The spin Hall effect (SHE) is one of the promising phenomena to utilize a spin current as spintronics devices, and the theoretical understanding of its microscopic mechanism is essential to know how to control its response. Although the SHE in multiorbital systems without inversion symmetry (IS) is expected to show several unique properties due to cooperative roles of orbital degrees of freedom and lack of IS, the theoretical understanding of the cooperative roles has been lacked. To clarify the cooperative roles, we study the spin Hall conductivity (SHC) derived by the linear-response theory for a $t_{2g}$-orbital tight-binding model of the [001] surface or interface of Sr$_2$RuO$_4$ in the presence of dilute nonmagnetic impurities. We find that the band anticrossing, arising form the combination of orbital degrees of freedom and lack of IS, causes the magnitude increase and sign change of the SHC at some nonmagnetic impurity concentrations. Since the similar mechanism for controlling the magnitude and sign of the response of Hall effects works in other multiorbital systems without IS, our mechanism gives a new method to control the magnitude and sign of the response of Hall effects in some multiorbital systems by introducing IS breaking and tuning the nonmagnetic impurity concentration.

PACS numbers: 72.25.Ba, 73.40.-c, 74.70.Pq

In general, the intrinsic SHE can be understood in terms of the effective flux, generated by an adiabatic process of an electron in the real space, since this effective flux gives the phase factor of that wave function which is similar to the Aharonov-Bohm phase in the presence of an external magnetic field. For example, in a multiorbital TM or TM oxide with inversion symmetry (IS), the spin-dependent effective flux is generated by the adiabatic process by using the $z$-component of the atomic spin-orbit interaction (SOI), the interorbital hopping integral between the orbitals connected by that $z$-component, and other hopping integrals [see Fig. 1(a)]. In addition, the SHE in the two-dimensional electron gas without IS whose electronic state is described by using the Rashba-type antisymmetric SOI can be understood by the adiabatic process by using the Rashba-type antisymmetric SOI and the intraorbital hopping integrals [see Fig. 1(b)].

Although there are a lot of studies about the SHE in multiorbital systems with IS (e.g., Refs. [8, 13]) or single-orbital systems without IS (e.g., Refs. [5, 22]), the characteristic properties of the SHE in a multiorbital system without IS have been little understood. In particular, cooperative roles of orbital degrees of freedom and lack of IS have not been clarified yet, although their combination will lead to several unique properties of the SHE.

To clarify these roles, it is necessary to study the SHE in a multiorbital system without IS by using a model considering both orbital degrees of freedom and IS breaking correctly. It should be noted that the correct treatment beyond the Rashba-type antisymmetric SOI is significant for multiorbital systems since the combination of these leads to completely different results of several electronic properties from the results for the Rashba-type antisymmetric SOI. Actually, the momentum dependence of the $d$-vector of a Cooper pair completely differs from that for the Rashba-type antisymmetric SOI due to the difference in the momentum dependence of the antisymmet-

![FIG. 1. (Color online) Schematic pictures of some adiabatic processes generating the effective flux, $\Phi_{\text{eff}}$, in (a) the $t_{2g}$-orbital model with IS and (b) the single-orbital Rashba model without IS.](image-url)
ric SOI; in the correct treatment, the antisymmetric SOI arises from the atomic SOI and the interorbital hopping integral due to the local parity mixing, induced by IS breaking [24]. Since the antisymmetric SOI is important even in discussing the SHE, a study about the SHE by using the correct treatment is highly desirable.

In this Letter, we study the SHE in a $t_{2g}$-orbital system without IS by using the correct treatment about orbital degrees of freedom and lack of IS beyond the Rashba-type antisymmetric SOI and reveal their cooperative roles in the SHE. In particular, we find that the band anticrossing due to the cooperative roles plays an important role in controlling the magnitude and sign of the spin Hall conductivity (SHC) of a multiorbital system without IS in the presence of dilute nonmagnetic impurities. After discussing applicability of the similar mechanism, we propose the ubiquitous method to control the magnitude and sign of Hall effects by using orbital degrees of freedom, IS breaking, and nonmagnetic impurity scattering.

To discuss the SHE in a multiorbital system without IS, we consider a $t_{2g}$-orbital tight-binding model of the [001] surface or interface of Sr$_2$RuO$_4$ [Fig. 2(a)] in the presence of dilute nonmagnetic impurities: the Hamiltonian is $H = H_0 + H_{LS} + H_{SB} + H_{imp}$, where $H_0 = \sum_{k} \sum_{\alpha, \beta} \varepsilon_{\alpha \beta}(k) c_{k \alpha}^\dagger c_{k \beta}$, $H_{LS} = \sum_k \sum_{\alpha, \beta} \varepsilon_{\alpha \beta}(k) c_{k \alpha}^\dagger c_{k \beta}$, $H_{SB} = \sum_k \sum_{\alpha, \beta} \varepsilon_{\alpha \beta}(k) c_{k \alpha}^\dagger c_{k \beta}$, $H_{SB} = \sum_k \sum_{\alpha, \beta} \varepsilon_{\alpha \beta}(k) c_{k \alpha}^\dagger c_{k \beta}$, and $H_{imp} = \sum_k c_{k \alpha}^\dagger \sum_{\alpha, \beta} \eta_{\alpha \beta} \sum_i \eta_{i, \alpha \beta} c_{k \alpha} c_{k \beta}$. We choose these parameters so as to reproduce the $t_{2g}$-orbital with $\lambda = 0.45$, $\varepsilon_{\alpha \beta}(k) = -2t_{2g} \cos k_x - 2t_{2g} \cos k_y - \mu$, $\varepsilon_{d_{z^2}}(k) = -2t_{d_{z^2}} \cos k_x - 2t_{d_{z^2}} \cos k_y - \mu$, and $\varepsilon_{d_{x^2-y^2}}(k) = -2t_{d_{x^2-y^2}} \cos k_x - 2t_{d_{x^2-y^2}} \cos k_y - \mu$. Therefore, we can treat its effects by the Born approximation; thus, $H_{SB}^{\text{int}}$ represents the interorbital hopping integral between the $d_{x^2-y^2}$ and $d_{z^2}$ orbitals due to the local parity mixing, induced by IS breaking near the [001] surface or interface [24].

$V_{\text{ISB}}(k) = V_{\text{ISB}}(k) = 2t_{\text{ISB}} \sin k_x \sin k_y$. Note that the second-order perturbation of $H_{LS}$ and $H_{SB}$ gives the Rashba-type antisymmetric SOI when the orbital degeneracy is lifted by the large crystal-electric-field energy [24]. In a case without IS, we set $t_{\text{ISB}} = 0.09$ eV to make $t_{\text{ISB}}$ larger than $\lambda$ and $t_5$; as we will show, this condition is essential to obtain the magnitude and sign changes of the SHC. For comparison, we also consider the case at $t_{\text{ISB}} = 0$ eV.

Diagonalizing $H_{\text{band}} = \sum_k H_{\text{band}}(k) = H_0 + H_{\text{LS}} + H_{\text{SB}}$, we obtain the band dispersions, $E_\alpha(k)$, and the unitary matrix $[U(k)]_{\eta \eta'}$, where $\eta$ is $\eta \equiv (\alpha, \sigma)$. Comparing the Fermi surfaces at $t_{\text{ISB}} = 0$ and $0.09$ eV, shown in Figs. 2(b) and 2(c), we see IS breaking causes the spin splitting of the $\alpha$, $\beta$, and $\gamma$ sheets. (At $t_{\text{ISB}} = 0$ eV, the $\alpha$ and $\beta$ sheets are formed mainly by the $d_{x^2}$ and $d_{z^2}$ orbitals, and the $\gamma$ sheet is formed mainly by the $d_y$ orbital [27].) In particular, the most drastic effect of IS breaking is the change of the curvature of the $\gamma$ sheet around $k \sim (\pi, \pi)$ due to the band anticrossing between the $d_{x^2}$ and $d_{z^2}$ orbitals. In other words, this band anticrossing causes the change of the main orbital forming the $\gamma$ sheet around $k \sim (\pi, \pi)$ from the $d_y$ orbital to the $d_{x^2}$ orbital. The necessary conditions for this band anticrossing are both that $t_{\text{ISB}}$ is larger than $\lambda$ and that several Fermi surface sheets are close to each other in a certain area of the Brillouin zone.

$H_{imp}$ represents the local scattering potential due to dilute nonmagnetic impurities. Assuming that the scattering is weak, we can treat its effects by the Born approximation; thus, $H_{imp}$ affects the SHC through the self-energy correction and the current vertex correction due to the four-point vertex function [8, 15, 23]. For simplicity, we assume that the main effects of $H_{imp}$ arise from the band-independent quasiparticle damping, $\gamma_{\text{imp}}$, due to the self-energy correction, which is proportional to the nonmagnetic impurity concentration, $n_{\text{imp}}$. We have checked the neglected terms do not qualitatively change the results shown below [28].

Then, we derive the SHC by the linear-response theory [13], $\sigma_{xy}^{S(I)} = \sigma_{xy}^{S(I)} + \sigma_{xy}^{S(II)}$, where $\sigma_{xy}^{S(I)}$ is the Fermi surface term, $\sigma_{xy}^{S(I)} = \frac{1}{N_k} \sum_{\{\eta\}} \int_{-\infty}^{\infty} \frac{d\omega}{2\pi} \left( \frac{\partial f(\omega)}{\partial \omega} \right) \left[ f_x(k) \right]_{\eta_1, \eta_2} \left[ f_y(k) \right]_{\eta_1, \eta_2}$ and $\sigma_{xy}^{S(II)}$ is the Fermi sea term, $\sigma_{xy}^{S(II)} = \frac{1}{N_k} \sum_{\{\eta\}} \int_{-\infty}^{\infty} \frac{d\omega}{2\pi} f(\omega) \text{Re} \left\{ \left[ f_x(k) \right]_{\eta_1, \eta_2} \left[ f_y(k) \right]_{\eta_1, \eta_2} \left[ f_y(k) \right]_{\eta_1, \eta_2} \right\}$. Here $\sum_{\{\eta\}}$ is $\sum_{\{\eta\}} = \sum_{\eta_1} \sum_{\eta_2} \sum_{\eta_3} f(\omega)$ is the Fermi distribution function, $f(\omega) = \frac{1}{e^{\beta} + 1}$, $G^{(A)}(k, \omega)$ is the retarded (advanced) Green’s function, $\left[ G^{(A)}(k, \omega) \right]_{\eta_1, \eta_2} = \sum_\omega \left[ U^0(k) \right]_{\eta_1} \left( \omega - E_k(k) + i \eta_{\text{imp}} \right) \left[ U^0(k) \right]_{\eta_2}$, $f_x(k)$ is the charge current operator, $f_x(k) = -e \partial \text{Re} f(\omega) \frac{\partial f(\omega)}{\partial \omega} \frac{\partial f(\omega)}{\partial \omega} \left[ f_y(k) \right]_{\eta_1, \eta_2}$, $f_y(k)$ is the

FIG. 2. (Color online) Schematic picture of (a) the situation considered, and Fermi surfaces at (b) $t_{\text{ISB}} = 0$ and (c) $0.09$ eV.
spin current operator, \( J^S_x(k) \equiv \frac{1}{2} \sum_{\sigma} \partial H_{\text{imp}}(k) / \partial k_x \), and \( (g \frac{\partial}{\partial \omega})h \) is \( g \frac{\partial}{\partial \omega} \). It should be noted, first, that in the clean limit, where \( \gamma_{\text{imp}} \) satisfies \( \gamma_{\text{imp}} < \triangle E(k) \) with \( \triangle E(k) \) being the band splitting giving the dominant contribution to \( \sigma^S_{xy} \), \( \sigma^S_{xy} \) is independent of \( \gamma_{\text{imp}} \) and is given mainly by the Berry curvature term, part of \( \sigma^S_{xy}(B) \); second, that in the dirty limit, where \( \gamma_{\text{imp}} \) satisfies \( \gamma_{\text{imp}} \gg \triangle E(k) \), \( \sigma^S_{xy} \) is proportional to \( \gamma^{-3} \) and is given mainly by \( \sigma^S_{xy}(T) \) and \( \sigma^S_{xy}(L) \).

We turn to results, obtained by using 20000 \times 20000 meshes of the first Brillouin zone. This size is necessary to suppress the finite-size effect in a clean region.

We first compare \( \gamma_{\text{imp}} \) dependence of \( \sigma^S_{xy} \) at \( t_{\text{SB}} = 0 \) and 0.09 eV in Fig. 3(a). We see three changes due to IS breaking: (i) an increase in a clean region \( (\gamma_{\text{imp}} < 0.45 \text{ meV}) \), (ii) a sign change in a slightly dirty region \( (0.45 \text{ meV} \leq \gamma_{\text{imp}} \leq 4.5 \text{ meV}) \), and (iii) the appearance of a minimum at \( \gamma_{\text{imp}} = 4.5 \text{ meV} \). These results suggest that the magnitude and sign of \( \sigma^S_{xy} \) can be controlled by using orbital degrees of freedom, IS breaking, and nonmagnetic impurity scattering. Note that in the range of \( \gamma_{\text{imp}} \) shown in Fig. 3(a), \( \sigma^S_{xy}(I) \) gives the main contribution to \( \sigma^S_{xy} \).

To clarify the origins of the above three changes, we analyze \( k \) dependence of \( \sigma^S_{xy}(k) \), defined as \( \sigma^S_{xy} = \frac{1}{\pi} \sum_k \sigma^S_{xy}(k) \), at \( t_{\text{SB}} = 0 \) and 0.09 eV. From the results at \( \gamma_{\text{imp}} = 0.045 \text{ meV} \) shown in Figs. 3(b) and 3(c), we see as \( t_{\text{SB}} \) changes from 0 eV to 0.09 eV, the main contribution to \( \sigma^S_{xy} \) in the clean region changes from the region around \( k \sim (\frac{7}{10}\pi, \frac{3}{10}\pi) \) to the region around \( k \sim (\frac{7}{10}\pi, 0) \). Since the latter main contribution is larger than the former one, change (i) arises from the evolution of the larger positive-sign contribution around \( k \sim (\frac{7}{10}\pi, 0) \). Then, from the results at \( \gamma_{\text{imp}} = 4.5 \text{ meV} \) shown in Figs. 3(d) and 3(e), we see the change of \( t_{\text{SB}} \) from 0 eV to 0.09 eV in the slightly dirty region leads to the sign-change of the main contribution to \( \sigma^S_{xy} \) around \( k \sim (\frac{7}{10}\pi, \frac{3}{10}\pi) \) from positive to negative. Thus, this sign change is the origin of change (ii).

Moreover, combining the results in the clean and the slightly dirty regions, we find change (iii) arises from the competition between the opposite-sign contributions around \( k \sim (\frac{7}{10}\pi, 0) \) and \( k \sim (\frac{3}{10}\pi, \frac{3}{10}\pi) \).

In addition, to understand how each \( t_{\text{SB}} \) orbital contributes to each \( k \) component of \( \sigma^S_{xy}(k) \) at \( t_{\text{SB}} = 0 \) and 0.09 eV, we analyze orbital-decomposed SHCs, obtained by the equations that the summations with respect to orbital and spin indices in Eqs. 1 and 2 are restricted.

Before the results at \( t_{\text{SB}} = 0.09 \text{ eV} \), we explain the relation between the main \( k \) component of \( \sigma^S_{xy}(k) \) and each \( t_{\text{SB}} \) orbital at \( t_{\text{SB}} = 0 \text{ eV} \). Considering all the possible orbital-decomposed SHCs which give the finite contribution to \( \sigma^S_{xy} \) and calculating these values, we find that in the clean and the slightly dirty regions, the main contribution to \( \sigma^S_{xy} \) arises from the term containing \( \{ f^S_x(k) \}_{d_{x^2}-d_{\gamma}} \) and \( \{ f^S_y(k) \}_{d_{\gamma}} \). Thus, all the \( t_{\text{SB}} \) orbitals around \( k \sim (\frac{7}{10}\pi, \frac{3}{10}\pi) \) near the Fermi level plays an important role in the SHE at \( t_{\text{SB}} = 0 \text{ eV} \).

We go on to explain the relation at \( t_{\text{SB}} = 0.09 \text{ eV} \). By using the similar analysis, we find in the clean region, the main contribution to \( \sigma^S_{xy} \) arises from the term containing \( \{ f^S_x(k) \}_{d_{x^2}-d_{\gamma}} \) and \( \{ f^S_y(k) \}_{d_{\gamma}} \). Thus, all the \( t_{\text{SB}} \) orbitals around \( k \sim (\frac{7}{10}\pi, \frac{3}{10}\pi) \) near the Fermi level plays an important role in the SHE at \( t_{\text{SB}} = 0 \text{ eV} \).
the adiabatic process shown in Fig. 4(a) is opposite sign against the main contribution [i.e., Fig. 4(c)], the former is smaller due to $t_{\text{SB}} > t_5$ and $t_{\text{SB}} > \lambda$. Thus, all the $t_{2g}$ orbitals around $k \sim (\frac{\pi}{2}, 0)$ and $(\frac{3\pi}{2}, \frac{\pi}{2})$, affected by the spin splitting due to IS breaking, near the Fermi level becomes important in the SHE at $t_{\text{SB}} = 0.09$ eV. In particular, the sign change of the contribution around $k \sim (\frac{\pi}{2}, \frac{3\pi}{2})$ due to the band anticrossing and its competition with the opposite-sign contribution around $k \sim (\frac{\pi}{2}, 0)$ are vital to obtain the magnitude and sign change of the SHE as a function of $\gamma_{\text{imp}}$.

We now discuss applicability of the similar mechanism for controlling the magnitude and sign of the SHE to other systems. If the following three conditions are satisfied, we can control the magnitude and sign of the SHE in a multiorbital system by introducing IS breaking and tuning $n_{\text{imp}}$. These conditions are, first, that there are several (at least two) same-sign contributions from $\sigma^S_{xy}(k)$ at some momenta in a case without IS in the absence of the band anticrossing; second, that the band anticrossing occurs at one of these momenta due to the combination of orbital degrees of freedom and IS breaking; third, that the value of the band splitting at momentum where the band anticrossing occurs differs from the values of the band splittings giving the other contributions. If the first and second conditions are satisfied, we have the different-sign components of $\sigma^S_{xy}(k)$ at some momenta since the band anticrossing causes the sign change of $\sigma^S_{xy}(k)$ at one of these momenta as a result of the change of the main orbital of the Fermi surface sheet. In addition, if the third condition is satisfied, we can control the magnitude and sign of $\sigma^S_{xy}(k)$ by tuning the value of $\gamma_{\text{imp}}$ through $n_{\text{imp}}$ since an increase of $\gamma_{\text{imp}}$ causes a larger decrease of the contribution of $\sigma^S_{xy}(k)$ arising from $\Delta E_1(k)$ than a decrease of the contribution of $\sigma^S_{xy}(k)$ arising from $\Delta E_2(k) > \Delta E_1(k)$. Since a large value of $t_{\text{SB}}$ is necessary for the band anticrossing and an increase of the potential of the local parity mixing due to IS breaking enhances $t_{\text{SB}}$, these three conditions can be achieved even in other multiorbital systems by introducing IS breaking and tuning $n_{\text{imp}}$. We think that the surface of SrTiO$_3$ [29, 30] and the interface between SrTiO$_3$ and LaAlO$_3$ [31–33] are the good candidates.

Moreover, the similar mechanism is applicable to other Hall effects (e.g., the anomalous Hall effect [14, 34]) since the band anticrossing due to the combination of orbital degrees of freedom and IS breaking drastically affects the main adiabatic processes generating the effective flux.

Thus, the mechanism we found gives a ubiquitous method to control the magnitude and sign of the response of Hall effects in multiorbital systems without IS.

In summary, we have studied the SHE in the $t_{2g}$-orbital tight-binding model of the [001] surface or interface of Sr$_2$RuO$_4$ in the presence of dilute nonmagnetic impurities on the basis of the linear-response theory. We have found when the band anti-crossing occurs at $k = (\frac{\pi}{2}, \frac{\pi}{2})$ due to the cooperative roles of orbital degrees of freedom and IS breaking and its contribution to the SHE becomes dominant compared with the contributions from $\sigma^S_{xy}(k)$ at the other momenta by increasing $\gamma_{\text{imp}}$, the SHE shows the magnitude increase and sign change as a function of $\gamma_{\text{imp}}$, being proportional to $n_{\text{imp}}$. Since the similar situation can be achieved in other systems by tuning the potential of the local parity mixing due to IS breaking and $n_{\text{imp}}$, we propose that the magnitude and sign of the response of Hall effects in some multiorbital systems can be controlled by introducing IS breaking and tuning $n_{\text{imp}}$.

We would like to thank M. Ogata, S. Murakami, and T. Okamoto for fruitful discussions and comments.

\[ \text{mizoguchi@host.phys.s.u-tokyo.ac.jp} \]

[1] M. I. Dyakonov, and V. I. Perel, Sov. Phys. JETP Lett. 13, 467 (1971); M. I. Dyakonov, and V. I. Perel, Phys. Lett. A35, 459 (1971).
[2] J. E. Hirsch, Phys. Rev. Lett. 83, 1834 (1999).
[3] S. A. Wolf, D. D. Awschalom, R. A. Buhrman, J. M. Daughton, S. von Molnár, M. L. Roukes, A. Y. Chichkov, and D. M. Treger, Science 294, 1488 (2001).
[4] S. Murakami, N. Nagaosa, and S. C. Zhang, Science 301, 1348 (2003).
[5] J. Sinova, D. Culcer, Q. Niu, N. A. Sinitsyn, T. Jungwirth, and A. H. MacDonald, Phys. Rev. Lett. 92, 126603 (2004).
[6] S. Murakami, N. Nagaosa, and S. C. Zhang, Phys. Rev. Lett. 93, 156804 (2004).
[7] C. L. Kane, and E. J. Mele, Phys. Rev. Lett. 95, 226801 (2005).
[8] H. Kontani, T. Tanaka, D. S. Hirashima, K. Yamada, and J. Inoue, Phys. Rev. Lett. 100, 096601 (2008).
[9] J. Smit, Physica 24, 39 (1958).
[10] L. Berger, Phys. Rev. B 2, 4559 (1970).
\[\text{[11]}\] A. Crépieux, and P. Bruno, Phys. Rev. B \textbf{64}, 014416 (2001).
\[\text{[12]}\] M. Onoda, and N. Nagaosa, J. Phys. Soc. Jpn. \textbf{71}, 19 (2002).
\[\text{[13]}\] P. Středa, J. Phys. C: Solid State Phys. \textbf{15}, L717 (1982).
\[\text{[14]}\] H. Kontani, T. Tanaka, and K. Yamada, Phys. Rev. B \textbf{75}, 184416 (2007).
\[\text{[15]}\] T. Tanaka, H. Kontani, M. Naito, T. Naito, D. S. Hirashima, K. Yamada, and J. Inoue, Phys. Rev. B \textbf{77}, 165117 (2008).
\[\text{[16]}\] See, e.g., J. J. Sakurai, \textit{Modern Quantum Mechanics} (Benjamin/Cummings, Menlo Park, CA, 1985).
\[\text{[17]}\] E. Saitoh, M. Ueda, H. Miyajima, and G. Tatara, Appl. Phys. Lett. \textbf{88}, 182509 (2006).
\[\text{[18]}\] T. Kimura, Y. Otani, T. Sato, S. Takahashi, and S. Maekawa, Phys. Rev. Lett. \textbf{98}, 156601 (2007).
\[\text{[19]}\] M. Morota, Y. Niimi, K. Ohnishi, D. H. Wei, T. Tanaka, H. Kontani, T. Kimura, and Y. Otani, Phys. Rev. B \textbf{83}, 174405 (2011).
\[\text{[20]}\] Y. K. Kato, R. C. Myers, A. C. Gossard, and D. D. Awschalom, Science \textbf{306}, 1910 (2004).
\[\text{[21]}\] J. Wunderlich, B. Kaestner, J. Sinova, and T. Jungwirth, Phys. Rev. Lett. \textbf{94}, 047204 (2005).
\[\text{[22]}\] E. I. Rashba, Sov. Phys. Solid State \textbf{2}, 1109 (1960).
\[\text{[23]}\] J. Inoue, G. E. W. Bauer, and L. W. Molenkamp, Phys. Rev. B \textbf{67}, 033104 (2003).
\[\text{[24]}\] Y. Yanase, J. Phys. Soc. Jpn. \textbf{82}, 044711 (2013).
\[\text{[25]}\] A. P. Mackenzie, S. R. Julian, A. J. Diver, G. J. McMullan, M. P. Ray, G. G. Lonzarich, Y. Maeno, S. Nishizaki, and T. Fujita, Phys. Rev. Lett. \textbf{76}, 3786 (1996).
\[\text{[26]}\] N. Arakawa, and M. Ogata, Phys. Rev. B \textbf{87}, 195110 (2013).
\[\text{[27]}\] I. I. Mazin, and D. J. Singh, Phys. Rev. Lett. \textbf{79}, 733 (1997).
\[\text{[28]}\] T. Mizoguchi, and N. Arakawa, in preparation.
\[\text{[29]}\] K. Ueno, S. Nakamura, H. Shimotani, A. Ohtomo, N. Kimura, T. Nojima, H. Aoki, Y. Iwasa, and M. Kawasaki, Nat. Mater. \textbf{7}, 855 (2008).
\[\text{[30]}\] P. D. C. King, S. McKeown Walker, A. Tamai, A. de la Torre, T. Eknapakul, P. Buaphet, S.-K. Mo, W. Meevasana, M. S. Bahramy, and F. Baumberger, Nat. Commun. \textbf{5}, 3414 (2014).
\[\text{[31]}\] A. Ohtomo, and H. Y. Hwang, Nature \textbf{427}, 423 (2004).
\[\text{[32]}\] M. Hirayama, T. Miyake, and M. Imada, J. Phys. Soc. Jpn. \textbf{81}, 084708 (2012).
\[\text{[33]}\] Y. Nakamura, and Y. Yanase, J. Phys. Soc. Jpn. \textbf{82}, 083705 (2013).
\[\text{[34]}\] R. Karplus, and J. M. Luttinger, Phys. Rev. \textbf{95}, 1154 (1954); J. M. Luttinger, Phys. Rev. \textbf{112}, 739 (1958).