An alternative formalism for the slave particle mean field theory of the $t$–$J$ model without deconfinement

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Received 28 March 2005, in final form 23 May 2005
Published 20 June 2005
Online at stacks.iop.org/SUST/18/1073

Abstract

An alternative formalism that does not require the assumption of the deconfinement phase of a $U(1)$ gauge field is proposed for the slave particle mean field theory. Starting from the spin-fermion model, a spinon field, which is either fermion or boson, is introduced to represent the localized spin moment. We find a $d$-wave superconductive state in the mean field theory in the case of the fermion representation of the localized spin moment that corresponds to the slave boson mean field theory of the $t$–$J$ model, whereas the $d$-wave superconductive state is absent in case of the Schwinger boson representation of the localized spin moments.

1. Introduction

As a theory for high temperature superconductivity, the slave particle mean field theory of the $t$–$J$ model has been studied intensively [1–5]. In the slave particle formalism, the electron operator is expressed in terms of auxiliary fermions and bosons. For instance, in the slave boson formalism the electron annihilation operator $c_{j\sigma}$ at site $j$ with spin $\sigma$ is given by

$$c_{j\sigma} = b^\dagger_{j\sigma} f_{j\sigma}.$$  

(1)

Here $b^\dagger_{j\sigma}$ is a boson operator and $f_{j\sigma}$ is a fermion operator. The boson created by $b^\dagger_{j\sigma}$ is called a holon and is supposed to carry the electron’s charge, while the fermion created by $f^\dagger_{j\sigma}$ is called a spinon and is supposed to carry the electron’s spin and carry no charge.

Application of the slave boson formalism to the $t$–$J$ model suggests a simple and attractive picture for $d_{x^2-y^2}$-wave superconductivity within a mean field theory. The $d_{x^2-y^2}$-wave superconductive state is described as the state with the spinon pairing with $d_{x^2-y^2}$-wave symmetry and the holon Bose–Einstein condensation [3, 4]. The state with the same spinon pairing without holon condensation is suggested to be a pseudogap phase [6]. $U(1)$ gauge field fluctuations around the mean field state lead to non-Fermi liquid-like behaviours, such as a $T$-linear resistivity law [7–9].

Although the picture suggested by the slave boson mean field theory of the $t$–$J$ model is intriguing, this theory is based on a crucial assumption. This assumption is related to the invariance of the electron operators under the $U(1)$ gauge transformation $b_j \rightarrow b_j \exp(i\theta_j)$ and $f_{j\sigma} \rightarrow f_{j\sigma} \exp(i\theta_j)$. The assumption of the theory is that the $U(1)$ gauge field associated with this gauge symmetry is in the deconfined phase. If the $U(1)$ gauge field is in the deconfined phase, then the holons and the spinons are independent particles. In other words, the system is in the spin–charge separation phase [10]. However, Polyakov showed that the pure compact $U(1)$ gauge field theory is always confining [11]. The situation is quite unclear when there are matter fields as in the slave particle theory. Nagaosa argued that dissipation effects associated with the presence of the Fermi surface of the spinons can lead to the deconfined phase through suppression of instanton proliferation [12]. However, in [13] it is argued that this mechanism can be suppressed by screening effects. Another approach based on duality mapping suggests that the system is confining [14]. The issue of the presence of the deconfined phase of the $U(1)$ gauge field theory is still controversial [15–18].

In this paper, we will not try to directly examine the possibility of the deconfined phase within the slave particle mean field theory of the $t$–$J$ model. Instead, we derive a similar theory by taking a different formalism. We start with the spin-fermion model [19–22]. If one takes the strong coupling limit, then the model is reduced to the $t$–$J$ model. In the slave particle formalism, fermion and boson fields are introduced in the $t$–$J$ model. Here we introduce them before taking the strong
coupling limit. Since the carrier fields and the localized spin moments are independent degrees of freedom, we can express the localized spin moment either using a fermion operator or using a boson operator, as independent fields. So it is not necessary to assume the deconfined phase to make the fields independent degrees of freedom.

2. Alternative formalism for the slave boson mean field theory

The model may be given by

\[
H = -t \sum_{\langle ij \rangle} (c_i^\dagger c_j + c_j^\dagger c_i) + \frac{1}{2} J_{K} \sum_{j} (c_j^\dagger \sigma c_j) \cdot S_j + J \sum_{\langle ij \rangle} S_i \cdot S_j.
\]

(2)

Here the carrier field is represented by a spinor \(c_i^\dagger = (c_i^\dagger, c_i^\dagger)\). The components of the vector \(\sigma = (\sigma_1, \sigma_2, \sigma_3)\) are the Pauli spin matrices. The field \(S_j\) describes the localized spin at site \(j\). Here, for simplicity, we do not distinguish the oxygen lattice sites and the copper lattice sites. (More precisely, the \(c_i^\dagger\) create states represented by the Wannier function of the symmetric state of four oxygen sites around a copper ion [23].)

In the model, the largest parameter is \(J_K\), which leads to a picture of the Zhang–Rice singlet [23]. If one takes the \(J_K \to \infty\) limit, then the model is reduced to the \(t-J\) model. In the slave particle formalism, fermion and boson fields are introduced after taking the \(J_K \to \infty\) limit. Here we first introduce either a fermion or a boson field that is independent of the carrier field in order to describe the localized spin moment. The \(J_K \to \infty\) limit will be considered in the final step.

There are several ways to represent the localized spin degrees of freedom. Here we consider Abrikosov’s pseudoferon representation and the Schwinger boson representation. In the former representation, the localized spin \(S_j\) is described by

\[
S_j = \frac{1}{2} f_j \sigma f_j,
\]

(3)

with the constraint \(\sum_\alpha f_{j,\alpha}^\dagger f_{j,\alpha} = 1\). In the path-integral formalism, the partition function is given by

\[
Z = \int D\varphi D\varphi^\dagger \exp(-S),
\]

where

\[
S = \int_0^\beta \text{d}t \left[ \sum_j \varphi_j (\partial_t - \mu) c_j + \sum_{\langle ij \rangle} f_j \varphi_i f_j \right.
\]

\[
+ \sum_j \lambda_j \left( \sum_\alpha f_{j,\alpha} f_{j,\alpha} - 2S \right) - t \sum_{\langle ij \rangle} (\varphi_i c_j + \varphi_j c_i)
\]

\[
+ \frac{J_K}{4} \sum_j (\varphi_j c_j) \cdot (\varphi_j^\dagger \sigma f_j)
\]

\[
+ \frac{J}{4} \sum_{\langle ij \rangle} (\varphi_j \sigma f_i) \cdot (\varphi_j^\dagger \sigma f_j).
\]

(4)

Due to strong Kondo coupling \(J_K\), the carriers can combine with the localized spin moments to form a singlet pair [23]. This picture is justified in the \(J_K \to \infty\) limit. In order to take into account this correlation effect, we rewrite the Kondo coupling term using Stratonovich–Hubbard (SH) fields as follows:

\[
\langle \varphi_j \sigma c_j \rangle \cdot (\varphi_j \sigma f_j) \rightarrow \frac{1}{2} f_j \varphi_j - \frac{1}{2} \varphi_j^\dagger f_j c_j - f_j \varphi_j c_j - \frac{1}{2} \varphi_j \varphi_j^\dagger f_j c_j,
\]

Here we have introduced the SH fields \(s_j\) and \(\varphi_j\) that describe the formation of a singlet pair from the carrier field and the pseudofermion. Note that the SH field for triplet pair formation has been omitted because its mean field values vanish at the saddle point, reflecting the Kondo coupling being antiferromagnetic.

Our next task is to rewrite the Heisenberg term. For the Heisenberg term in equation (2), we introduce the SH fields \(D_{ij}\) and \(\varphi_{ij}\) as follows:

\[
\langle \varphi_j \sigma f_i \rangle \rightarrow \frac{1}{2} f_i \varphi_i - \frac{1}{2} \varphi_i^\dagger f_i c_i - f_i \varphi_i c_i - \frac{1}{2} \varphi_i \varphi_i^\dagger f_i c_i.
\]

(5)

Here \(D_{ij}\) is identical to the mean field used for \(d\)-wave pairing of spinons in the slave boson mean field theory of the \(t-J\) model [3, 4]. In addition, we introduce the SH fields \(\varphi_{ij}\) and \(\varphi_{ij}\) as follows:

\[
\langle \varphi_j \sigma f_i \rangle \rightarrow \frac{1}{2} f_i \varphi_i - \frac{1}{2} \varphi_i^\dagger f_i c_i - f_i \varphi_i c_i - \frac{1}{2} \varphi_i \varphi_i^\dagger f_i c_i.
\]

The fields \(\varphi_{ij}\) and \(\varphi_{ij}\) are identical to the mean fields used by Nagaosa and Lee [7, 8] to describe the normal state properties of the system.

We consider the states with uniform values of \(s_j = s_0\), \(D_{i+j} = -D_{i+j}\) and \(\varphi_{i-j} = \varphi_{i-j}\). In the momentum and the Matsubara frequency space, the action is given by

\[
S = \sum_k \left[ \left( \bar{c}_{k+} \cdot c_{k-} \right) \begin{pmatrix} -i \omega_k + \xi_k & 0 \\ 0 & -i \omega_k - \xi_k \end{pmatrix} \begin{pmatrix} c_{k+} \ \\
\bar{c}_{k-} \end{pmatrix} + \left( \bar{f}_{k+} \cdot f_{k-} \right) \begin{pmatrix} -i \omega_k + \chi_k + \lambda & \Delta_k \\ -\Delta_k & -i \omega_k - \chi_k - \lambda \end{pmatrix} \begin{pmatrix} f_{k+} \ \\
\bar{f}_{k-} \end{pmatrix} \right.
\]

\[
\times \left( \bar{c}_{k+} \cdot c_{k-} + \left( \bar{f}_{k+} \cdot f_{k-} \right) \begin{pmatrix} 0 & 0 \\ 0 & -\eta \end{pmatrix} \begin{pmatrix} f_{k+} \ \\
\bar{f}_{k-} \end{pmatrix} + \left( \bar{f}_{k+} \cdot f_{k-} \right) \begin{pmatrix} 0 & 0 \\ 0 & -\eta \end{pmatrix} \begin{pmatrix} c_{k+} \ \\
\bar{c}_{k-} \end{pmatrix} \right]\right]
\]

(5)

where \(k = (k, \omega_k)\), \(\xi_k = -2t(\cos k_x + \cos k_y) - \mu\), \(\chi_k = -\omega_0 J(\cos k_x + \cos k_y)\), \(\Delta_k = \omega_0 J(\cos k_x - \cos k_y)\) and \(\eta = 3 \sqrt{2} J_K s_0/4\). The characteristic temperature for \(D_{ij}\) is of the order of \(J\), which is identical to that of the slave boson mean field theory of the \(t-J\) model. For \(s_0\), the characteristic temperature is given by \(k_B T \sim W \exp(-1/(J_K N_F))\) with \(W\) the bandwidth and \(N_F\) the density of states at the Fermi surface of the conduction electrons within the weak coupling theory, while \(k_B T \sim J_K\) in the strong coupling limit. By integrating out the fermion fields \(\bar{f}_{k+}\) and \(f_{k-}\), we find

\[
S = \sum_k \left[ \left( \bar{c}_{k+} \cdot c_{k-} \right) \begin{pmatrix} -i \omega_k + \xi_k & 0 \\ 0 & -i \omega_k - \xi_k \end{pmatrix} \begin{pmatrix} c_{k+} \ \\
\bar{c}_{k-} \end{pmatrix} + \frac{1}{(\omega_k)^2 - (\chi_k + \lambda)^2 - |\Delta_k|^2} \begin{pmatrix} \eta^2 \Delta_k \ \\
\eta \Delta_k \end{pmatrix} \begin{pmatrix} -i \omega_k - \chi_k - \lambda \end{pmatrix} \right.
\]

\[
\times \left( \bar{c}_{k+} \cdot c_{k-} \right) \begin{pmatrix} c_{k+} \ \\
\bar{c}_{k-} \end{pmatrix} + \frac{1}{(\omega_k)^2 - (\chi_k + \lambda)^2 - |\Delta_k|^2} \begin{pmatrix} \eta \Delta_k \ \\
\eta^2 \Delta_k \end{pmatrix} \begin{pmatrix} -i \omega_k - \chi_k - \lambda \end{pmatrix} \begin{pmatrix} c_{k+} \ \\
\bar{c}_{k-} \end{pmatrix} \right]\right]
\]

(6)

The quasiparticle excitation energy spectrum \(E_k\) is determined from the following equation:

\[
E_k^2 = E_k^2(\chi_k + \lambda)^2 - \Delta_k^2 |\eta^2 \Delta_k|^2 \leq \left( E_k^2 - (\chi_k + \lambda)^2 - |\Delta_k|^2 \right)^2 \leq \left( \eta^2 \Delta_k \right)^2 \Delta_k \Delta_k.
\]

This equation has a simple solution at the \(J_K \to \infty\) limit, provided that \(s_0 \neq 0\), as follows:

\[
E_k = \pm \sqrt{(\chi_k + \lambda)^2 + \Delta_k \Delta_k}.
\]

(7)

This is the quasiparticle spectrum of the \(d_{x^2-y^2}\)-wave superconductivity and corresponds to that of the slave boson mean field theory. On the other hand, the phase with \(\eta = 0\) corresponds
to the pseudogap phase of the slave boson mean field theory. In this phase, we find that $E_1 = \xi_1$. So there is no excitation gap for the carriers. But there is an excitation energy gap with d-wave symmetry in the localized spin moment system. Therefore, we have obtained the same picture for the d-wave superconductive phase and the pseudogap phase.

3. Alternative formalism for the slave fermion mean field theory

Now we shall discuss an alternative formalism for the slave fermion mean field theory of the $t$–$J$ model. The most crucial difference is the absence of the $d$–$s$–$p$–wave superconductive state in the mean field theory, as we shall see below. Instead of the Abrikosov pseudosfermion, we use the Schwinger boson to represent the localized spin:

$$S_j = \frac{i}{2} \chi_j \sigma_3 z_j,$$

with $\langle \sum_{\sigma=\uparrow,\downarrow} z_{j\sigma} z_{j\sigma} \rangle = 2\tilde{\Delta}$. The Kondo coupling term is rewritten as $(\chi_j \sigma_3 z_j) \cdot (\bar{\chi_j} \sigma_3 z_j) \rightarrow \bar{\chi_j} \chi_j (z_j^+c_j^+ - z_j^−c_j^− - \bar{z}_j^−c_j^+)$. Here the SH fields for the triplet pair formation have been omitted as in the Abrikosov pseudosfermion case. The SH fields $\chi_j$ and $\bar{\chi_j}$ are Grassmann numbers. The term for the Heisenberg model is rewritten as follows: $(\chi_j \sigma_3 z_j) \cdot (\bar{\chi_j} \sigma_3 z_j) = -2(\bar{\chi_j}^\dagger z_j^+ - \bar{\chi_j}^\dagger z_j^−)(z_j^+\bar{z}_j^+ - z_j^−\bar{z}_j^−) + \text{const}$. The action is given by

$$S = \int_0^\beta d\tau \left[ \sum_{j,\sigma} \chi_j \sigma \bar{c}_j \chi_j - t \sum_{\langle j, k \rangle} \chi_j \sigma \bar{c}_j \chi_k + \sum_{j} \tilde{\Delta} \chi_j \sigma \bar{c}_j \right] + \frac{1}{4} J_k \sum_{\langle j, k \rangle} \bar{\chi}_j \chi_k + \frac{1}{2} \sum_{\langle j, \langle k \rangle \rangle} A_{ijk} A_{ikj}$$

$$- \frac{2}{3} J_k \sum_{\langle j \rangle} \bar{\chi}_j \chi_j (z_j^+\bar{z}_j^+ - z_j^−\bar{z}_j^−) + \chi_j \bar{\chi}_j (z_j^+ - \bar{z}_j^−) + \text{const},$$

where SH fields $A_{ijk}$ and $\bar{A}_{ijk}$ are complex numbers.

We consider the uniform solution of $\chi_j$, $\bar{\chi}_j$, $\chi_j$, and $A_{ijk}$ for the saddle point equations. This corresponds to the slave fermion mean field theory of the $t$–$J$ model. In the momentum and the Matsubara frequency space, we obtain

$$S = \sum_k \bar{\chi}_k (z_k^+c_k^+ - z_k^−c_k^−) + \sum_k \tilde{\Delta} \bar{\chi}_k (z_k^+c_k^+ - z_k^−c_k^−)$$

$$- \frac{1}{2} J_k \sum_k \chi_k (z_{-k}^+c_{-k}^+ - z_{-k}^−c_{-k}^−)$$

$$+ \frac{1}{2} \sum_k \chi_k (z_{-k}^+c_{-k}^+ - z_{-k}^−c_{-k}^−)$$

$$- \frac{1}{2} A_k \chi_k (z_{-k}^+c_{-k}^+ - z_{-k}^−c_{-k}^−),$$

with $\chi_k = (\sin k_x \pm \sin k_y) / 2$ where the plus sign is for $A_{i,\pm \pm} = A_{i,\pm \mp} = A$ and the minus sign is for $A_{i,\pm \pm} = -A_{i,\pm \mp} = A$.

Integrating out Schwinger bosons yields

$$S_{\text{eff}} = \frac{9}{16} \sum_k \frac{1}{\chi_k^2 (\chi_0 - \lambda)}$$

$$\times \left( \begin{array}{cc} \chi_k^+ & c_{-k}^+ \\ \chi_k^{-} & c_{-k}^{-} \end{array} \right),$$

It has been used that the square of a Grassmann variable, such as $\chi_k^2$, is zero. Note that there is no off-diagonal term at the mean field level. Therefore, the superconductive state is absent at least within the mean field theory.

4. Holon operator

In the above formalism, we have represented the localized spin either by the Abrikosov pseudosfermion or by the Schwinger boson. This corresponds to representing the spinon by a fermion or a boson, respectively. Now we show that the SH fields $s_j$ and $\chi_j$ can be seen as the holon annihilation operator.

We first consider the case of the Abrikosov pseudosfermion representation. The basis states of the Hilbert space at site $j$ are represented as $| c, f \rangle = | c \rangle \otimes | f \rangle$, where $| c \rangle$ denotes the carrier state and $| f \rangle$ denotes the localized spin state. We assume that the doubly occupied state at the copper site is excluded due to the strong on-site Coulomb repulsion. Then, the constraint is $\sum_c f_c^+ f_c \leq 1$. Thus the basis of the Hilbert space is $\{ 0, | \uparrow \rangle, | \downarrow \rangle, | \uparrow \downarrow \rangle \}$. Here $0$ represents the unoccupied state etc. In this basis, states like $| \uparrow \downarrow \rangle$ have been excluded because of the strong antiferromagnetic Kondo coupling.

By constructing the matrix representation of $s_j$, which has the form of $s_j = \frac{1}{\sqrt{2}} (f_j^+ c_{j\uparrow} - f_j^+ c_{j\downarrow})$ at the saddle point, in the above Hilbert space from those of $f_c$ and $c_f$, we find that the operation of $s_j$ does not change the total spin at site $j$ but annihilates the charge by one at site $j$. Therefore, $s_j$ is like the holon annihilation operator with bosonic statistics. However, the identification of $s_j$ with the holon operator is incomplete. For the full correspondence, we need to assume that the annihilation of the spin $\sigma$ spinon is equivalent to the creation of the spin $-\sigma$ spinon. This point may be more clarified if we formulate the theory in terms of the $SU(2)$ doublets introduced in [24]. Representing the spinons in terms of those $SU(2)$ doublets would give us an alternative to the $SU(2)$ formalism of the slave particle theory [25]. Similarly, we can show that $\chi_j$ is like the holon annihilation operator with fermionic statistics. The dynamics of the holons is determined by the action that is derived by integrating out the carrier fields and the localized spin moment fields. In the case of the Abrikosov pseudosfermion representation, the resulting action is $S_h = \sum_{q, r, \Omega_2} K(q, i\Omega_2) \bar{\eta}(q, i\Omega_2) \eta(q, i\Omega_2)$, where $K(q, i\Omega_2) = (1/\beta) \sum_{k, \omega} \text{Im} \sum_{\omega} \text{Im} (\tilde{\Omega}_0 (\omega_0^+ + i\Omega_2) - (\chi_k + \lambda) \tilde{\Omega}_k^q) \sqrt{|(\omega_0^2 - E_k^2)|}$.

5. Gauge field and deconfinement

Now we discuss the $U(1)$ gauge field. The phase fluctuation of $\chi_j$ corresponds to the $U(1)$ gauge field considered in [7, 8]. Indeed, the gauge field is associated with the phase fluctuations of the same SH field for the spinons. As argued in [7, 8], the transition temperature, which is of the order of $J$ within the mean field analysis, can be suppressed by the gauge field fluctuations. Since the effective gauge field propagator depends on the dynamics of the holon, which is governed by $S_h$, some physical properties of the slave particle formalism, such as the $T$-linear resistivity law, could be different. This point is left for future study.

Now we consider the deconfinement nature of the gauge field in the present formalism. Contrary to the slave boson
theory case [7, 8], the spinon fields and the carrier fields are independent fields because we have started from a spin-fermion-like model in which the carriers and the localized spin moments are independent. In this sense the spinons are not confined to the electrons. However, one can show that in the Néel ordered phase the $U(1)$ gauge field associated with the phase fluctuations of $\chi_{ij}$ leads to a logarithmic confining potential between the Abrikosov pseudofermions. In other words, there are no low lying spin 1/2 excitations in the localized spin system, whereas in the disordered phase the logarithmic potential is replaced by a potential of exponential decay. Therefore, the gauge field is not confining in this phase and the fermions are no longer confined to pairs.

6. Conclusion

In this paper, we have proposed an alternative formalism for the slave particle theory of the $t$–$J$ model. Here we have started from the spin-fermion model which is a multiband model and is reduced to the $t$–$J$ model by taking the strong Kondo coupling limit. The strong coupling limit is taken after introducing fields that describe the localized spin moments. Since the spin-fermion model is a multiband model, the carrier fields and the fields describing the localized spin moments are independent fields.

The main results concerning the superconductive states of the slave particle theory at the mean field level are reproduced by the present formalism. The gauge field fluctuations lead to a confinement phase of the spin 1/2 particles that describe the localized spin moments in the Néel ordering phase. Meanwhile, those particles are not confined in the disordered phase.

Acknowledgments

This work was partially supported by a Grant-in-Aid for the 21st Century COE ‘Center for Diversity and Universality in Physics’ from the Ministry of Education, Culture, Sports, Science and Technology (MEXT) of Japan.

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