0 \gamma G_0 + \text{h.c.}) + G_0^T \gamma G_0^T W_x G_0 \right\} (\gamma + \gamma^T) \]  

and for the excess Hall conductivity from (2.33)

\[ \sigma_{yx} = \frac{ie^{s}TB}{d} \sum_\omega \sum_k \text{tr} \left\{ [k^2_y W_y (1 - 2k^2_x G_0 W_x) G_0 W_x G_0 (W_y G_0 W_x G_0 \gamma + \gamma G_0 W_y) - \text{h.c.}] \right. \]
\[ - \left\{ k^2_x k^2_y W_y G_0 W_x G_0 W_y G_0 W_x G_0 \gamma + \gamma G_0 W_x G_0 W_y \right\} - \text{h.c.} \]
\[ + \left\{ k^2_x k^2_y (W_y G_0^T W_x G_0^T W_y G_0 W_x G_0 \gamma - \text{h.c.}) G_0 (\gamma + \gamma^T) G_0^T \right\} \]

where \( G_0 \) means \( G_0(k; \omega) \).

From the above expression we find that both \( \sigma_{xx} \) and \( \sigma_{yx} \) are even functions of \( m_i^{-1} \) and therefore we obtain \( \sigma_{yy} = \sigma_{xx} \), \( \sigma_{xy} = -\sigma_{yx} \); no anisotropy arises in the fluctuation conductivity tensor in contrast to recent theoretical predictions about vortex state below the critical temperature. [10, 11]

Since the above expression for conductivity tensor is complicated, it may be helpful to compare it with the former studies in the case of \( m_i^{-1} = 0 \). In this case the coupling of two components disappears and the problem must reduce to that in the conventional single component superconductor, which has been studied Aslamasov and Larkin(AL) [10] and Fukuyama, Ebisawa and Tsuzuki(FET) [11]. Then all matrices become diagonal and actually one can confirm

\[ \sigma_{xx} = \frac{e^{s}T}{d} \sum_{i=d,s} \sum_\omega \sum_k (\gamma_i \gamma^*_i) [-k^2_x (m_i^{-2} \gamma_i g_i^* g_i^3 + \text{c.c.}) + m_i^{-1} \gamma_i g_i^2 g_i] \]
\[ = \sum_{i=d,s} \frac{e^{s}T |\gamma_i|^2}{2\pi d \alpha_i \gamma_i \gamma_i^*} \]  

and

\[ \sigma_{yx} = \frac{ie^{s}TB}{d} \sum_{i=d,s} \sum_\omega \sum_k (\gamma_i \gamma^*_i) [2k^2_y m_i^{-3} \gamma_i g_i^4 g_i^* - 6k^2_x k^2_y m_i^{-4} \gamma_i g_i^5 g_i^* + k^2_x k^2_y m_i^{-4} \gamma_i g_i^3 g_i^3 - \text{c.c.}] \]
\[ = - \sum_{i=d,s} \frac{e^{s}TB |\gamma_i|^2 \gamma_i \gamma_i^*}{6\pi d \alpha_i \gamma_i \gamma_i^*} \]  

(2.36)
where we introduced single component Green’s functions

\[ g_i = \frac{1}{i\gamma_i \omega + \eta_i}, \quad \eta_i = \frac{k^2}{2m_i} + \alpha_i. \]  

(2.38)

and

\[ \gamma_{i1} = \Re \gamma_i, \quad \gamma_{i2} = \Im \gamma_i \]  

(2.39)

These results correspond to AL and FET, respectively.

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§3. Numerical results

In this section we show the numerical results of paraconductivity (2.34) and excess Hall conductivity (2.35). In Fig.1 we show the plot of them as functions of temperature for various

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Fig. 1. Temperature dependence of (a) paraconductivity \( \sigma_{xx} \) and (b) excess Hall conductivity \( \sigma_{yz} \) for the coupling parameter

\[ m_v^{-1} = 0 (\text{solid line}), m_v^{-1} = 3 (\text{long dashed line}), m_v^{-1} = 4 \text{ (short dashed line)} \]  

where \( T_d = 20, T_s = 16, m_d^{-1} = 3, m_s^{-1} = 6, \gamma_d = 12 + 4i, \gamma_s = 4 - 10i. \) Here we incorporated the temperature dependence by setting the GL parameters \( \alpha_d = T - T_d \) and \( \alpha_s = T - T_s. \)
coupling parameter $m_v^{-1}$. Here the critical temperature of the $d$-wave component is chosen to be higher than that of $s$-wave, and as a result, both $\sigma_{xx}$ and $\sigma_{yx}$ diverges at $T_d$ even in the presence of the coupling of the two component order parameters. It seems that $\sigma_{xx}$ is enhanced but $\sigma_{yx}$ is suppressed due to the coupling of the two component order parameters. In order to see the coupling effects more clearly, we show the above results as functions of $m_v^{-1}$ in Fig. 2. In the light of GL free energy (2.15), $\det W_x > 0$ is required for the stability of the system against the spatial variation of order parameters. This condition is reflected in the singularities at $m_v^{-1} = \pm \sqrt{m_d^{-1}m_s^{-1}}$ in the fluctuation conductivity tensor.

Fig. 2. Coupling parameter dependence of (a) paraconductivity $\sigma_{xx}$ and (b) excess Hall conductivity $\sigma_{yx}$ at fixed temperature $T = 22$ with the same set of parameters as in Fig.1.

§4. Summary and Discussions

In this work we gave the formulation for the paraconductivity and the excess Hall conductivity in the presence of the coupling of two-component order parameter in the limit of weak external
fields. The phenomenological formulations based on the stochastic TDGL equation have been done by several authors and these correspond to AL process microscopically. Our calculation is considered to be their extension. We also numerically calculated the dependences on temperature and coupling parameter with the typical parameters and found that the coupling of two components enhances paraconductivity but reduces excess Hall conductivity. Such tendency does not change even for the other set of parameters.

We paid attention to the stability condition about GL coefficients and saw the singular behavior of conductivity tensor at critical value $m^{-1} = \pm \sqrt{m_d^{-1} m_s^{-1}}$, though it could have not been attained if one had resorted to perturbative expansions with respect to $m^{-1}$.

It is well known that excess Hall conductivity vanishes identically in the absence of imaginary part of order parameter relaxation time within the conventional framework of TDGL theory. On the other hand, several recent studies investigating vortex states based on two-component TDGL equations have pointed out that the Hall effect is also caused by the anisotropy of crystal lattice even if the relaxation times are real. However, our calculation in the fluctuation regime resulted in vanishing Hall conductivity in the case of $\text{Im} \gamma_d = \text{Im} \gamma_s = 0$, and moreover, we found that no anisotropy arose in the conductivity tensor even for finite value of $m^{-1}$.

Finally we mention an aspect of multiple transition in unconventional superconductors from the viewpoint of superconducting fluctuation. If the transition temperatures are close to one another in the presence of the coupling between components, the induction effect in superconducting fluctuation would gain its significance. In this context, the recent experiments through microwave surface impedance which suggest multi-component superconductivity can present an interesting system if those components correspond to order parameters of different symmetry. To compare with such experiments, we need to carry out the extention so as to include the frequency dependence in paraconductivity, which will be given elsewhere.

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excess Hall conductivity

(b)

temperature
paraconductivity

(a)
(b)