Spin polarons in the t-J model: Shape and backflow

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The self-consistent Born approximation for the self-energy and the wave function of a single hole moving in a two-dimensional quantum antiferromagnet is used to calculate various spin-correlation functions that characterize the perturbation induced by the hole. The results for these correlation functions agree with the available small-cluster calculations. The hole-spin correlation functions decay asymptotically with the distance from the hole as a power law and exhibit a dipolar distortion of the antiferromagnet. The quasiparticle spectral weight remains finite.

The two-dimensional (2D) t-J model is considered to provide a relevant description of the low-energy physics of high-T_c superconductors, where the moving charge carriers strongly interact with the Cu d^9 spins which form a 2D spin-$\frac{1}{2}$ quantum antiferromagnet (AFM).[1] In particular, the problem of a single mobile vacancy (hole) in an AFM plane was extensively investigated using various analytical approaches as well as with exact diagonalization of small clusters.[2–7] Consensus has been achieved regarding, e.g., the ground-state energy and the quasiparticle (QP) bandwidth. The distortion of the AFM background around the vacancy and the corresponding distribution of spin was only partially studied quantitatively.[3,8,9] Despite these studies there are still different opinions whether the QP spectral weight is finite or vanishes in the thermodynamic limit.[10] The issue of debate is whether a finite spectral weight as expected from the scaling of exact diagonalization results[11] or from approaches relying on the Born approximation[7] is consistent with a power-law decay of certain correlation functions which describe the deformation around the hole.

The aim of this work is to present a more detailed picture of the shape and size of the quasiparticle. We investigate various types of correlation functions which characterize the perturbation of the spin background induced by the hole. Among these are the change of staggered magnetization, the change of nearest-neighbor spin correlations, and the spin-backflow current associated with the moving hole. The explicit form of the quasiparticle wave function allows us to extract the exponents which determine the power-law decay of these correlation functions as well as their symmetry.

Our approach is based on a spinless-fermion–Schwinger-boson representation for the t-J Hamiltonian. By means of linear-spin-wave (LSW) theory we introduce the proper collective excitations of the Heisenberg antiferromagnet. After these steps the t-J Hamiltonian[5–7]

$$H=\frac{1}{\sqrt{N}}\sum_{\mathbf{kq}}(M_{\mathbf{kq}}h^\dagger_{\mathbf{k}-\mathbf{q}}h_{\mathbf{k}}\alpha^\dagger_{\mathbf{q}}+\text{H.c.})+\sum_{\mathbf{q}}\omega_{\mathbf{q}}\alpha^\dagger_{\mathbf{q}}\alpha_{\mathbf{q}} \qquad (1)$$

has many similarities with the Hamiltonian known from the polaron problem. Here $h^\dagger_{\mathbf{k}}$ is the creation operator for a (spinless) hole in a Bloch state and $\alpha^\dagger_{\mathbf{q}}=u_{\mathbf{q}}a^\dagger_{\mathbf{q}}-v_{\mathbf{q}}a_{\mathbf{q}}$ creates an antiferromagnetic magnon with energy $\omega_{\mathbf{q}}$. The fermion-magnon coupling is given by $M_{\mathbf{kq}}=4t(u_{\mathbf{q}}\gamma_{\mathbf{k}-\mathbf{q}}+v_{\mathbf{q}}\gamma_{\mathbf{k}})$, where $\gamma_{\mathbf{q}}=(\cos q_x+\cos q_y)/2$, $u_{\mathbf{q}}$ and $v_{\mathbf{q}}$ are the usual Bogoliubov factors in linear spin wave theory, and N is the number of lattice sites.

We will calculate the Green's function for the hole, $G_{\mathbf{k}}(\omega)=[\omega-\Sigma_{\mathbf{k}}(\omega)]^{-1}$, in self-consistent Born approximation (SCBA), i.e., $\Sigma_{\mathbf{k}}(\omega)=\frac{1}{N}\sum_{\mathbf{q}}M^2_{\mathbf{kq}}G_{\mathbf{k}-\mathbf{q}}(\omega-\omega_{\mathbf{q}})$. This approximation amounts to the summation of noncrossing diagrams to all orders. The QP dispersion and spectral weight calculated within the SCBA (Ref. 7) agree very well with the exact diagonalization results for small clusters.[11] The success of the SCBA has roots in the vanishing of low order vertex corrections as pointed out by several authors.[6,7,12] Recently Bonfim and Reiter[13] derived the quasiparticle wave function corresponding to the SCBA,

$$\begin{aligned}|\Psi^{(n)}_{\mathbf{k}}\rangle = a_{\mathbf{k}}\Big[h^\dagger_{\mathbf{k}} &+ N^{-1/2}\sum_{\mathbf{q}_1}M_{\mathbf{kq}_1}G_{\bar{\mathbf{k}}_1}(\bar\omega_1)h^\dagger_{\bar{\mathbf{k}}_1}\alpha^\dagger_{\mathbf{q}_1}\\ &\cdots + N^{-n/2}\sum_{\mathbf{q}_1,\ldots,\mathbf{q}_n}M_{\mathbf{kq}_1}G_{\bar{\mathbf{k}}_1}(\bar\omega_1)\cdots M_{\bar{\mathbf{k}}_{n-1}\mathbf{q}_n}\\ &\times G_{\bar{\mathbf{k}}_n}(\bar\omega_n)\, h^\dagger_{\bar{\mathbf{k}}_n}\alpha^\dagger_{\mathbf{q}_1}\ldots\alpha^\dagger_{\mathbf{q}_n}\Big]|0\rangle. \end{aligned} \qquad (2)$$

Here $\bar{\mathbf{k}}_m=\mathbf{k}-\mathbf{q}_1\cdots-\mathbf{q}_m$, $\bar\omega_m=\epsilon_{\mathbf{k}}-\omega_{\mathbf{q}_1}\cdots-\omega_{\mathbf{q}_m}$, and $|0\rangle$ is the vacuum for fermion and magnon operators. In the limit $n\to\infty$, $|\Psi^{(n)}_{\mathbf{k}}\rangle$ is a solution of the Schrödinger equation $H|\Psi^{(\infty)}_{\mathbf{k}}\rangle=\epsilon_{\mathbf{k}}|\Psi^{(\infty)}_{\mathbf{k}}\rangle$ where $\epsilon_{\mathbf{k}}=\Sigma_{\mathbf{k}}(\epsilon_{\mathbf{k}})$ is the quasiparticle energy and $\Sigma_{\mathbf{k}}$ is the self-energy in the SCBA.[13] The terms which are omitted in Eq. (2) correspond to the crossing diagrams also neglected in the Green's function. Because of AFM long-range order the quasiparticle moves on one sublattice visiting the other sublattice only virtually, therefore $\epsilon_{\mathbf{k}}=\epsilon_{\mathbf{k}+(\pi,\pi)}$. For the calculation of correlation functions this implies that one should use proper linear combinations of the wave function (2), i.e., $|\Psi^{(n)}_{\mathbf{k}\pm}\rangle\equiv 2^{-1/2}[|\Psi^{(n)}_{\mathbf{k}}\rangle\pm|\Psi^{(n)}_{\mathbf{k}+(\pi,\pi)}\rangle]$, corresponding to the QP on $\uparrow$ and $\downarrow$ sublattice, respectively.

The QP spectral weight is given by $(a_{\mathbf{k}})^2=[1-\partial\Sigma_{\mathbf{k}}(\omega)/\partial\omega|_{\epsilon_{\mathbf{k}}}]^{-1}$. It can be shown[14] that this also guarantees the proper normalization of the wave function Eq. (2), i.e., $\langle\Psi^{(\infty)}_{\mathbf{k}\pm}|\Psi^{(\infty)}_{\mathbf{k}\pm}\rangle=1$. This establishes the full internal consistency of the two schemes for the Green's function and for the wave function.

In Fig. 1 we present $\mathcal{N}_{\mathbf{k}} = \langle \Psi^{(n)}_{\mathbf{k}\pm} | \Psi^{(n)}_{\mathbf{k}\pm} \rangle$ as a function of J/t for various numbers of magnons $n \leq 3$ kept in the wave function at the minimum of the QP band, i.e., $\mathbf{k} = (\frac{\pi}{2}, \frac{\pi}{2})$. The spectral weight $(a_{\mathbf{k}})^2$ is given by $\mathcal{N}_{\mathbf{k}}$ for $n = 0$. The numerical study was performed using a $N = 16 \times 16$ unit cell. Convergency was checked by calculations up to $N = 32 \times 32$. The deviation of $\mathcal{N}_{\mathbf{k}}$ from unity gives an estimate of the importance of the higher order magnon terms not included in $|\Psi^{(n)}_{\mathbf{k}\pm}\rangle$. We found that for $J/t \geq 0.3$ it was sufficient to include three-magnon terms in Eq. (2), and in the following study we used $n = 3$. (In the Ising case using analytical methods we were able to calculate up to $n = 40$.)

Figure 2(a) shows the distribution of bosons around the hole, $N_{\mathbf{R}} = \langle n_0 (a^\dagger_{\mathbf{R}} a_{\mathbf{R}}) \rangle$, calculated in the Ising limit for $J_z/t = 0.002$ and $k = 0$. Here $n_i = h^\dagger_i h_i$ is the hole density operator at $\mathbf{R}_i = \mathbf{0}$. Within the LSW approximation $\langle n_0 (a^\dagger_{\mathbf{R}} a_{\mathbf{R}}) \rangle$ is equivalent to $\langle n_0 \, (S^+_{\mathbf{R}} S^-_{\mathbf{R}}) \rangle$. The correlation function $N_{\mathbf{R}}$ can therefore be used as a suitable definition of the *spin polaron* and consequently of its spatial size. The total number of bosons for $J/t < 0.2$ is close to the result obtained in Ref. 15, with $n_{\text{tot}} \equiv \sum_{\mathbf{R}} N_{\mathbf{R}} \sim 1.4 (t/J)^{1/3}$. For the spatial dependence we found $N_{\mathbf{R}} \propto \exp[-(R/R_s)^{3/2}]$. The size of the polaron can be characterized by a radius R_p which encloses a given fraction $p = \sum_{R<R_p} N_{\mathbf{R}}/n_{\text{tot}}$ of the polaron. For $p \gtrsim 0.5$ and $J/t \geq 0.002$ we obtained the scaling $R_p \propto (t/J)^{1/6}$ with $R_{0.9} \sim 2$ for $J/t = 0.2$. The total number of bosons and the scaling of the polaron size as well as its asymptotical R dependence are consistent with the picture of the random motion of the hole confined by a string potential. There the average path length[4] scales as $l \propto n_{\text{tot}} \propto J^{-1/3}$ and for the random walk of the hole follows $R_p \propto l^{1/2}$. In the SCBA the spin polarons of the t-J_z model do not depend on $\mathbf{k}$ and have the symmetry of the square lattice.

This is different in the t-J model. In Fig. 2(b) $N_{\mathbf{R}} = \langle n_0 (a^\dagger_{\mathbf{R}} a_{\mathbf{R}}) \rangle - N_{\text{AFM}}$ is shown for $J/t = 0.4$ and $\mathbf{k} = (\frac{\pi}{2}, \frac{\pi}{2})$. Here we have subtracted the large contribution $N_{\text{AFM}} = 0.197$ from quantum fluctuations in the ground state in the absence of the hole. The shape of the polaron is extended in the direction of the QP momentum which reflects a quasi-one-dimensional motion of the polaron. This is consistent with the asymmetry of the QP energy band in the "hole pocket," where the effective next-nearest-neighbor hopping for the $(1,1)$ direction is $\sim 5\times$ that of the $(1,-1)$ direction.[6,7] The asymmetry is most pronounced at $\mathbf{k} = (\frac{\pi}{2}, \frac{\pi}{2})$, and gradually vanishes away from the QP energy minimum and disappears at $k = 0$ and $\mathbf{k} = (\pi, 0)$. Since the magnetic excitations $\omega_{\mathbf{q}}$ vanish linearly with q, $N_{\mathbf{R}}$ decays as a power law, as predicted by Schmitt-Rink *et al.*[5] Below we will show that $N_{\mathbf{R}} \sim R^{-2}$ which implies that the definition of the polaron size used for the t-J_z model cannot be used here.

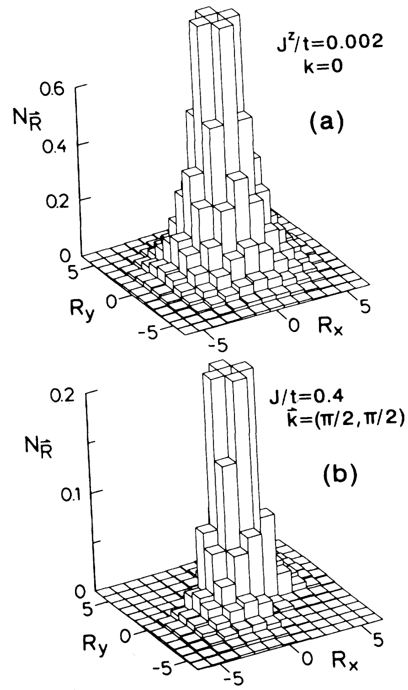

FIG. 2. Distribution of spin deviations around the hole, $N_{\mathbf{R}}$. (a) Ising limit of the model, $J_z/t = 0.002$, $k = 0$ and $N \to \infty$, $n = 40$. (b) Isotropic model, $J/t = 0.4$, $\mathbf{k} = (\frac{\pi}{2}, \frac{\pi}{2})$ and $N = 16 \times 16$, $n = 3$.

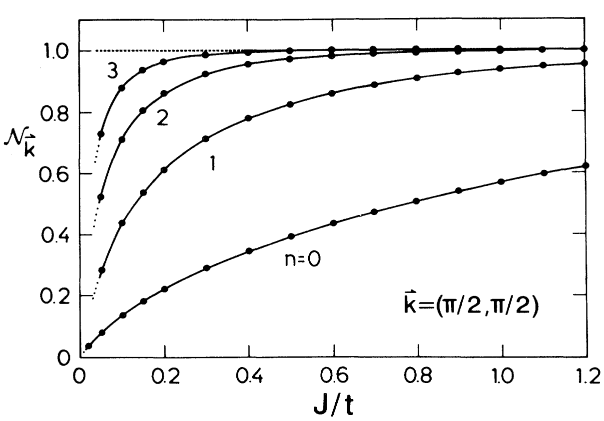

FIG. 1. The norm $\mathcal{N}_{\mathbf{k}} = \langle \Psi^{(n)}_{\mathbf{k}\pm} | \Psi^{(n)}_{\mathbf{k}\pm} \rangle$ as a function of J/t for $n = 0, ..., 3$ at $\mathbf{k} = (\frac{\pi}{2}, \frac{\pi}{2})$ and the unit cell $N = 16 \times 16$.

In Fig. 3 we display the strong $\mathbf{k}$ dependence of three different correlation functions which measure different aspects of the perturbation introduced by the hole: (a) The local z component of spin $S_{\mathbf{R}} = \langle n_0 S^z_{\mathbf{R}} \rangle$. To calculate this, $S^z_{\mathbf{R}_j} = (-1)^{(R_x+R_y)_j} (\frac{1}{2} - a^\dagger_{\mathbf{R}_j} a_{\mathbf{R}_j})$ must be first expressed in terms of magnon operators $\alpha_{\mathbf{q}}$. This also shows that the x and y components of spin vanish because $\langle \Psi^{(n)}_{\mathbf{k}\pm} | n_i \alpha_{\mathbf{q}} | \Psi^{(n)}_{\mathbf{k}\pm} \rangle \equiv 0$. (b) The effect of the hole on the AF correlations and the energy of the spin system is measured with the correlation function $C_{\mathbf{R}} = \langle n_0 \, (\mathbf{S}_{\mathbf{R}_1} \cdot \mathbf{S}_{\mathbf{R}_2}) \rangle$ defined on bonds between two neighboring sites (1–2), $\mathbf{R} = (\mathbf{R}_1 + \mathbf{R}_2)/2$. (c) Another interesting aspect of the deformation of the spin background is contained in the bond-spin currents $\mathbf{j}_{\mathbf{R}} = \langle n_0 (\mathbf{S}_{\mathbf{R}_1} \times \mathbf{S}_{\mathbf{R}_2})^z \mathbf{\Delta} \rangle$, where $\mathbf{\Delta}$ is a unit vector $\mathbf{\Delta} = \mathbf{R}_2 - \mathbf{R}_1$. This quantity follows from the equation of motion for the spin density,

$$\dot{\mathbf{S}}_{\mathbf{R}} = it \sum_{\Delta, ss'} (\hat{\sigma}_{ss'} c^\dagger_{\mathbf{R},s} c_{\mathbf{R}+\Delta,s'} - \text{H.c.}) - 2iJ \sum_{\Delta} \mathbf{S}_{\mathbf{R}} \times \mathbf{S}_{\mathbf{R}+\Delta},$$

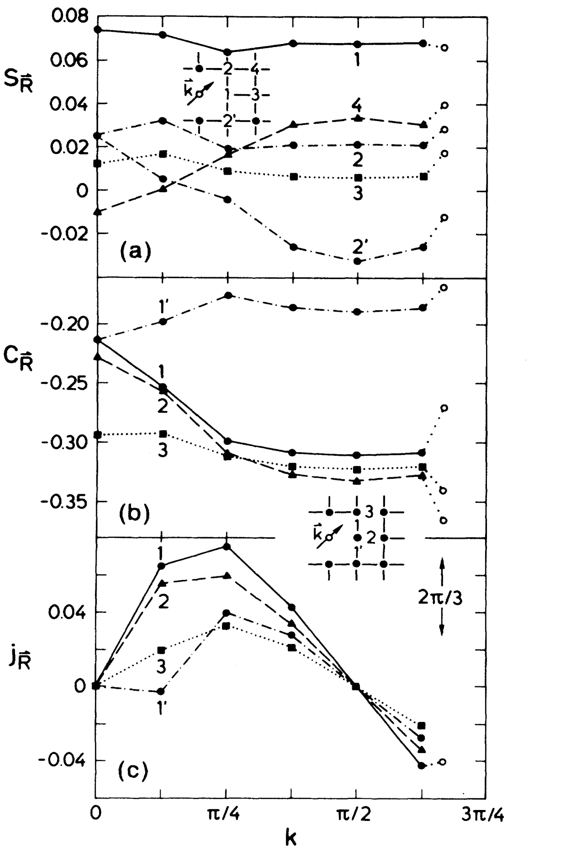

FIG. 3. Momentum dependence of various correlation functions along $\mathbf{k} = (k,k)$ for $J/t = 0.4$: (a) $S_{\mathbf{R}}$, (b) $C_{\mathbf{R}}$, (c) $j_{\mathbf{R}}$. Always $N = 16 \times 16$ and $n = 3$ was used. The labeling of the bond vectors $\mathbf{R}$ for (c) is the same as for (b). The lines connecting the symbols are a guide to the eye only. The open circles ○ give a comparison with exact diagonalization results for an 18 site cluster (Ref. 8).

where $\hat{\sigma}$ are Pauli spin matrices. Here the first term is the spin current induced by the hopping of the hole and the second term ($\sim \mathbf{j}_{\mathbf{R}}$) describes the backflow in the spin system. Due to the broken symmetry total spin is not a good quantum number; therefore we consider only the z component of the current. As total S^z is conserved, the space integral over the two terms cancels.

In Fig. 3 all three correlation functions are presented for $J/t = 0.4$ along $\mathbf{k} = (k,k)$ for different positions $\mathbf{R}$. The distribution of the z component of spin $S_{\mathbf{R}}$, Fig. 3(a), exhibits at $\mathbf{k} = (\frac{\pi}{2}, \frac{\pi}{2})$ a similar symmetry as $N_{\mathbf{R}}$, Fig. 2(b) [compare $S_{\mathbf{R}}$ for $\mathbf{R} = (1, \pm 1)$ labeled with 2 and 2′, respectively]. This asymmetry disappears at $k = 0$ as it should. The nearest-neighbor spin correlation function $C_{\mathbf{R}}$, Fig. 3(b), remains negative. Hence AF correlations persist in the vicinity of the hole contrary to a ferromagnetic polaron. $C_{\mathbf{R}}$ is asymmetric as can be seen from the bonds $\mathbf{R} = (1, \pm\frac{1}{2})$, labeled with 1 and 1′. In Fig. 3(c) $j_{\mathbf{R}}$ is presented for various bonds $\mathbf{R}$ defined in Fig. 3(b). $j_{\mathbf{R}}$ is an odd function with respect to the wave vector (at $k = 0$). Because of the symmetry it vanishes also at $\mathbf{k} = (\pi, 0)$. Since the ground state has AFM long-range order, the points $\mathbf{k}$ and $\mathbf{k} + (\pi, \pi)$ are equivalent, and therefore, $j_{\mathbf{R}}$ vanishes also at $\mathbf{k} = (\frac{\pi}{2}, \frac{\pi}{2})$ similar to Ref. 9 but distinct from Ref. 4.

To test the validity of the approximations made we compared with results from exact diagonalization studies of small clusters. In Fig. 3 we compare with an 18 site cluster with one hole[8]. The ground state is at $\mathbf{k} = (\frac{2}{3}\pi, \frac{2}{3}\pi)$ for $J/t = 0.4$ [$\mathbf{k} = (\frac{\pi}{2}, \frac{\pi}{2})$ is absent in this system]. Due to the high symmetry of the 4×4 cluster[3] such subtle asymmetries of the polaron cannot be studied there. Although finite size effects in such small clusters are expected to be quite large, the agreement of $S_{\mathbf{R}}$ with the exact results for $N = 18$ is surprisingly good. Similarly $C_{\mathbf{R}}$ agrees with the exact values and the nontrivial asymmetry (e.g., between the bonds 1 and 1′) is correctly reproduced. The somewhat larger deviation of $j_{\mathbf{R}}$, Fig. 3(c), from the exact result is probably due to the strong $\mathbf{k}$ dependence of $j_{\mathbf{R}}$ which makes it very sensitive on the boundary conditions in the finite cluster. A systematic finite size analysis of such data would be required to test, e.g., the symmetry of $j_{\mathbf{R}}$ around $\mathbf{k} = (\frac{\pi}{2}, \frac{\pi}{2})$.

In Fig. 4 the spatial dependence of the bond-spin currents is shown for $\mathbf{k} = (\frac{5}{8}\pi, \frac{5}{8}\pi)$ and $J/t = 0.4$. Elser *et al.*[8] emphasize that at $J/t = 0.2$ their results for $j_{\mathbf{R}}$ exhibit strong finite size effects. Nevertheless, all results in Ref. 8 are at least qualitatively consistent with the results obtained from Eq. (2), which are free from such size effects.

To get more insight into the symmetry and asymptotic dependence of the correlation functions we simplify the wave function Eq. (2) by keeping only the one-magnon contribution, $n = 1$. If we consider the wave function only at large distances from the hole, we obtain

$$|\Psi^{(1)}_{\mathbf{k}-}\rangle = a_{\mathbf{k}} \sqrt{\frac{2}{N}} \left[\sum_{\mathbf{R}_i \in \downarrow} e^{-i\mathbf{k}\cdot\mathbf{R}_i} h^{\dagger}_{\mathbf{R}_i} + \right.$$

$$\left. + \sum_{\mathbf{R}_i \in \uparrow} e^{-i\mathbf{k}\cdot\mathbf{R}_i} h^{\dagger}_{\mathbf{R}_i} \sum_{\mathbf{R}} (\phi_0 + i\,\phi_1) S^{-}_{\mathbf{R}_i+\mathbf{R}} \right] |0\rangle, \quad (3)$$

where $\phi_0 = -2\sqrt{2}\gamma_{\mathbf{k}} t/(JR)$, $\phi_1 = -4(\mathbf{v}_{\mathbf{k}} \cdot \hat{\mathbf{R}})t/(JR)$, $\mathbf{v}_{\mathbf{k}} = \nabla_{\mathbf{k}}\gamma_{\mathbf{k}}$, and $\hat{\mathbf{R}} = \mathbf{R}/R$. The ϕ_1 term with dipolar symmetry vanishes at $\mathbf{k} = (0,0)$ and (π,π). At $(\frac{\pi}{2}, \frac{\pi}{2})$

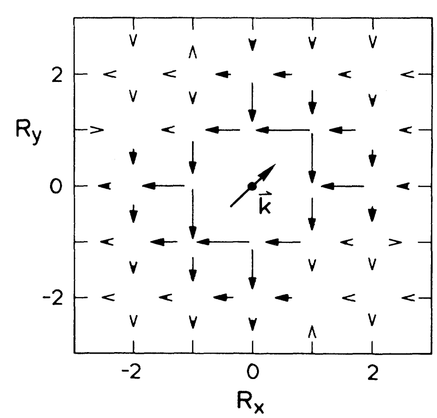

FIG. 4. Bond-spin currents $\mathbf{j}_{\mathbf{R}}$ for $J/t = 0.4$ and $\mathbf{k} = (\frac{5}{8}\pi, \frac{5}{8}\pi)$. The unit cell was $N = 16 \times 16$ and $n = 3$. The length of arrows is proportional to $|\mathbf{j}_{\mathbf{R}}|$.

ϕ_1 has its maximum while the monopole contribution ϕ_0 vanishes. This wave function has similarity to that for a ^{3}He atom moving in superfluid ^{4}He (Ref. 16) as was pointed out in Ref. 17. In this limit the correlation functions are then given by

$$\begin{aligned} N_{\mathbf{R}} &= \phi_0^2 + \phi_1^2 = 8\,[\gamma_{\mathbf{k}}^2 + 2(\mathbf{v_k}\cdot\hat{\mathbf{R}})^2]t^2/(J^2R^2), \\ S_{\mathbf{R}} &= (-1)^{R_x+R_y}(\tfrac{1}{2} - N_{\rm AFM} - N_{\mathbf{R}}), \\ C_{\mathbf{R}} &= C_\infty + \tfrac{1}{2}[(\nabla\phi_0)^2 + (\nabla\phi_1)^2], \quad C_\infty = -0.329, \\ j_{\mathbf{R}} &= (\phi_0\nabla\phi_1 - \phi_1\nabla\phi_0)\cdot\mathbf{\Delta}. \end{aligned} \tag{4}$$

These results strictly hold only asymptotically, but still reflect all symmetries found at short distances in the numerical treatment. The correlation functions decay as a power law. In particular, $N_{\mathbf{R}} \propto R^{-2}$. Although the average number of excited magnons is finite, it turns out that the change in the total number of spin deviations $\sum_{\mathbf{R}} N_{\mathbf{R}}$ diverges logarithmically. This is not unphysical, however, since the number of spin deviations per site induced by the hole still goes to zero. We note that two-magnon processes left out in Eq. (3) give corrections of order $(t/J)^2$; however, they do not remove this divergency. The results for $C_{\mathbf{R}}$–$C_\infty = 4\,(\gamma_{\mathbf{k}}^2+2|\mathbf{v_k}|^2)t^2/(J^2R^4)$ and for $j_{\mathbf{R}} = 8\sqrt{2}\gamma_{\mathbf{k}}\,[\mathbf{v_k} - (\mathbf{v_k}\cdot\hat{\mathbf{R}})\,\hat{\mathbf{R}}]\cdot\mathbf{\Delta}\,t^2/(J^2R^3)$ also do not contain quadrupolar corrections. The spin-backflow current $\mathbf{j_R}$ decays as R^{-3} and vanishes in the ground state at $\mathbf{k} = (\frac{\pi}{2}, \frac{\pi}{2})$. The missing quadrupolar contribution to $j_{\mathbf{R}}$ ($\propto R^{-4}$) becomes essential only along the line $|k_x| + |k_y| = \pi$ in the Brillouin zone where $\gamma_{\mathbf{k}} = 0$. A detailed discussion of these corrections will be given elsewhere.[14] The distortion of the AFM given by Eqs. (4) is determined by the interplay of ϕ_0 and ϕ_1 which have monopole and dipolar symmetry, respectively. We stress that the detailed form of the correlation functions differs from the results calculated by Shraiman and Siggia[4] for a wave function with a twisted order parameter.

To conclude, we have investigated various spin-correlation functions around the spin vacancy using the self-consistent Born approximation for the hole self-energy[5–7] and the corresponding wave function introduced recently by Bonfim and Reiter.[13] Our results are consistent with existing small-cluster diagonalization studies in the vicinity of the hole. Using the same method we found that the correlation functions far from the hole decay as a power law, while the QP spectral weight is finite. Remaining corrections beyond linear-spin-wave theory and the SCBA are not expected to lead to qualitative changes. This also pertains to the power-law decay of correlation functions which is determined by the linear dispersion of spin waves. The agreement of the short-range behavior with the exact diagonalization results shows that the violation of local constraints is not crucial, at least not for $J/t \geq 0.3$.

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