I. INTRODUCTION

Recent development in engineering of two-dimensional crystalline electron systems has provided a new field in the superconducting research [1-3]. In particular, various phenomena unique to noncentrosymmetric superconductors have been observed in SrTiO$_3$ heterostructures [4] and transition metal dichalcogenides [5-7]. Interplay of antisymmetric spin-orbit coupling (ASOC) and magnetic field has been focused on in these systems. On the other hand, fabrication of artificial superlattices containing strongly correlated electron systems, such as CeCoIn$_5$/YbCoIn$_5$ [8], CeCoIn$_5$/CeRhIn$_5$ [9], and YbRhIn$_5$/CeCoIn$_5$/YbCoIn$_5$ [10], makes interplay of ASOC and strong electronic correlations to be a fascinating topic. For instance, proposals of topological superconductivity [11-16] and enhanced Edelstein effect [17] shed light on potential impact of this topic on topological science and spintronics research.

Bulk CeCoIn$_5$ is one of the typical unconventional superconductors in the vicinity of the antiferromagnetic (AFM) quantum critical point [18]. Non-Fermi liquid behaviors [19-23] and $d_{x^2-y^2}$-wave superconductivity [24], which are characteristic of magnetic criticality, have been established. Artificially-engineered superlattice containing a few layer CeCoIn$_5$ naturally realizes two-dimensional $d_{x^2-y^2}$-wave superconductivity [25]. At the interface of heterostructures Rashba-type ASOC arises from polar inversion symmetry breaking [26], and therefore, the superlattice containing heavy ions is expected to be affected by the Rashba ASOC. Depending on the superlattice structure, staggered or uniform Rashba ASOC appears, and accordingly, locally [27] or globally [10] noncentrosymmetric superconductivity have been supported by experimental results for the heavy fermion superlattices [25]. Unique superconducting phases are expected to be realized there owing to the interplay of two-dimensional magnetic fluctuations and Rashba ASOC.

Noncentrosymmetric structures can also be found in bulk materials. Indeed, vast studies of noncentrosymmetric superconductivity were triggered by the discovery of superconductivity in CePt$_3$Si [29]. Furthermore, a recent experiment uncovered Rashba-type ASOC in a high-temperature cuprate superconductor Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ [30] whose crystal structure can be regarded as a naturally-formed superlattice. Controllability of artificial superlattices as well as spin-momentum locking uncovered in bulk materials generate renewed interest on noncentrosymmetric superconductivity in strongly correlated electron systems.

Motivated by these considerations, we study superconductivity in the two-dimensional Rashba-Hubbard model. Although this model has been analyzed as a minimal model for strongly-correlated electron systems lacking inversion symmetry [14, 31-39], most of theoretical studies are based on weak-coupling approaches such as the perturbation theory or the random phase approximation (RPA) [31-39]. In particular, analysis based on a theoretical method appropriate in quantum critical region has not been conducted. To clarify the superconducting phase stabilized by the interplay of critical magnetic fluctuations and Rashba ASOC, in this paper we adopted fluctuation exchange (FLEX) approximation which appropriately reproduces critical behaviors of self-consistent renormalization theory [41].

An electronic structure characteristic of the Rashba-Hubbard model is spin-splitting due to the Rashba ASOC and resulting type-II van Hove singularity which is positioned away from the time-reversal invariant momenta.
We may expect unusual properties due to a large density of states when the Fermi surface is close to the van Hove singularity. Indeed, a recent theoretical study proposed the ferromagnetic (FM) spin fluctuation and spin-triplet $f$-wave superconductivity [38, 39]. In order to examine this proposal and to provide a thorough study of unconventional superconductivity in the Rashba-Hubbard model, we calculate the Fermi surfaces (FSs), magnetic susceptibility, and superconducting gap functions in a wide range of the filling. We show that FSs are robust against critical magnetic fluctuations in contrast to previous theory [42]. Furthermore, we show that strongly parity-mixed superconductivity with dominant $d_{x^2-y^2}$-wave pairing is robust in a whole parameter range in contrast to the proposal in Ref. [38]. Interestingly, the parity mixing is enhanced near the van Hove singularity. Indeed, a recent theoretical study proposed $\alpha$-wave pairing is attributed not to the FM fluctuation but ferromagnetic (CAFM) order, and the spin-triplet gap is included in $\varepsilon(k)$. The second term in the free part of Hamiltonian, Eq. (2), describes the ASOC which appears in crystals lacking inversion symmetry. The g-vector, $g(k)$, characterizes the structure of ASOC [29], and it is Rashba type in polar noncentrosymmetric systems. We here assume a Rashba type g-vector represented as $g(k) = \left( -\frac{\partial \varepsilon(k)}{\partial k_y}, \frac{\partial \varepsilon(k)}{\partial k_x}, 0 \right)$. The ASOC shows a form of the momentum-dependent Zeeman field. Therefore, spin degeneracy of the band is split by the ASOC, and bands with negative and positive helicity have distinct energy,

$$E_\lambda(k) = \varepsilon(k) + \lambda \alpha|g(k)|,$$

where $\lambda = \pm$ is the helicity index. Unless stated otherwise, we set a temperature $T = 0.01$, $t' = 0.3$, and $\alpha = 0.5$ with a unit of energy $t = 1$.

**B. Green function and susceptibility**

The noninteracting Green functions for $U = 0$ are expressed by the $2 \times 2$ matrix form in the spin basis,

$$\hat{G}^{(0)}(k, i\omega_n) = \left( i\omega_n I - \varepsilon(k) I - \alpha g(k) \cdot \sigma \right)^{-1},$$

where $\omega_n = (2n + 1)\pi T$ are fermionic Matsubara frequencies. The noninteracting Green functions in the helicity basis,

$$G_{\lambda}^{(0)}(k, i\omega_n) = \frac{1}{i\omega_n - \varepsilon(k) - \lambda \alpha |g(k)|},$$

are obtained by unitary transformation with $\hat{V}$ which diagonalizes $H_0$ [Eq. (7)],

$$\hat{V}^\dagger H_0 \hat{V} = \begin{pmatrix} E_+(k) & 0 \\ 0 & E_-(k) \end{pmatrix}.$$
Within the FLEX approximation, the self-energy is expressed as

\[ \hat{\Sigma}(\mathbf{k}, i\omega_n) = \sum_{\lambda = \pm} \frac{1}{2} \left( \hat{I} + \lambda \frac{\mathbf{g}}{g^*} \cdot \mathbf{\sigma} \right) G^{(0)}_\lambda(\mathbf{k}, i\omega_n). \]  

In the interacting case \( U \neq 0 \), the dressed Green functions contain a self-energy, \( \hat{\Sigma}(\mathbf{k}, i\omega_n) \),

\[ \hat{G}(\mathbf{k}, i\omega_n) = \left( i\omega_n \hat{I} - \varepsilon(\mathbf{k}) \hat{I} - \alpha \mathbf{g}(\mathbf{k}) \cdot \mathbf{\sigma} - \hat{\Sigma}(\mathbf{k}, i\omega_n) \right)^{-1}. \]

Within the FLEX approximation, the self-energy is expressed with use of an effective interaction, \( \hat{\Gamma}^{\text{n}}(\mathbf{k}, i\nu_n) \), as

\[ \Sigma_{\sigma\sigma'}(\mathbf{k}, i\omega_n) = T \sum_{\mathbf{q}, i\nu_n} \Gamma^{\text{n}}_{\sigma\sigma'}(\mathbf{q}, i\nu_n) G_{\xi\eta}(\mathbf{k} - \mathbf{q}, i\omega_n - i\nu_n), \]

and the effective interaction is given by

\[ \hat{\Gamma}^{\text{n}}(\mathbf{k}, i\nu_n) = \hat{U} \hat{\chi}(\mathbf{k}, i\nu_n) \hat{U}, \]

where

\[ \hat{U} = \begin{pmatrix} 0 & 0 & 0 & -U \\ 0 & U & 0 & 0 \\ 0 & 0 & U & 0 \\ -U & 0 & 0 & 0 \end{pmatrix}, \]

\( \hat{\chi}(\mathbf{k}, i\nu_n) \) is the generalized susceptibility, and \( i\nu_n \) are bosonic Matsubara frequencies. We introduce the bare susceptibility

\[ \chi^{(0)}_{\sigma_1\sigma_2\sigma_3\sigma_4}(\mathbf{q}, i\nu_n) = -T \sum_{\mathbf{k}, i\omega_n} G_{\sigma_1\sigma_3}(\mathbf{k}, i\omega_n) G_{\sigma_4\sigma_2}(\mathbf{k} - \mathbf{q}, i\omega_n - i\nu_n), \]

and compute the generalized susceptibility by

\[ \hat{\chi}(\mathbf{q}, i\nu_n) = \left( \hat{I} - \chi^{(0)}(\mathbf{q}, i\nu_n) \hat{U} \right)^{-1} \chi^{(0)}(\mathbf{q}, i\nu_n). \]

According to Eqs. (12)-(17), \( \hat{G}, \hat{\Sigma}, \hat{\Gamma}, \hat{\chi}^{(0)}, \hat{\chi} \) depend on each other, and therefore, we self-consistently determine these functions. As a consequence of the self-consistent condition, the FLEX approximation is a conserving approximation in which several conservation laws are satisfied in the framework of the Luttinger-Ward theory.

Introducing the vector representation of the self-energy

\[ \hat{\Sigma} = \Sigma_0 \hat{I} + \mathbf{\Sigma} \cdot \mathbf{\sigma}, \]

and carrying out analytic continuation, we represent the renormalized retarded Green functions as

\[ \hat{G}^{\text{R}}(\mathbf{k}, \omega) = \left( \omega \hat{I} - \varepsilon'(\mathbf{k}, \omega) \hat{I} - \alpha \mathbf{g}'(\mathbf{k}, \omega) \cdot \mathbf{\sigma} \right)^{-1}, \]

where \( \varepsilon' = \varepsilon + \Sigma^R_0 \) and \( \alpha \mathbf{g}' = \alpha \mathbf{g} + \text{Re} \Sigma^R \). Since \( \text{Im} \Sigma^R \) is proportional to \( T^2 \) in a Fermi liquid state, we dropped it and obtain

\[ \hat{G}^R(\mathbf{k}, \omega) = \sum_{\lambda = \pm} \frac{1}{2} \left( \omega \hat{I} + \lambda \frac{\mathbf{g}'}{g'} \cdot \mathbf{\sigma} \right) G^R_\lambda(\mathbf{k}, \omega), \]

where

\[ G^R_\lambda(\mathbf{k}, \omega) = \frac{1}{\omega - \varepsilon'(\mathbf{k}, \omega) - \lambda \alpha \alpha |\mathbf{g}'(\mathbf{k}, \omega)|}. \]

From Eq. (21), we determine FSs of interacting systems by solving

\[ \varepsilon'(\mathbf{k}, 0) - \lambda \alpha |\mathbf{g}'(\mathbf{k}, 0)| = 0. \]

In this calculation, a static function \( A(\mathbf{q}, 0) \) is evaluated by an approximation justified at low temperatures,

\[ A(\mathbf{q}, 0) \sim \frac{A(\mathbf{q}, i\pi T) + A(\mathbf{q}, -i\pi T)}{2}. \]

C. Linearized \'Eliashberg equation

To investigate superconductivity, we numerically solve the linearized \'Eliashberg equation which is given by

\[ \lambda \Delta_{\sigma\sigma'}(\mathbf{k}) = T \sum_{\mathbf{k}' \not\sigma' \Sigma} \Gamma_{\sigma\sigma'}(\mathbf{k} - \mathbf{k}') F_{\sigma_1 \sigma_2}(\mathbf{k}'), \]

\[ F_{\sigma_1 \sigma_2}(\mathbf{k}') = G_{\sigma_1 \sigma_3}(\mathbf{k}') \Delta_{\sigma_3 \sigma_4}(\mathbf{k}') G_{\sigma_4 \sigma_2}(-\mathbf{k}'), \]

where \( \Delta \) is the gap function and \( \hat{\Gamma} \) is obtained by

\[ \hat{\Gamma}(\mathbf{k} - \mathbf{k}') = \hat{U} + \hat{\Gamma}^{\text{n}}(\mathbf{k} - \mathbf{k}'). \]

Here we adopted abbreviated notation \( \mathbf{k} = (\mathbf{k}, i\omega_n) \).

Evaluating \( \lambda \), eigenvalues of the linearized \'Eliashberg equation, we determine the critical temperature \( T_c \) from the criterion \( \lambda = 1 \).

Even when \( \lambda \neq 1 \), we can identify the leading superconducting instability by comparing \( \lambda \) for irreducible representations of a given point group. Since the point group symmetry of the Rashba-Hubbard model is \( C_{4v} \), the gap function can be classified into irreducible representations of \( C_{4v} \). We numerically calculate eigenvalues for each irreducible representations and conclude that the \( B_1 \) representation gives the largest eigenvalue \( \lambda \) in the whole parameter range.

III. FERMI SURFACES

As is known by vast previous works, the topology and shape of FSs are closely related to magnetic fluctuations and superconductivity. Therefore, we begin with discussions about the FSs of interacting systems. In Fig. 1(a)-(d), we compare the FSs in the noninteracting
IV. MAGNETIC FLUCTUATIONS

Next, we discuss magnetic fluctuations. Dynamical spin susceptibility tensor is given by the generalized susceptibility as

\[
\chi^{\mu\nu}(q, i\nu_n) = \sum_{s_1s_2s_3s_4} \sigma_{s_1s_2} \sigma_{s_3s_4}(q, i\nu_n) \alpha_{s_3s_4} \quad (27)
\]

We illustrate static longitudinal spin susceptibility, \(\chi^{zz}(Q, 0)\), at \(n = 0.65\) and \(n = 0.85\) for various fillings. Although \(\chi^{zz}\) and \(\chi^{+}\) are equivalent in the absence of the spin-orbit coupling, magnetic anisotropy is induced by the Rashba ASOC in this model.

The transverse spin susceptibility shows qualitatively similar momentum dependence for all fillings [Figs. 2(a)-[h]]. The incommensurate antiferromagnetic (IAFM) fluctuation appears in a wide range of filling, \(0.65 \lesssim n \lesssim 0.95\). A weak CAFM fluctuation with the modulation vector \(Q = (\pi, \pi)\) also develops at \(n = 0.65\) near the type-II van Hove singularity.

On the other hand, we observe a significant enhancement of the CAFM fluctuation in the longitudinal spin susceptibility at \(n = 0.65\) [Fig. 2(a)], whereas the IAFM fluctuation is dominant for \(0.75 \lesssim n \lesssim 0.95\) [Figs. 2(b)-[d]]. The longitudinal spin susceptibility not only reveals the AFM fluctuations but also implies the FM spin fluctuation when the Fermi level is close to the type-I or type-II van Hove singularity [39], as we see a weak peak at \(Q = (0, 0)\) [Figs. 2(b) and 2(c)] as well as Fig. 5(a). However, consistent with the previous analysis based on the RPA [39], the FM fluctuation is weakened in the strong coupling region. Indeed, we see only a weak signature of the FM spin fluctuation. As we show later, this FM spin fluctuation is almost unrelated to the superconductivity.

For all fillings in Fig. 2, the longitudinal spin correlation is stronger than the transverse one. Thus, the Ising-type AFM spin fluctuation with dominant \(c\)-axis component is implied. The magnetic anisotropy is enhanced when the filling \(n\) is decreased and the Fermi level approaches to the type-II van Hove singularity.

Growth in the maximum value of the longitudinal spin susceptibility, \(\chi^{zz}(Q, 0)\), at \(n = 0.65\) suggests that the system is in the vicinity of the CAFM order. It should be noticed that in Fig. 2 we choose \(U = 2.4\) for \(n = 0.65\) while \(U = 5\) for other fillings. This is because \(U = 5\) is larger than the critical interaction for the AFM order at \(n = 0.65\). In fact, the critical interaction is approximately \(U_c = 3.3\) for \(n = 0.65\), whereas \(U_c > 6\) for other fillings. Thus, it is indicated that the AFM order develops when the FSs are close to the type-II van Hove singularity. Such filling dependence of magnetic fluctuations is qualitatively different from the conventional Hubbard.
model without the Rashba ASOC. In the ordinary Hubbard model, the magnetic correlations are enhanced near the half-filling. On the other hand, the magnetic correlations are mainly determined by the type-II van Hove singularity in the Rashba-Hubbard model with a large ASOC.

V. SUPERCONDUCTIVITY

Here we study the superconductivity. Superconducting phases are classified based on irreducible representations of the point group symmetry of the system, that is, \( C_{4v} \). We calculate eigenvalues of the linearized E\'liashberg equation for all the irreducible representations, \( A_1, A_2, B_1, B_2, \) and \( E \). For instance, Fig. 3 shows the interaction dependence of \( \lambda \) and reveals that the superconducting phase of \( B_1 \) representation is the most stable. We confirmed that the \( B_1 \) superconducting phase is stable in the whole filling range investigated in this paper, that is, from half-filling to the type-II van Hove singularity.

Fig. 3 shows the filling dependence of the maximum eigenvalue \( \lambda \) for the \( B_1 \) representation. The results suggest superconductivity with a high transition temperature near the half-filling, whereas the transition temperature decreases in the low-filling region. It should be noticed that the magnetic fluctuation grows near the type-II van Hove singularity more strongly than near the half-filling (Fig. 2). Our results indicate a weak tendency to superconductivity near the type-II van Hove singularity in spite of a strong instability to the CAFM order.

This is partly because the magnetic fluctuation is significantly localized in the momentum space: \( \chi^{zz}(q,0) \) shows a sharp peak around the commensurate wave vector \( q = Q \). It makes total weight of the spin fluctuation, \( \int dq \chi^{zz}(q,0) \), to be small. A strong magnetic anisotropy also favors magnetic order rather than superconductivity. Because both longitudinal and transverse spin fluctuations mediate an attractive interaction for spin-singlet pairing \[49\], an isotropic spin fluctuation may give rise to higher superconducting transition temperatures than the Ising spin fluctuation.

The superconducting order parameter of \( B_1 \) representation contains spin-singlet \( d_{x^2-y^2} \)-wave pairing as well as spin-triplet pairing with either \( p \)-wave or \( f \)-wave symmetry. Because of the Rashba ASOC, superconducting order parameters with distinct space inversion parity coexist. The gap functions are decomposed into the spin-singlet component \( \psi(k) \) and spin-triplet component \( d(k) \) in a standard manner,

\[
\tilde{\Delta}(k) = (\psi(k) + d(k) \cdot \sigma) i\sigma_y. \quad (28)
\]

Fig. 3 shows gap functions of the most stable \( B_1 \) state for various fillings. In the whole parameter range, a strongly parity-mixed superconducting state is stabilized. Although the \( d_{x^2-y^2} \)-wave pairing is always dominant, the subdominant spin-triplet pairing component changes the momentum dependence as a function of the filling. The \( d_{x^2-y^2} + f \)-wave state is stabilized for \( n = 0.65 \), whereas the \( d_{x^2-y^2} + p \)-wave state is stable for other fillings. As we have shown in Sec. III, the longitudinal spin susceptibilities show qualitatively different behaviors between \( n = 0.65 \) and other fillings. The correspondence between Figs. 2 and 4 implies that the CAFM fluctuation favors the \( d_{x^2-y^2} + f \)-wave pairing whereas the IAFM fluctuation favors the \( d_{x^2-y^2} + p \)-wave pairing. This is consistent with the previous RPA analysis where the CAFM fluctuation arising from the strong nesting \((t' = 0 \text{ and } n \sim 1)\) stabilizes a \( d_{x^2-y^2} + f \)-wave state \[35\].

Here we compare our results with a previous work which investigated similar parameter range within the RPA \[35\]. The authors of Ref. \[35\] claimed that various superconducting states with different symmetry are stabilized depending on the filling. In particular, the spin-triplet \( f \)-wave pairing state near the type-II van Hove singularity has been illustrated, and its origin was attributed to the FM spin fluctuation. In contrast to their results, our calculation based on the FLEX approximation with the linearized E\'liashberg equation shows that the gap functions of superconductivity are essentially independent of the filling and the \( d_{x^2-y^2} \)-wave pairing is
of the magnitudes which is evaluated by
\[ r = \frac{\sum_k |\psi(k)|^2}{\sum_k |d(k)|^2}. \tag{29} \]

We see \( r > 1 \) in the whole parameter range, indicating the dominant \( d_{x^2-y^2} \)-wave pairing. However, the value of \( r \) is close to unity, and therefore, strongly parity-mixed superconducting states are concluded. The parity mixing is particularly enhanced around \( n = 0.75 \), in which the FSs lie between the type-I and type-II van Hove singularities.

VI. CONCLUSIONS AND DISCUSSIONS

In summary, we have conducted a thorough study on superconductivity in the two-dimensional Rashba-Hubbard model, a minimal model for strongly-correlated noncentrosymmetric electron systems. With use of the FLEX approximation combined with linearized Eliashberg equation, we have clarified interplay of Rashba spin-orbit coupling and critical magnetic fluctuations in a wide range of filling from type-II van Hove singularity to half-filling. Our results reveal robust FSs against the critical magnetic fluctuation, enhancement of IAFM fluctuation, and stabilization of strongly parity-mixed superconducting state in a wide parameter range. The obtained gap functions show the \( d_{x^2-y^2} + p \)-wave superconductivity for the filling \( n = 0.75, 0.85, \) and 0.95. On the other hand, for \( n = 0.65 \), we have observed impacts of the type-II van Hove singularity near the Fermi level. The FSs undergo Lifshitz transition due to the electron correlation, the CAFM fluctuation is strongly enhanced, and the \( d_{x^2-y^2} + f \)-wave superconducting state is stabilized. Strong parity mixing in the gap functions has been observed in the whole parameter range. In particular, magnitude of spin-triplet gap function is comparable to that of spin-singlet one when the Fermi level lies between the type-I and type-II van Hove singularities.
This work resolved unsettled issues on the strongly correlated Rashba systems [37, 39, 12], and elucidated a microscopic mechanism to stabilize a strongly parity-mixed superconducting phase. This finding opens a way to realize intriguing phenomena arising from parity mixing in superconducting order parameters. For instance, as proposed by recent theoretical studies, noncentrosymmetric superconductors with strong parity mixing may be a platform of fractional flux quanta [44], nonreciprocal electric current [43], and topological superconductivity [11–16].

Another class of superconducting phases with strong parity mixing may be stabilized by a critical fluctuation of structural transitions breaking the space inversion symmetry, that is named odd-parity electric multipole fluctuations [51, 52]. In this case, phonons coupled to dynamical spin-orbit coupling mediate pairing interaction in both spin-singlet and spin-triplet channels. In contrast, in our proposal anisotropic magnetic fluctuations naturally lead to a strongly parity-mixed superconducting state in quasi-two-dimensional electron systems with strong electron correlations. The candidates may be naturally-formed or artificially-engineered heterostructures of cuprate superconductors [30, 53] or heavy fermion superconductors [8–10, 25, 27, 28].

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