Self-energy enhancements in doped Mott insulators

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We analyze enhancements in the magnitude of the self-energy for electrons far away from the Fermi surface in doped Mott insulators using the dynamical cluster approximation to the Hubbard model. For large onsite repulsion, U, and hole doping, the magnitude of the self-energy for imaginary frequencies at the top of the band (k = (π, π)) is enhanced with respect to the self-energy magnitude at the bottom of the band (k = (0, 0)). The self-energy behavior at these two k-points is switched for electron doping. Although the hybridization is much larger for (0, 0) than for (π, π), we demonstrate that this is not the origin of this difference. Isolated clusters under a downward shift of the chemical potential, μ < U/2, at half-filling reproduce the overall self-energy behavior at (0, 0) and (π, π) found in low hole doped embedded clusters. This happens although there is no change in the electronic structure of the isolated clusters. Our analysis shows that a downward shift of the chemical potential which weakly hole dopes the Mott insulator can lead to a large enhancement of the (π, π) self-energy which is not necessarily associated with electronic correlation effects, even in embedded clusters.

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I. INTRODUCTION

Understanding the electronic properties of two-dimensional metals close to the Mott insulator transition remains a formidable theoretical challenge. A remarkable example is found in the cuprates in which high-Tc superconductivity arises when doping the Mott insulator. Although these systems have been studied intensively over the past decades there is a lack of consensus on the mechanism of superconductivity. The simplest electronic correlated model which can capture the electronic properties and phase diagram of the cuprates is the Hubbard model on a square lattice. Recent progress in numerical approaches12,15,20 to strongly correlated electron systems allows for an accurate determination of the electron spectra of the Hubbard model even in the relevant but difficult regime of a large onsite Coulomb repulsion, U. The electron spectra obtained from these approaches can be compared with ARPES experiments16 testing the validity of the model.

The self-energy in imaginary frequencies, Σk(iωn), is the key quantity encoding the strength of electron correlations. determining the Greens function through Dyson’s equation:

\[ G_k(i\omega_n) = \frac{1}{i\omega_n + \mu - e_k - \Sigma_k(i\omega_n)}, \]

where: \( \omega_n = \frac{(2n+1)\pi}{\beta} \), are Matsubara frequencies, n is an integer, and \( \mu \) the chemical potential, \( \beta \) the inverse of the temperature, \( \beta = 1/T \). Deviations from independent electron behavior due to Coulomb interactions can be monitored through the quasiparticle weight, \( Z_k \), obtained from17:

\[ Z_k = \frac{1}{1 - \frac{\text{Im}\Sigma_k(i\omega_n)}{\omega_n}} \bigg|_{\omega_n \to 0}. \]

In the dynamical mean field theory (DMFT), the magnitude of the local self-energy, \( |\text{Im}\Sigma(i\omega_n)| \), is enhanced with U until \( |\text{Im}\Sigma(i\omega_n)| |\omega_n \to 0 | \to -\infty \) as \( U \to U_{c2} \), the critical value for the Mott transition from the metallic phase. Quasiparticles disappear16 uniformly, \( Z_k = Z \to 0 \), and a Mott-Hubbard gap opens over the whole Fermi surface. This scenario is modified by non-local electron correlations in cluster extensions of DMFT such as cellular-DMFT (CDMFT)10,11,14,25, and the dynamical cluster approximation12 (DCA) which rather obtain anisotropic self-energy enhancements over the Fermi surface with larger \( |\text{Im}\Sigma_K(i\omega_n)| \) in the antinodal region around the coarse-grained momentum: \( K = (\pi, 0) \), than the nodal region for \( K = (\pi/2, \pi/2) \).12,13,15,20,23,26,27 This leads to a pseudogap12,14,26 in the electron spectra consistent with ARPES experiments on cuprates16. The origin of the pseudogap at \( (\pi, 0) \) in embedded cluster calculations has unambiguously been identified with spin fluctuations12,26 based on the fluctuation diagnostics approach18.

Electronic correlation effects at the Fermi surface are then signaled by large enhancements of \( |\text{Im}\Sigma_K(i\omega_n)| \) regardless of their origin. In order to fully characterize the ground state and excitations of the Hubbard model it is useful to quantify the strength of electron correlations not only at the Fermi surface but also in regions of the first Brillouin zone which are away from it. Recent work has reported large self-energy enhancements far away from the Fermi surface10,12,22,24 in hole doped Mott insulators. At low hole dopings and large-U, \( |\text{Im}\Sigma_{(\pi, \pi)}(i\omega_n)| \) can be
comparable or even larger than $|\text{Im} \Sigma_{(\pi,0)}(i\omega_n)|$ and much larger than $|\text{Im} \Sigma_{(0,0)}(i\omega_n)|$. Motivated by these recent findings we focus on self-energies far away from the Fermi surface. More specifically we would like to understand the origin of the large enhancement in $|\text{Im} \Sigma_{(\pi,\pi)}(i\omega_n)|$ found in embedded clusters: do self-energy enhancements away from the surface necessarily correspond to electronic correlation effects? Our study adds relevant information to previous works which have concentrated on electronic correlation effects at the Fermi surface. We consider the Hubbard model on a square lattice:

$$H = t \sum_{\langle ij \rangle} (c_{i\sigma}^\dagger c_{j\sigma} + c_{j\sigma}^\dagger c_{i\sigma}) + U \sum_i n_{i\uparrow} n_{i\downarrow} - \mu \sum_{i\sigma} n_{i\sigma},$$

where $t$ is the nearest neighbors hopping integral and $U$ the onsite Coulomb repulsion and $n_{i\sigma} = c_{i\sigma}^\dagger c_{i\sigma}$. We analyze the half-filled, $n = 1$, and the doped Hubbard model relevant to the cuprates.

The model is solved using DCA from which the different self-energies are obtained. We consider $N_c = 4,8$ clusters embedded in a self-consistent bath. The quantum cluster problem is solved using the Hirsch-Fye algorithm.

### III. Self-Energy Enhancements in DCA

We have performed DCA calculations on $N_c = 4,8$ clusters for $t = -1$, $U = 8$ and $\beta = 8$. The DCA self-energies obtained for $N_c = 4$ are shown in Fig. 1. For $n = 1$, the $K = (\pi,0)/(0,\pi)$ self-energies display divergent behavior, $\text{Im} \Sigma_{(\pi,0)}(i\omega_n)|_{\omega_n \to -\infty} \to -\infty$, due to the opening of a Mott-Hubbard gap. The $(0,0)$ and $(\pi,\pi)$ cluster self-energies are identical due to particle-hole symmetry and at low frequencies: $\text{Im} \Sigma_{(0,0)/(\pi,\pi)}(i\omega_n)|_{\omega_n \to 0} \to 0$. Under weak hole doping, $n = 0.94$, such divergence disappears and $\text{Im} \Sigma_{(0,0)/(\pi,\pi)}(i\omega_n)|_{\omega_n \to 0} \to 0$ as $\omega_n \to 0$. On the other hand, $(0,0)$ and $(\pi,\pi)$ become inequivalent so that $|\text{Im} \Sigma_{(\pi,\pi)}(i\omega_n)| > |\text{Im} \Sigma_{(0,0)}(i\omega_n)|$. Such asymmetry is robust against an increase in the size of the cluster as shown in Fig. 1(c) and d) for $N_c = 8$.

In order to elucidate the origin of the differences between the $(0,0)$ and $(\pi,\pi)$ self-energies we first analyze the behavior of sector populations. In Fig. 2 the sector populations per spin, $n_{K,\sigma}$, in a half-filled system, $n = 1$, are compared with the low hole doped system: $n = 0.94$. As expected, the doped holes mainly populate the $(\pi,0)/(0,\pi)$ sectors: $n_{(\pi,0),\sigma} = n_{(0,\pi),\sigma} < 0.5 $ since these are closest to the Fermi energy. On the other hand the $n_{(\pi,\pi),\sigma}$ and $n_{(0,0),\sigma}$ populations are weakly affected by doping as shown in Fig. 2. Since these populations satisfy: $n_{(\pi,\pi),\sigma} \approx 1 - n_{(0,0),\sigma}$, we would expect that $K = (\pi,0)$ and $K = (\pi,\pi)$ behave in a similar way under hole doping as for $n = 1$ shown in Fig. 1 a).

Hence, from the values of the $n_{(\pi,\pi)}$ and $n_{(0,0)}$ populations under hole doping we would have expected that for $n = 0.94$, $\text{Im} \Sigma_{(0,0)}(i\omega_n) \approx \text{Im} \Sigma_{(\pi,\pi)}(i\omega_n)$ in contrast to the DCA results plotted in Figs. 1 b)-d). We note that for electron doping the behavior of the $(0,0)$ and $(\pi,\pi)$ self-energies is switched so that (not shown): $|\text{Im} \Sigma_{(0,0)}(i\omega_n)| > |\text{Im} \Sigma_{(\pi,\pi)}(i\omega_n)|$.

Changes in ground state properties can be monitored by evaluating correlation functions, $C_{K,\sigma,K',\sigma'} = \langle n_{K,\sigma} n_{K',\sigma'} \rangle - \langle n_{K,\sigma} \rangle \langle n_{K',\sigma'} \rangle$. For instance, the opening of a pseudogap in the spectral function has been related with the formation of an RVB state in the cluster through the dependence of $C_{K,\sigma,K',\sigma'}$ on $U$. The formation of a RVB state in the cluster is signaled by $C_{K,\sigma,K',\sigma'} > 0$ with $K = K' = (\pi,0)$ or $(0,\pi)$. Here, we analyze the behavior of $C_{K,\sigma,K',\sigma'}$ involving the $(0,0)$ and $(\pi,\pi)$ sectors as shown in Fig. 3 for $n = 1$ and $n = 0.94$. For $n = 1$, $C_{(\pi,0),\uparrow,(\pi,0),\downarrow} = C_{(\pi,0),\uparrow,(\pi,\pi),\downarrow} = C_{(\pi,0),\uparrow,(0,\pi),\downarrow}$ as expected from particle-hole symmetry. At low hole dopings, this behavior is modified so that $(0,0)$ and $(\pi,\pi)$ become inequivalent since $C_{(\pi,0),\uparrow,(\pi,\pi),\downarrow} > C_{(0,0),\uparrow,(0,\pi),\downarrow}$ and $C_{(\pi,0),\uparrow,(0,\pi),\downarrow} > C_{(\pi,\pi),\uparrow,(0,\pi),\downarrow}$. However, these differences seem to be too small to explain the corresponding significant differences between the $(0,0)$ and $(\pi,\pi)$ self-energies shown in Fig.
The different magnitudes of the (0, 0) and (π, π) self-energies could arise from the different coupling strengths of the two sectors to the bath. We now analyze the doping dependence of the bath-cluster hybridization functions, \( \Gamma_{K}(i\omega_{n}) \) for \( K = (0,0) \) and \( (\pi,\pi) \). These can be obtained from

\[
\text{Im}\Gamma_{K}(i\omega_{n}) = \frac{\text{Im}G_{0K}(i\omega_{n})}{\text{Re}G_{0K}(i\omega_{n})^2 + |\text{Im}G_{0K}(i\omega_{n})|^2 + \omega_{n}},
\]

with \( G_{0K}(i\omega_{n}) \) the cluster excluded Greens function.

In Table I DCA results for \( \text{Im}\Gamma_{K}(i\omega_{n}) \) on half-filled \( N_c = 4 \) clusters, \( n = 1 \), are compared with hole doped clusters with \( n = 0.94 \). At half-filling both (0, 0) and (π, π) sectors are coupled with identical strengths due to particle-hole symmetry. However, for hole doping the (0, 0) hybridization is stronger than the (π, π) hybridization as shown in Table I. The asymmetry in \( \text{Im}\Gamma_{K}(i\pi/\beta) \) is related to the downward shift of \( \mu \) needed for hole doping the system. As \( \mu \) is decreased below \( U/2 \), \( \text{Im}G_{0(\pi,\pi)}(i\pi/\beta) \) is suppressed while \( \text{Im}G_{0(0,0)}(i\pi/\beta) \) is increased from their corresponding values at half-filling. Similar behavior of the hybridization functions is found for \( N_c = 8 \). The fact that \( |\text{Im}\Gamma_{(0,0)}(i\pi/\beta)| > |\text{Im}\Gamma_{(\pi,\pi)}(i\pi/\beta)| \) could naturally explain the larger magnitude of the (π, π) self-energy as compared to (0, 0) found in DCA.

The role played by the different (0, 0) and (π, π) hybridizations on the corresponding self-energies can be checked by fixing: \( \Gamma_{(0,0)}(i\omega_{n}) = \Gamma_{(\pi,\pi)}(i\omega_{n}) \). The results shown in Fig. I demonstrate how, even in this situation, \( |\text{Im}\Sigma_{(\pi,\pi)}(i\omega_{n})| > |\text{Im}\Sigma_{(0,0)}(i\omega_{n})| \) as in the actual self-consistent DCA calculations containing the different hybridizations of Table I. This result indicates that the larger (π, π) self-energy enhancement: \( |\text{Im}\Sigma_{(\pi,\pi)}(i\omega_{n})| > |\text{Im}\Sigma_{(0,0)}(i\omega_{n})| \), observed in DCA calculations is not due to the (π, π) hybridization being weaker than the (0, 0) hybridization.
Under an upward shift the behavior in the ter shown in Fig. 5 it is useful to calculate the spectral shift of bath-cluster hybridizations equal on a $N_c = 4$ cluster. The plot shows how the larger ($\pi, \pi$) self-energy enhancement: $|{\text{Im}}\Sigma_{\pi,\pi}(i\omega_n)| > |{\text{Im}}\Sigma_{\pi,\pi}(0,0)(i\omega_n)|$, shown in Fig. 4 is present even when we artificially fix: $\Gamma_{\pi,\pi}(i\omega_n) = \Gamma_{\pi,\pi}(i\omega_n)$. The parameters are: $t = -1$, $U = 8$, $\beta = 8$ and $n = 0.94$.

IV. SELF-ENERGY ENHANCEMENTS IN ISOLATED CLUSTERS

We have found above that the coupling to the bath plays a minor role on the different behavior of the $(0,0)$ and $(\pi, \pi)$ DCA self-energies. Motivated by this fact we have decoupled the cluster from the bath and have analyzed the self-energies of isolated $N_c = 4$ clusters. The hole (electron) doping in DCA calculations is simulated in an isolated cluster by just shifting the chemical potential: $\mu < U/2$ ($\mu > U/2$) with constant occupation, $n = 1$.

The imaginary part of the self-energy of an isolated $N_c = 4$ cluster is shown in Fig. 4. The $(\pi,0)$ and $(0,\pi)$ sectors display divergent behavior: $|{\text{Im}}\Sigma_{\pi,0}(0,\pi)(i\omega_n)| \to \infty$ as $\omega_n \to 0$ associated with the Mott-Hubbard gap. The imaginary part of the $(0,0)$ and $(\pi, \pi)$ self-energies coincide due to particle-hole symmetry when $\mu = U/2$. In Fig. 4 b) and c) we show how under a downward shift of $\mu$, they become different: $|{\text{Im}}\Sigma_{\pi,\pi}(i\omega_n)| > |{\text{Im}}\Sigma_{\pi,\pi}(0,0)(i\omega_n)|$ as found in DCA calculations (see Fig. 4). Under an upward shift the behavior in the $(0,0)$ and $(\pi, \pi)$ self-energies is switched so that: $|{\text{Im}}\Sigma_{\pi,\pi}(i\omega_n)| > |{\text{Im}}\Sigma_{\pi,\pi}(0,0)(i\omega_n)|$ as shown in Fig. 4 d).

In order to understand the different behavior of the $(0,0)$ and $(\pi, \pi)$ self-energies in the $N_c = 4$ isolated cluster shown in Fig. 4 it is useful to calculate the spectral functions: $A_K(\omega) = -\frac{1}{\pi} {\text{Im}}G_K(\omega + i\delta^+)$. In Fig. 4 we show $A_K(\omega)$ for: $\mu = U/2$ which is compared with the: $\mu = 2 < U/2$ case. The main difference between the two cases is a rigid shift of the spectra by $\mu$ with no associated redistribution of weight nor changes in the relative peak positions. This is in contrast to the larger enhancement: $|{\text{Im}}\Sigma_{\pi,\pi}(i\omega_n)| > |{\text{Im}}\Sigma_{\pi,\pi}(0,0)(i\omega_n)|$, for $\mu = 2 < U/2$ from which we would have naively interpreted that $(\pi, \pi)$ are more strongly correlated than $(0,0)$ electrons.

The two-peak structure of the spectral density, $A_K(\omega)$, shown in Fig. 4 is adequately described through a single-pole description of the self-energy:

$$\Sigma_K(i\omega_n) = -\frac{U}{2} = \frac{E_K}{i\omega_n + \Delta - E_K},$$

with $\Delta = \frac{U}{2} - \mu$ and the self-energy pole position $F_K$ and the constant $E_K$ independent of $\mu$. We use $n = 1$ in the Hartree contribution to the self-energy since we are at half-filling.

By introducing Eq. 5 in the Greens function of Eq. 4 and performing analytical continuation to the real axis:

$$A_K(\omega) = -\frac{1}{\pi} i {\text{Im}}G_K(\omega + i\delta^+) = \frac{\epsilon^M_K - F_K}{\epsilon^M_K - \epsilon^S_K} \delta(\omega - \Delta - \epsilon^M_K) + \frac{\epsilon^S_K - F_K}{\epsilon^M_K - \epsilon^S_K} \delta(\omega - \Delta - \epsilon^S_K),$$

with the location of the two peaks given by:

$$\epsilon^M_K = \frac{\epsilon_K + F_K}{2} - \sqrt{\left(\frac{\epsilon_K - F_K}{2}\right)^2 + E_K},$$

$$\epsilon^S_K = \frac{\epsilon_K + F_K}{2} + \sqrt{\left(\frac{\epsilon_K - F_K}{2}\right)^2 + E_K},$$

FIG. 4: (Color online) DCA self-energies with $(0,0)$ and $(\pi, \pi)$ bath-cluster hybridizations equal on a $N_c = 4$ cluster. The plot shows how the larger $(\pi, \pi)$ self-energy enhancement: $|{\text{Im}}\Sigma_{\pi,\pi}(i\omega_n)| > |{\text{Im}}\Sigma_{\pi,\pi}(0,0)(i\omega_n)|$, shown in Fig. 4 is present even when we artificially fix: $\Gamma_{\pi,\pi}(i\omega_n) = \Gamma_{\pi,\pi}(i\omega_n)$. The parameters are: $t = -1$, $U = 8$, $\beta = 8$ and $n = 0.94$.

FIG. 5: (Color online) Dependence of the imaginary part of the self-energy with $\mu$ in isolated $N_c = 4$ clusters. In a) the chemical potential is $\mu = U/2$ and there is particle-hole symmetry. As $\mu$ is shifted downwards for $\mu = 3$ in b) and $\mu = 2$ in c), then $|{\text{Im}}\Sigma_{\pi,\pi}(i\omega_n)|$ is enlarged with respect to $|{\text{Im}}\Sigma_{\pi,\pi}(0,0)(i\omega_n)|$. In d) we show a case $\mu > U/2$ in which $\mu$ is shifted upwards showing how the $(\pi, \pi)$ and the $(0,0)$ self-energies are just switched from the $\mu = 2$ situation. We note that for the $\mu$ values considered, there is no change in the cluster population which is half-filled: $n = 1$. The parameters used are $U = 8$ and $t = -2$ appropriate for the $N_c = 4$ embedded DCA calculations in Fig. 4.
isolated half-filled µ

However, if |ImΣ0) and satellite (S) peaks of the (0, 0) and (π, π) spectral functions are displayed for clarity. The occupation is at half-filling, n = 1, in all cases. The parameters used are U = 8 and t = −2 and Δ = U/2 − µ.

for K = (0, 0) whereas the M and S peak labels are switched for K = (π, π). The analytical expression for AK(ω) shows a two-peak structure as expected. Since the weights of the delta peaks in Eq. 6 are independent of Δ, it is evident that a shift in µ just leads to a rigid shift with no redistribution of the spectrum as found in the isolated cluster. The locations and weights of the two peaks occurring in A(0,0)(ω) and A(π,π)(ω) shown in Fig. 6 are faithfully reproduced by using: 

F(0,0) = −F(π,π) = 6.95, and E(0,0) = E(π,π) = 15.25 in the self-energy of Eq. 4. These parameters are obtained using the exact peak locations for εK and εS in Eq. 4.

We finally analyze how the shift in µ modifies the (0, 0) and (π, π) self-energies based on the single-pole form. The imaginary part of the self-energy in Eq. 5 reads:

\[ \text{Im} \Sigma_K(i\omega_n) = -\frac{E_K}{(\Delta + F_K)^2 + \omega_n^2}. \]  

In the symmetric case, µ = U/2 (Δ = 0), we have that: ImΣ(0,0)(ωn) = ImΣ(π,π)(ωn) since the self-energy poles are symmetrically located: |Δ + F(0,0)| = |Δ + F(π,π)|. However, if µ is shifted downwards (Δ = U/2 − µ > 0) then |Δ + F(π,π)| (|Δ + F(0,0)|) is suppressed (enhanced). From Eq. 5, this leads to an enhancement of |ImΣ(π,π)(iπ/β)| and a suppression of |ImΣ(0,0)(iπ/β)| with respect to their Δ = 0 values. Taking the FK, EK parameters used above to reproduce the spectral densities of Fig. 6 the self-energy behavior of Fig. 5 by which |ImΣ(π,π)(iπ/β)| > |ImΣ(0,0)(iπ/β)| under a downward shift of µ is correctly reproduced. We note that the single pole functional form of Eq. 5 has been previously introduced in single site DMFT24, for analyzing the atomic limit of the Mott insulator and more recently for interpreting particle-hole asymmetries in the electronic properties of the doped Hubbard model.25.

It is now worth analyzing the relevance of the single-pole self-energy functional of Eq. 5 to embedded clusters. Does the electronic structure of (0, 0) and (π, π) approximately behaves as in isolated clusters discussed above? To answer this question, we have obtained the DCA spectral functions at low hole dopings which are shown in Fig. 7. The (π, 0) spectral function displays a pseudogap at the chemical potential in agreement with previous works12,13,18,23,27 which we don’t discuss further here. The (0, 0) and (π, π) spectral functions are essentially gapped and mainly consist of a peak containing most of the spectral weight and a satellite peak with much more smaller weight. The main effect of the downward chemical potential shift, µ = 2.2 < U/2, which slightly hole dopes the system, n = 0.97, is to rigidly shift AK(ω) such that the main peak in A(0,0)(ω) becomes much closer to the chemical potential than the main peak in A(π,π)(ω). Hence, the overall behavior of A(0,0)(ω) and A(π,π)(ω) in DCA at low hole dopings is consistent with a rigid upward shift of the main spectral function structures similarly to the overall behavior found in isolated clusters. There are, however, some features in A(0,0)(ω) and A(π,π)(ω) intrinsic to DCA spectral functions not found in isolated clusters. The main peak in A(π,π)(ω) is more broad as compared to the main peak in A(0,0)(ω) when µ < U/2. This behavior is reasonable since there are more decay possibilities for electrons excited further away from the Fermi surface. Apart from this broadening there are some smaller structures occurring around the chemical potential in DCA which do not occur in isolated clusters. Due to the small weight of these features, the overall behavior of the (0, 0) and (π, π) self-energies is dominated by the position of the main peaks in A(0,0)(ω) and A(π,π)(ω) in consistent agreement with the isolated cluster analysis.

V. CONCLUSIONS AND DISCUSSION

In the present work we have analyzed the doping dependence of DCA self-energies, ΣK(iωn), for K away from the Fermi surface in a Hubbard model on the square lattice. For hole doping, we find larger enhancements of the self-energy magnitudes at the top (K = (π, π)) than at the bottom (K = (0, 0)) of the band which
would naively imply that electron correlation effects are stronger at $(\pi, \pi)$ than at $(0,0)$. However, the DCA populations satisfy: $n_{(0,0)} \approx 1 - n_{(\pi, \pi)}$, which would naively suggest similar electron correlation effects for $(0,0)$ and $(\pi, \pi)$ momenta.

In order to understand the origin of such self-energy differences we have first clarified the role played by bath-cluster hybridizations. Our DCA analysis shows that the self-energy difference: $|\Sigma_{(0,\pi)}(i\omega_n)| > |\Sigma_{(0,0)}(i\omega_n)|$ is not related to the $(0,0)$ hybridization being stronger than the $(\pi, \pi)$ hybridization. Indeed, such self-energy difference is also found in half-filled $(n = 1)$ isolated clusters but with $\mu < U/2$. Based on the equivalent spectral weight distributions at $K = (0,0)$ and $K = (\pi, \pi)$, we conclude that electron correlation effects acting at $K = (\pi, \pi)$ and at $(0,0)$ should be similar in the isolated cluster. This is in contrast to the conclusion we would have reached by just looking at the larger self-energy enhancement at $(\pi, \pi)$ compared to $(0,0)$.

Our analysis indicates that DCA self-energies at $k$-

points far away from the Fermi surface contain "apparent" enhancements which do not necessarily correspond to electronic correlation effects. These enhancements are a consequence of shifting the chemical potential from $\mu = U/2$ in a Mott insulator with the nearly constant occupation, $n \approx 1$. From an experimental point of view it would be interesting to compare the spectral functions of electrons close to the bottom of the band with electrons at the top of the band in hole doped cuprates. This would require probing unoccupied electronic states using angular resolved inverse photoemission (ARIPES) in combination with the more popular ARPES experiments probing occupied states.

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