1. Introduction

The discovery of pnictide superconductors (SCs) has triggered lots of attention from the condensed matter community. Unlike cuprates, where only the Cu $3d_{x^2-y^2}$ orbital plays a significant role, local-density approximation (LDA) calculations [1, 2] indicate that pnictide SCs have several active 3d orbitals near the Fermi surface (FS). Consequently, it is widely believed that such SCs should be understood in terms of multi-orbital models instead of the single-orbital ones [3–5]. Regarding the minimal model capable of capturing the essential physics of pnictide SCs, some authors proposed more realistic three- and five-orbital models [4, 5], while others argued that the main physics of pnictide SCs is contained in two-orbital models [3, 6, 7]. Because of their relative simplicity, as well as the fact that [6] the correct FS shape can be reproduced in both the doped and undoped cases, it is crucial to find the properties of these two-orbital models.

Most previous theoretical work on two-orbital models for pnictides has been based on mean-field-like approximations, such as the random phase approximation (RPA), fluctuation exchange (FLEX) and functional renormalization group (fRG) calculations. The main results from these studies give a good understanding of pnictides while still giving very different pairing pictures. For instance, Graser et al proposed two nearly degenerate competing pairing states with $A_{1g}$ and $B_{1g}$ symmetries because of the near nesting of FS sheets [8]. Although other RPA studies also suggested a competing pairing picture, different pairing channels were proposed, for example, the competition between singlet $d$-wave and triplet $p$-wave states [9] or $s_{\pm}$-wave and $d$-wave states [10]. Aside from RPA studies, a FLEX calculation demonstrated an $s_{\pm}$-wave or $d_{xy}$-wave pairing depending on whether the intraband antiferromagnetic (AFM) spin fluctuation was stronger than the interband one or not [11]. Meanwhile, the fRG approach...
revealed $s_\pm$-wave and sub-dominant $d$-wave pairings [12]. Moreover, Dai et al suggested a spin triplet pairing state after using the BCS mean field method [13].

From the above discussions, it seems to be very difficult to justify the pairing symmetry of the two-orbital models. While it is unrealistic to take the full quantum fluctuations into account in the usual theoretical methods, the important role of electronic correlations in magnetism and superconductivity has not been fully recognized. Despite the fact that there exist some unbiased numerical investigations of the two-orbital model [6], [14–16], these results are insufficient for understanding the electronic correlations in two-dimensional systems because they were obtained for either an eight-site lattice [6, 14, 15] or a diagonal ladder [16]. To give a comprehensive and systematic understanding of the two-orbital model [3, 14], a quantum Monte Carlo (QMC) method is employed in this paper.

However, it is known that there exist several difficulties in QMC simulations of two-orbital models: One is the severe limitation of the cluster size. For example, the computational demands of a two-orbital model with a given cluster are much higher than those of a one-orbital model with a double cluster size. Another problem and probably the toughest, is the insurmountable Fermi sign problem [17, 18]. Unlike a one-orbital model, the usual discrete Hubbard–Stratonovich transformation [19] for the Hund’s coupling and pair-hopping terms, which is specific to the two-orbital models, leads to a more serious sign problem. Sakai et al [18] proposed a new type of transformation for the Hund’s coupling and pair-hopping interaction, which can effectively alleviate the sign problem. Based on this progress, we developed a feasible constrained-path Monte Carlo (CPMC) method [20] for two-orbital models, which works well in the weak and intermediate correlation regimes.

Using the CPMC method, we compute the magnetic structure factors and the pairing correlations of the two-orbital model as functions of the doping density $\rho$ and interaction strength $U$. We find a $(\pi, 0)/(0, \pi)$ magnetic order that is enhanced by the Coulomb repulsion $U$ and the Hund’s coupling $J$ in the undoped case. Because of the particle–hole asymmetry of the two-orbital model, this magnetic order behaves differently in the electron- and hole-doped cases. We also find that the doping has a much stronger effect than that of the Coulomb repulsion on the pairing correlations. In the electron-doped case, a nodeless $s_\pm$ pairing is dominant, whereas in the hole-doped case, two degenerate-like $A_{1g}$ and $B_{1g}$ intraband singlet pairings compete with each other and become the dominant ones.

Our paper is organized as follows: In section 2, we briefly introduce the two-orbital model under investigation and discuss the proper choice of model parameters. Some modifications to the original CPMC algorithm and definitions of the calculated physical quantities are presented as well. In section 3, we show in detail the simulation results for the magnetic and pairing properties of the model with various parameters. Our main conclusions are summarized in section 4.

2. Model and numerical approach

On the basis of LDA calculations, Mazin et al [2] advocated that the band structure of pnictides involves only three Fe 3d orbitals, $d_{yz}$, $d_{zx}$ and $d_{xy}$ (or $d_{x^2−y^2}$), near the Fermi level. Accordingly, Raghu et al [3] introduced a minimal multi-orbital model for pnictide SCs with the further approximation that a next-neighbour hybridization between the $d_{yz}$ and $d_{zx}$ orbitals can be equated to the role of the $d_{xy}$ or $d_{x^2−y^2}$ orbital. As described in [14], the kinetic part of the two-orbital model Hamiltonian is given by

$$H_0 = \sum_{i,\sigma} \left( \begin{array}{l} d_{i,\sigma}^\dagger d_{i,\sigma} + \frac{1}{2} t_{1,\sigma} (d_{i,\sigma}^\dagger d_{i+\sigma,\sigma} + d_{i+\sigma,\sigma}^\dagger d_{i,\sigma}) + h.c. \end{array} \right)$$

where $x$ and $y$ represent the $d_{yz}$ and $d_{zx}$ orbitals, respectively. The operator $d_{i,\sigma}$ creates an electron on orbital $\alpha$ in Fe site $i$ with spin $\sigma$ and the index $\hat{\mu} (\hat{\nu}) = \hat{x}$ or $\hat{y}$ denotes a unit vector linking the nearest-neighbour sites. When estimating the hopping amplitudes to recover the right topology of the FS and band features given by DFT [21, 22], the band-structure calculation [3] and the Slater–Koster tight-binding scheme [6] recommended different hopping amplitudes; however, the Lanczos study of an eight-site cluster suggested these two schemes give similar physics [6]. Following the band-structure calculation, the hopping parameters will always be taken as $t_1 = -1.0$, $t_2 = 1.3$ and $t_3 = t_4 = -0.85$ in our calculations.

The interaction terms [14, 23, 24], with a Hubbard repulsion in the same orbital, a repulsion $U'$ for different orbitals, a ferromagnetic Hund’s coupling $J$ and pair-hopping terms, can be expressed as

$$H_{\text{int}} = \sum_i \left( H_i^1 + H_i^2 + H_i^3 + H_i^4 \right),$$

$$H_i^1 = J \sum_{\alpha \neq \alpha'} \left( d_{i,\alpha}^\dagger d_{i+\alpha,\alpha'}^\dagger d_{i+\alpha',\alpha} d_{i,\alpha'} + d_{i+\alpha',\alpha}^\dagger d_{i,\alpha}^\dagger d_{i+\alpha,\alpha'} d_{i,\alpha'} \right),$$

$$H_i^2 = (U' - J) \sum_\sigma n_{i,\sigma} n_{i,\sigma},$$

$$H_i^3 = U' \sum_\sigma n_{i,\sigma} n_{i+\sigma},$$

$$H_i^4 = U' \sum_\sigma n_{i,\sigma} n_{i−\sigma},$$

where $\alpha$ denotes the $d_{yz}$ or $d_{zx}$ orbital and $U'$ satisfies the constraint $U' = U - 2J$ due to the rotational invariance [25]. Throughout this work, the correlation strength is taken up to the intermediate range, i.e., $1/|U'| \leq 2$ for both undoped and...
doped cases, which is believed to be correct for pnictides SCs [26].

In equation (3), $H_i$ can be transformed [18] as

$$e^{-\Delta t\phi} = \frac{1}{2} \left[ \sum_{j=1}^{n} e^{i \phi (N_{j-1} + N_{j} + b_{j-1}N_{j} + b_{j}N_{j-1})} \right]$$

(7)

with

$$f_{i,a} = d_{i,a}^\dagger a_i + d_{i,a} a_i,$$

(8)

$$N_{i,a} = n_{i,a} + n_{i,a} - 2n_{i,a} n_{i,a}$$

(9)

where $a$, $b$, and $\phi$ are some parameters depending on Hund’s coupling $J$ and Trotter interval $\Delta \tau$ and $\gamma = \pm 1$ is the newly introduced auxiliary field.

Due to the property that $N_{i,a} = N_{i,a}$, the factor $e^{N_{i,a} N_{i,a}}$ in equation (7) can be further decoupled into a product of single $e^{N_{i,a}}$-like terms using the discrete Hubbard–Stratonovich transformation [19]. Then all the terms containing $e^{N_{i,a}}$, which are independent of the introduced field $\gamma$ in equation (7), can be combined with $H_i$ in equation (4) in the ordinary CPMC treatment. However, after this recombination, we can see that the remaining factor in equation (7), $e^{\phi(N_{i,-1} - N_{i+1})}$, unlike other interactions, which use the number operator $n_{i,a}$, involves some hopping-like terms. So some adjustment must be made for this new item $e^{\phi(N_{i,-1} - N_{i+1})}$.

Recalling that in the standard QMC algorithm, the matrix form of the interaction term, such as the Hubbard repulsion $H_i$, always has the form:

$$e^{\mu_i I + A},$$

(10)

where $A$ is sparse with one element in the diagonal and $I$ is the identity matrix. Consequently, the determinant division $L(I+A)R/\det LR$ and the matrix inverse $(L(I+A)R)^{-1}$ can be calculated using a fast updating algorithm [27, 28].

We find that the matrix form of $e^{\mu_i I + A}$ can be cast into a similar form as equation (10):

$$e^{\phi(N_{i,a} N_{i,a} + h.c.)} = I + B$$

(11)

but with $B$ having four non-zero elements

$$B = \begin{pmatrix} 
0 & \cdots & \cdots & 0 \\
\vdots & \ddots & \ddots & \vdots \\
\vdots & \ddots & \ddots & \vdots \\
0 & \cdots & \cdots & 0 
\end{pmatrix}$$

(12)

where $b_{nn} = b_{mm} = \frac{e^{-i\phi} + e^{i\phi}}{2}$, $b_{mn} = b_{mn} = -\frac{e^{-i\phi} + e^{i\phi}}{2}$, $m$, $n$ refer to the positions of orbitals $d_{i,a}$, $d_{i,a}$, which are acted on by $d_{i,a}^\dagger$, $d_{i,a}^\dagger$ in $e^{\phi d_{i,a}^\dagger d_{i,a}}$. If we insert the unitary matrix $UU^\dagger(I+B)UU^\dagger$ to make $U(I+B)U = I + B$ with $B$ the desired diagonal form as $A$ in equation (10), the determinant division $L(I+BR)\det LR$ and the matrix inverse $(L(I+B)R)^{-1}$ can then be written as

$$\frac{\det L(I+B)R}{\det LR} = \frac{\det L(I+B')R'}{\det LR'},$$

(13)

$$\frac{(L(I+B)R)^{-1} = (L'(I+B')R')^{-1}}{R'},$$

(14)

where $L' = LU$ and $R' = U^{-1}R$. Now the standard CPMC algorithm can be applied with the new formulas of equation (13).

In order to investigate the magnetic properties, we examine the magnetic correlations through the static magnetic structure factor

$$S(k) = \frac{1}{N} \sum_{\gamma} e^{\gamma_\gamma(k) \phi} ((n_{i_1} - n_{i_2}) (n_{j_1} - n_{j_2})),$$

(15)

where $n_{ia} = n_{i,a} + n_{i,a}$.

For the pairing properties, the classification of possible pairing symmetries in [29] is followed (see table 1). In multi-orbital systems, the pairing operators have both spatial and orbital degrees of freedom [30]. The singlet and triplet (with projection 1) pairing operators, $\Delta_i(k)$ and $\Delta'_i(k)$, can be defined as

$$\Delta_i(k) = \frac{1}{\sqrt{2}} f(k) (\tau_i)_{aa'} \left( d_{k,a}^\dagger d_{k,a'} - d_{k,a'}^\dagger d_{k,a} \right),$$

(16)

$$\Delta'_i(k) = f(k) (\tau_i)_{aa'} \left( d_{k,a}^\dagger d_{k,a'}^\dagger - d_{k,a}^\dagger d_{k,a'} \right),$$

(17)

where $d_{i,a}$ creates an electron in orbital $a$ with momentum $k$ and spin $\sigma$ and $f(k)$ is the form factor that transforms according to one of the irreducible representations of the symmetry group [30] (for concrete forms see table 1), while the $\tau_i$’s are the Pauli matrices ($i = 1, 2, 3$) or identity matrix ($i = 0$). Using the Fourier transformation, we can get the pairing operator in coordinate space $\Delta(i)$. The corresponding pairing correlation function is defined as

$$P(r = |i-j|) = \langle \Delta(i) \Delta(j) \rangle.$$  

(18)

In the simulations, we use the free-electron wave function as the trial wave function. Our CPMC code is checked by comparing our results to the Lanczos results for the $2 \times 2$ and $3 \times 2$ clusters and also to a previous eight-site cluster Lanczos simulation [6]. The maximum differences in total energy and magnetic moment were less than 3% up to $U = 2.5$. 

### Table 1. The possible nearest-neighbour-bond pairing basis matrices of the two-orbital models used in our simulations [29]. The first column is the index number, the second and third columns list the representations and the basis matrices $f(k)$, the last column shows the spin parities where $S$ refers to singlet and $T$ to triplet. Note that a nodeless $s_\alpha$ is also listed in the first row.

| No. | IR | $f(k)$ | Spin |
|-----|----|-------|------|
| 1   | $A_{1g}$ | $(\cos k_x \cos k_y, \sin k_x \sin k_y)$ | $S$ |
| 2   | $A_{1g}$ | $(\cos k_x + \cos k_y, \sin k_x - \sin k_y)$ | $S$ |
| 3   | $A_{1g}$ | $(\cos k_x - \cos k_y, \sin k_x + \sin k_y)$ | $S$ |
| 4   | $A_{2g}$ | $(\cos k_x, \sin k_y)$ | $S$ |
| 5   | $A_{1g}$ | $(\cos k_x \cos k_y, \sin k_x \sin k_y)$ | $T$ |
| 6   | $B_{2g}$ | $(\cos k_x + \cos k_y, \sin k_x - \sin k_y)$ | $T$ |
| 7   | $B_{1g}$ | $(\cos k_x - \cos k_y, \sin k_x + \sin k_y)$ | $T$ |
| 8   | $B_{2g}$ | $(\cos k_x, \sin k_y)$ | $T$ |
| 9   | $E_g$ | $\sin k_x \sin k_y$ | $T$ |
| 10  | $E_g$ | $\sin k_x \sin k_y$ | $T$ |
| 11  | $E_u$ | $\sin k_x \sin k_y$ | $T$ |
3. Results and discussions

3.1. Magnetic Property

First we discuss the magnetic order in the undoped system. As shown in figure 1(a), the magnetic structure factor $S(k)$ is presented at half filling (one electron per orbital) for different Coulomb repulsions $U$ and Hund’s couplings $J$ for a $6 \times 6$ lattice. It is obvious that the sharp peak at $(\pi, 0)/(0, \pi)$ persists at various $U$ and $J$, signifying a robust $(\pi, 0)/(0, \pi)$ magnetic order. In addition, this stable spin order still persists for an $8 \times 8$ lattice (see figure 2(b)). It is worth noting that the $(\pi, 0)/(0, \pi)$ peak in $S(k)$ cannot be viewed as a criterion for the formation of the striped AFM order [31, 32], as we will discuss later that another proposed magnetic order, the orthomagnetic (OM) order [33], also has a similar magnetic structure.

In figure 1(a), we see that on increasing the Coulomb repulsion $U$, the magnetic order is enhanced. Since the strength of the Coulomb repulsion (in units of $|t_{ij}|$) can be viewed as a measurement of the electronic correlation strength, this $U$-induced enhancement implies the important role of electronic correlations for the investigated magnetic order. Similarly, an enhancement in the magnetic order is again observed on increasing the Hund’s coupling $J$ at fixed $U = 2.0$ (see figure 1(b)), considering that $J$ favours the local magnetic moments, which also signals possible contributions of the local moments to this magnetic order. Within the same argument, the robust $(\pi, 0)/(0, \pi)$ peak at $U = 0.0$ (see figure 1(a)) indicates that the magnetic order not only relates to the electronic correlations and local moments, but also to other factors, such as the FS nesting.

Next we discuss the effects of doping on the magnetic order. On electron doping, as shown in figure 2(a), the $(\pi, 0)$ peak seems to be unaffected compared with the undoped case initially, but the values of $S(k)$ along the $(\pi, \pi)-(0, 0)$ direction are strongly suppressed. As a result, the magnetic order is relatively enhanced. When more electrons are doped, $S(0, 0)$ starts falling and a probable incommensurate magnetic structure around $(\pi, 0)$ arises. Unlike previous studies, we find that the effect of electron doping on the magnetic order is not a monotonic suppression and there may exist a small regime close to half filling where the magnetic order is enhanced or at least unaffected by doping. Similar phenomena are also observed for the $8 \times 8$ lattice as shown in figure 2(b).

In the hole-doped case, however, because of the particle–hole asymmetry of the two-orbital model, the behaviour of $S(k)$ with doping is different. In figure 3, the $(\pi, 0)$ peak is directed suppressed even at very low doping densities and the values along the $(\pi, \pi)-(0, 0)$ direction are relatively insensitive to the doping concentration. Interestingly, as reflected in figure 4, the different behaviours of the magnetic order for different dopants seem to be closely associated with their different FS evolutions upon doping: With exactly the same doping density, the electron pocket is notably diminished by hole doping while that of the electron-doped system is just slightly enlarged. On the other hand, in both cases the hole pockets almost remain unchanged. These facts may imply that the FS nesting remains in good condition at low electron doping while it weakens for strong electron or hole doping. This explains why the enhancement of magnetic order is observed only at low electron
doping. Therefore, we propose that, at least in the intermediate interaction regime, FS nesting plays an important role in the magnetism of a two-orbital system.

Now we analyse the competing magnetic orders of the two-orbital model at half filling. As proposed in [33], the OM order, in which the magnetic moments on nearest-neighbour sites are at right angles, is recommended for a two-orbital model. Numerically, it is rather difficult to distinguish the striped AFM and the OM orders: both have similar magnetic structure factors, negative next-nearest-neighbour spin–spin correlations and almost-zero expectations of the nearest-neighbour spin–spin correlations [6]. In order to identify the competing magnetic orders at half filling, we calculate the expectation values of the four-spin-operator \(\langle S_i \cdot S_{i+x} \rangle\). If \(U\) favours the OM order, \(\langle S_i \cdot S_{i+x} \rangle\) should grow slower than \(\langle S_i^2 \rangle\). As a result, \(\langle S_i^2 \rangle - \langle S_i \cdot S_{i+x} \rangle\) should increase on increasing \(U\). In figure 5, a clear \(U\)-dependent enhancement of \(\langle S_i^2 \rangle - \langle S_i \cdot S_{i+x} \rangle\) is observed for the 6 × 6 and 8 × 8 lattices, which implies a strong tendency for the formation of the OM order as \(U\) is increased. In addition, this tendency becomes stronger when the lattice size is enlarged from 6 × 6 to 8 × 8.

To substantiate this argument, we also calculate \(\langle (S_i \cdot S_{i+x+y})^2 \rangle - \langle (S_i \cdot S_{i+x})^2 \rangle\). Similarly, if \(U\) favours the OM order, the nearest-neighbour spin–spin correlation ought to grow slower than the next-nearest-neighbour one. Then, \(\langle (S_i \cdot S_{i+x+y})^2 \rangle - \langle (S_i \cdot S_{i+x})^2 \rangle\) should also be enhanced by \(U\), which is demonstrated by the results presented in the inset of figure 5.

From the above discussions, we conclude that, at least in the weak to intermediate electronic correlation regime, the magnetic order at half filling in the two-orbital model tends to be the OM order. Similar conclusions are drawn from unrestricted Hartree–Fock [33] and Density Matrix Renormalization Group (DMRG) [16] studies of the same model for other lattices.

### 3.2. Pairing symmetry

Since the pairing symmetry is intricately related to the pairing mechanism, it is essential to clarify the dominant pairing channel among all the possible candidates. In this section, the long-range pairing correlations of the possible nearest-neighbour-bond pairing states [29] and a proposed nodeless \(s_\pm\) pairing state [6, 15] are discussed (see table 1) and subsequently, the effects of doping density \(\rho\) and...
The non-on-site pairing correlations $P(r)$ as functions of the pairing distance $r$ for the hole-doped case: (a) and (c): intraband singlet pairings 2, 7 versus all the interband singlet pairings. (b) and (d): intraband singlet pairing 2, 7 versus all the triplet pairings. As in figure 6, the dashed line represents the $s_{z}$ pairing. In this case, four holes are doped in the $6 \times 6$ lattices and eight holes in the $8 \times 8$ lattices with $U = 1.4$, $J = 0.25U$.

Figure 7.

The average of the long-range pairing correlation $P_{\text{ave}}$ for 2, 7 and $s_z$ as functions of the doping density $\rho(a)$ and (c) and the Coulomb repulsion $U(b)$ and (d) for the $6 \times 6$ lattice. (a) and (b) correspond to the electron-doped cases and (c) and (d) the hole-doped cases. Here, eight electrons are doped in (b) and 12 holes in (d).

The number of pairs, is plotted in figure 8 as functions of $\rho$ and $U$ for the electron- and hole-doped cases. From figures 8(a) and (c), we observe that in both the electron- and hole-doped cases, the $s_z$ pairing, together with pairings 2 and 7, are suppressed on increasing the doping density $\rho$. Obviously, the $s_z$ pairing is dominant in the electron-doped case, whereas in the hole-doped case the suppression of the $s_z$ pairing is more drastic than that of the pairings 2 and 7 and unlike the electron-doped case, pairings 2 and 7 become dominant when $\rho > 0.06$.

Finally, with a fixed doping density $\rho$, we study the effect of the Coulomb repulsion $U$ on the pairing correlations, as presented in figures 8(b) and (d). Overall, the effect of $U$ is much weaker than that of doping—the pairing properties are almost unchanged as $U$ is increased. The reason for this is that the doping modifies the structure of the FS significantly and is closely related to the nesting. Therefore, at least in the weak to intermediate interaction regimes, other factors, such as FS nesting rather than $U$, play a more important role for the pairing correlations.

Consistent with figures 8(a) and (c), the $s_z$ pairing prevails over the pairings 2 and 7 in the electron-doped case and the latter become the leading channels in hole doping. Thus our results demonstrate dopant-dependent pairing symmetries in the two-orbital model.

4. Conclusions

In this paper, we have systematically studied the magnetic and pairing properties of the two-orbital model for pnictides...
at half filling and for electron- and hole-doped cases. We found that the \((\pi, 0)/(0, \pi)\) magnetic order is robust at half filling in the weak to intermediate interaction regime. On increasing the Coulomb repulsion \(U\), the magnetic order is enhanced and the system tends to be in the OM order, which is consistent with unrestricted Hartree–Fock and DMRG studies [16, 33].

When the system is doped away from half filling, the magnetic order behaves differently for electron and hole doping: it is relatively enhanced for electron doping and suppressed eventually, while for the hole-doped case, the magnetic order is directly suppressed. This difference is closely relevant to the different evolution of the FS when electrons and holes are doped into the system—the FS nesting remains in good condition for light electron doping while for the hole-doped case the electron pocket shrinks significantly and thus nesting can hardly be realized.

The strong effects of doping on the long-range pairing correlations were also observed for the two-orbital model. For the electron-doped case, an \(s_a\) pairing state dominates the possible nearest-neighbour-bond pairing channels, while two nearly degenerate intraband singlet pairing channels with \(A_{1g}\) and \(B_{1g}\) symmetries take over in the hole-doped case, which illustrates the dopant-dependent pairing property of the two-orbital model.

Acknowledgments

We thank Adriana Moreo for useful discussions. This work was supported by NSFC under Grants Nos. 11174072 and 91221103 and by SRFDP under Grant No. 20104208110001.

References

[1] Vildosola V, Pourrovskii L, Arita R, Biermann S and Georges A 2008 Phys. Rev. B 78 064518
[2] Mazin I I, Singh D J, Johannes M D and Du M H 2008 Phys. Rev. Lett. 101 057003
[3] Raghu S, Qi X L, Liu C X, Scalapino D J and Zhang S C 2008 Phys. Rev. B 77 220503
[4] Daghofer M, Nicholson A, Moreo A and Dagotto E 2010 Phys. Rev. B 81 014511
[5] Kuroki K, Onari S, Arita R, Usui H, Tanaka Y, Kontani H and Aoki H 2008 Phys. Rev. Lett. 101 087004
[6] Moreo A, Daghofer M, Riera J A and Dagotto E 2009 Phys. Rev. B 79 134502
[7] Hu J and Hao N 2012 Phys. Rev. X 2 021009
[8] Graser S, Maier T A, Hirschfeld P J and Scalapino D J 2009 New J. Phys. 11 025016
[9] Qi X L, Raghu S, Liu C X, Scalapino D J and Zhang S C 2008 arXiv:0804.4332v2
[10] Bang Y and Choi H Y 2008 Phys. Rev. B 78 134523
[11] Yao Z J, Li J X and Wang Z D 2009 New J. Phys. 11 025009
[12] Wang F, Zhai H, Ran Y, Vishwanath A and Lee D H 2008 arXiv:0805.3343v3
[13] Dai X, Fang Z, Zhou Y and Zhang F C 2008 Phys. Rev. Lett. 101 057008
[14] Daghofer M, Moreo A, Riera J A, Arrigoni E, Scalapino D J and Dagotto E 2008 Phys. Rev. Lett. 101 237004
[15] Nicholson A, Ge W, Zhang X, Riera J, Daghofer M, Oleš A M, Martins G B, Moreo A and Dagotto E 2011 Phys. Rev. Lett. 106 217002
[16] Berg E, Kivelson S A and Scalapino D J 2010 Phys. Rev. B 81 172504
[17] Held K and Vollhardt D 1998 Eur. Phys. J. B: Condens. Matter Complex Syst. 5 473
[18] Sakai S, Arita R and Aoki H 2004 Phys. Rev. B 70 172504
[19] Hirsch J E 1983 Phys. Rev. B 28 4059
[20] Zhang S, Carlson J and Gubernatis J E 1997 Phys. Rev. B 55 7464
[21] Singh D J and Du M H 2008 Phys. Rev. Lett. 100 237003
[22] Xu G, Ming W, Yao Y, Dai X, Zhang S C and Fang Z 2008 Europhys. Lett. 82 67002
[23] Bascones E, Calderón M J and Valenzuela B 2010 Phys. Rev. Lett. 104 227201
[24] Luo Q, Martins G, Yao D X, Daghofer M, Yu R, Moreo A and Dagotto E 2010 Phys. Rev. B 82 104508
[25] Dagotto E, Hotta T and Moreo A 2001 Phys. Rep. 344 1
[26] Dai F, Hu J and Dagotto E 2012 Nature Phys. 8 709–718
[27] White S R, Scalapino D J, Sugar R L, Loh E Y, Gubernatis J E and Scalatter R T 1989 Phys. Rev. B 40 506
[28] Hanke W and Kopae V (eds) 1992 Electronic Phase Transitions (Modern Problems in Condensed Matter Sciences) (Amsterdam: North-Holland)
[29] Wan Y and Wang Q H 2009 Europhys. Lett. 85 57007
[30] Moreo A, Daghofer M, Nicholson A and Dagotto E 2009 Phys. Rev. B 80 104507
[31] Dong J et al 2008 Europhys. Lett. 83 27006
[32] Paglione J and Greene R L 2010 Nature Phys. 6 645–658
[33] Lorenzana J, Seibold G, Ortiz C and Grilli M 2008 Phys. Rev. Lett. 101 186402