scipost 202407 00001v1: Two-particle self-consistent approach for broken symmetry phases:

L. Del Re

September 11, 2024

Reply to Referee

The present work develops a two-particle self-consistent (TPSC) approach for broken symmetry phases and applies it to the three-dimensional Hubbard model. A benchmark against the Quantum Monte Carlo (QMC) data of Ref. [13] looks promising.

We thank the Referee for taking the time for reviewing our manuscript.

(I) - However, the manuscript clearly addresses a specialist audience. For example, there is a big leap from Eq. (1) to Eq. (2), i.e., as of the second page of the text, the author is talking to specialists only. Moving parts of Appendix A into the main text may help, but I expect that the manuscript will remain very technical nevertheless.

We understand the referee's comment and have decided to significantly improve the paper to reach a broader audience.

In response to the referee's suggestion, we have integrated the Appendix into the main text. This adjustment will help less experienced readers better understand the formalism of TPSC.

Additionally, based on the referee's advice and with the aim of reaching a broader audience, we have added two new figures (Figure 1 and Figure 2 in the revised version). These figures will aid readers in developing a better physical intuition about the method, even if they are not specialists.

Figure 1 shows a diagrammatic expression of the electron self-energy, which resembles the one seen in Hartree-Fock (HF) theory. We clarify at the beginning of Section 3 that this assumption simplifies the expressions for the Green's function and two-particle susceptibilities, which can then be analytically obtained using a formula similar to the Random Phase Approximation (RPA). In the following paragraph, we explain how TPSC, despite its similarities to mean-field approaches, goes beyond a simple mean-field description. Finally, we direct the reader to Figure 2, which presents a flow diagram summarizing the main steps and equations of the TPSC method.

The revised paper structure, now supported by the additional figures, provides a comprehensive overview of the method from the outset, making it accessible even to non-specialist readers.

(II)- Another point if the exponent $\beta = 1/2$ mentioned at the beginning of section 4. For the three-dimensional Heisenberg (O(3)) universality class, it should rather be $\beta \sim 0.369$ see

for example M. Campostrini et al., Phys. Rev. B 63, 214503 (2001) (also mentioned in Ref. [13]). To me, this is a strong indication that TPSC remains a mean-field theory after all, and I would find a related discussion plus references appropriate.

We thank the referee for the insightful comment, as it allows us to highlight a particularly subtle and important aspect of the theory. The exponent value for the order parameter $\beta = \frac{1}{2}$ might suggest that TPSC is a mean-field theory. However, the method actually belongs to a different universality class.

The critical exponents for TPSC, as well as those for other theories based on λ-Moriya corrections [see, for example, Phys. Rev. B 99, 045137], fall within the $O(N)$ universality class in the limit $N \to \infty$ [see Phys. Rev. B 53, 14236. This class is distinct from the mean-field theory, which is obtained in the limit of infinite spatial dimensions and is characterized by the exponents $\nu = \frac{1}{2}$, $\gamma = 1$, and $\beta = \frac{1}{2}$. In contrast, in three dimensions, the critical exponents for $O(\infty)$ are $\nu = 1$, $\gamma = 2$, and $\beta = \frac{1}{2}$.

In two dimensions, the critical temperature vanishes, as predicted by the Mermin-Wagner theorem, which also holds for the $O(3)$ case. This can be understood by considering the divergence of the sum in Eq.(14) at finite temperature in 2D within the broken symmetry phase, while it remains finite at $T = 0$, where the discrete sum over Matsubara frequencies is replaced by an integral over a continuous variable. Such a divergence prevents the system from ordering at finite temperature, and the correlation length of AF fluctuations grows exponentially as the temperature is reduced.

Our results for the critical exponent $\beta = \frac{1}{2}$ confirm that TPSC belongs to the $O(\infty)$ universality class, as expected. However, it is important to note that the critical exponents in TPSC are approximate and do not exactly match those of the $O(3)$ universality class, as the referee correctly pointed out.

We have added this discussion and the following Phys. Rev. B 65, 144520 reference in the revised version of the manuscript.

(III)-Furthermore, there are other (lattice) QMC investigation of the Néel transition in the half-filled three-dimensional Hubbard model. P. R. C. Kent et al., Phys. Rev. B 72, 060411(R) (2005) and S. Fuchs et al., Phys. Rev. Lett. 106, 030401 (2011) may be a good starting point to access further literature (actually, both references are also cited in Ref. [13]).

We thank the Referee for pointing out these important works, which we have now included in our references.

From our understanding, these studies estimate the Néel temperature and critical behavior of the Hubbard model from the symmetric side of the phase diagram, rather than within the broken symmetry phase. In this sense, they complement the works we originally cited, where the authors used Diagrammatic Monte Carlo to directly address the Hubbard model within the broken symmetry phase.

We believe that the inclusion of the works suggested by the Referee has enhanced the quality of our paper, and we have added them to our references accordingly.

Inspired by the referee's advice, we have expanded our literature review including additional references on Monte Carlo techniques. We have also added a recent preprint that presents a comprehensive Monte Carlo study of the Hubbard model in the symmetric phase.

(IV)-To conclude, I believe that there is some interest in this work such that it ultimately

merits publication in some form. However, my impression at least of the present manuscript is a minor technical progress that is of interest mainly to specialists. Therefore, I recommend transfer of a suitably revised version to SciPost Physics Core.

We recognize that the previous structure of the paper made it difficult to fully appreciate the potential of extending TPSC to treat magnetic phases of the Hubbard model.

For example, the approach we present here can be applied to models hosting altermagnetism [see e.g. Phys. Rev. X 12, 040501], a recently identified category of broken-symmetry phases. Group theory predictions suggest that a large number of such materials may exist in three dimensions [see Phys. Rev. X 14, 031039], where our algorithm can be readily applied. We anticipate that many new magnetic materials will soon be theoretically proposed and experimentally realized. Given the increasing complexity of these materials, it is crucial to develop efficient algorithms capable of capturing quantum fluctuations, which our method does.

Additionally, TPSC has already been successfully integrated with ab-initio methods, though only for symmetric phases [see Phys. Rev. Lett. 123, 256401]. This opens up exciting possibilities for extending our method to broken symmetry phases in combination with DFT (Density Functional Theory) for realistic electronic structure calculations.

We also demonstrated that TPSC is an effective tool for studying the amplitude (Higgs) mode, which is often elusive in most mean-field theories [see e.g. Phys. Rev. Lett. 115, 157002]. This paves the way for theoretical calculations of amplitude collective modes in altermagnets, providing not only a reference for future experimental investigations but also insights into fundamental questions –such as how the topological properties of altermagnets electronic structures are reflected in their collective modes.

We have included this discussion, along with new references, in the conclusions as part of the outlook, which should clarify the potential of the approach introduced in this work.

Extending TPSC to broken symmetry phases of the Hubbard model marks a significant step in describing complex magnetic materials. We believe this advancement will interest a broader audience, a point reflected in the expanded literature in this latest version of the manuscript. Scipost Physics offers a more suitable venue for sharing our findings, as it will help us reach a wider readership. Additionally, the paper's revised structure makes it more accessible to non-specialists.

Reuqested changes:

1-Move part of Appendix A to the beginning of section 3 in order to make the discussion more self-contained.

We followed the Referee's instructions and improved the paper structure inspired by the Referee's comments. [See our answer to point (I)]

2-Add discussion to the result $\beta = 1/2$ at the beginning of section 4 and cite relevant references, such as Phys. Rev. B 63, 214503 (2001) .

We added the discussion as well as the reference where the universality class of the Heisenberg models have been estimated. [See our answer to point (II)]

3-Mention further investigations of the half-filled three-dimensional Hubbard model such as P. R. C. Kent et al., Phys. Rev. B 72, 060411(R) (2005) and S. Fuchs et al., Phys. Rev. Lett. 106, 030401 (2011) and compare, e.g., the Néel temperature T_N with these.

We have added these references to our literature [see our answer to point (III)]. The comparison of methods for determining the Néel temperature using QMC techniques has already been addressed in previous works, which we have cited in the main text. Since the Néel temperature determined in our study aligns with the results of earlier research –where TPSC was used to estimate the critical behavior of the magnetic phases of the Hubbard model from the symmetric side of the phase diagram– we opted not to include this redundant information.

4-I liked the last three sentences of the Introduction (section 1) and hope that I have passed the test.

We have removed these sentences.

5-I believe that many equations would fit on one line, which would render the manuscript more readable. Examples: Eqs. (4), (10), (16), (18), (19), (21), (28), (37), (40), and (42). Same for a reduction of Eqs. (29) and (41) from three to two lines.

We reduced the number of lines as the Referee suggested.

6-Between Eqs. (5) and (6) there is an abbreviation "BSE" that has not been introduced.

In the revised version of the manuscript, the acronym BSE has now been defined.

7-Appendix C is quite short such that the cross-references generate unnecessary overhead. Why not move the content of the appendix to the appropriate place in section 4?

We agree with the referee's comment and have incorporated Appendix C into the Results section.

8-In panels (c) and (d) of Fig. 3, clarity would be improved of the labels "TPSC" and "RPA" had the same colors as the corresponding lines.

We have changed the colors of the lines that are now the same as the corresponding labels.

9-I can guess what the bar means, e.g., in the $\delta_{a\bar{b}}$ below Eq. (17), but I believe that an explanation would be helpful.

We have added an explanation of this notation to the text.

10-If the meaning of the bar in the $\overline{y}$ in Eq. (18) and below was specified, I have missed this description.

We decided to replace $\overline{y}$ with y' and explained in the text that y' is a repeated index and that summation is implied over repeated indices.

11-Below Eq. (41), there is a reference to "The last equation", but Eq. (41) actually is not an equation.

We fixed this.

12-Some preprint references are actually published. Maybe some appeared only after submission of the manuscript, but I still recommend an update. Specifically: [13] is published in Phys. Rev. Lett. 132, 246505 (2024). [19] is published in Phys. Rev. B 109, 045155 (2024). [28] is published in Phys. Rev. B 108, 075144 (2023). [34] is published in Phys. Rev. B 109, 075143 (2024).

We have updated all these references in the new version of the manuscript.

13-The URLs in Refs. [29,30] are redundant and could be omitted.

We removed the URLs from these references.

14-There are issues with the English text, such as: a) The manuscript mixes British and American English. For example, there are occurrences of "magnetisation" (British English) and "magnetization" (American English). I recommend that the author settles on one version and runs the manuscript through an appropriate spellchecker. b) There are several instances of "as following" which in my opinion should read "as follows". c) There is a duplicate "for for" in the caption of Fig. 3. Overall, I believe that the manuscript would benefit from careful proofreading, preferably also from somebody else than the author, and ideally a native English speaker.

We have addressed points (a), (b), and (c) and followed the referee's advice. We believe that the level of English in our paper has significantly improved.

Two-particle self-consistent approach for broken symmetry phases (DIFF)

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Abstract

Spontaneous symmetry breaking of interacting fermion systems constitutes a major challenge for many-body theory due to the proliferation of new independent scattering channels once absent or degenerate in the symmetric phase. One example is given by the ferro/antiferromagnetic broken symmetry phase (BSP) of the Hubbard model, where vertices in the spin-transverse and spin-longitudinal channels become independent with a consequent increase in the computational power for their calculation. Here we generalize the formalism of the non-perturbative Two-Particle-Self-Consistent method (TPSC) to treat broken SU(2) magnetic phases of the Hubbard model, providing with a efficient yet reliable method. We show that in the BSP, the sum-rule enforcement of susceptibilities must be accompanied by a modified gap equation resulting in a renormalisation of the order parameter, vertex corrections and the preservation of the gap-less feature of the Goldstone modes. We then apply the theory to the antiferromagnetic phase of the Hubbard model in the cubic lattice at half-filling. We compare our results of double occupancies and staggered magnetisation to the ones obtained using Diagrammatic Monte Carlo showing excellent quantitative agreement. We demonstrate how vertex corrections play a central role in lowering the Higgs resonance with respect to the quasi-particle excitation gap in the spin-longitudinal susceptibility, yielding a well visible Higgs-mode.

Contents

1 Introduction

The ~~characterization~~ characterisation of broken symmetry phases (BSP) in correlated quantum systems remains a formidable challenge for many-body theory. In fact, determining the precise ground state of spin Hamiltonians, such as the 3D-Heisenberg model with antiferromagnetic exchange, remains an open question to this day. Even ~~if the precise knowledge of the groundstate might remain out of reach~~though the precise ground state may remain elusive, it is possible to improve mean-field predicted groundstates , e.g. the Néel state, including quantum corrections encoded in the long-range and low-energy Goldstone modes [1, 2, 3, 4, 5], e.g. spin-waves in antiferromagnets [6].

The situation becomes richer when interacting electrons in solids ~~get~~ are strongly correlated. A minimal model to describe ~~correlated materials is given by~~ such materials is the Hubbard model [7], where electrons interact ~~via~~ through on-site Coulomb repulsion~~that enhances~~, enhancing electron localisation [8]. The theoretical challenge with strongly correlated BSP ~~consists in taking into account at the same time the~~ lies in simultaneously accounting for long-range fluctuations encoded in ~~the~~ Goldstone modes and the localisation of electrons.

Such an ambitious task could be achieved by employing cluster [9] or diagrammatic [10] extensions of Dynamical Mean Field Theory (DMFT) [11], ~~or Diagrammatic Monte Carlo (DiagMC) [12, 13]~~as well as Monte Carlo techniques [14, 15, 12, 16, 13, 17]. However, the inclusion of long-range modes for cluster theories would be limited by the maximum size of the cluster used in the calculations, even if clever clustering schemes that permit an optimal finite-size scaling analysis are available [18]. In diagrammatic approaches, the proliferation of independent vertex components [19, 20, 21, 22, 23, 24, 25], once absent or degenerate in the symmetric phase, strongly increases the computational power needed for their numerical evaluation.

Hence, it is of great interest to develop efficient algorithms ~~requiring less computational~~

~~resources but that at the same time are able to include~~ that require fewer computational resources while still accurately including correlation effects. In this ~~realm~~context, the Two-Particle-Self-Consistent (TPSC) approach [26, 27, 28, 29, 30, 31, 32, 33] has ~~been~~ proven to be a reliable and efficient method ~~to describe~~ for describing the physics of the Hubbard model in the weak-to-intermediate interaction regime. Given its reduced computational complexity, TPSC has already been successfully extended to multi-orbital models [30], interfaced with *ab-intio* calculations [31] and applied to non-equilibrium [32]. ~~However, TPSC formulations available today can only treat~~ However, current TPSC formulations are limited to symmetric phases, ~~which prevents the application of the theory~~ preventing their application to parameter regimes where materials ~~are found in BSP. Furthermore, since~~ exhibit broken symmetry phases (BSP). Additionally, because TPSC uses Moriya corrections to two-particle ~~propagators masses [34, 35, 36, 37, 38] for including~~ propagator masses [34, 35, 36, 37, 38] to include correlation effects, a straightforward generalisation of TPSC equations ~~could~~ might violate Goldstone's theorem~~introducing~~ , leading to an unphysical energy gap ~~to~~ in the Goldstone modes. In this work, we ~~show how to properly~~ extend the TPSC formalism to ~~the case of~~ handle spontaneous symmetry breaking ~~, which correctly preserves~~ while correctly preserving the Goldstone modes.

We apply the new formulation to the antiferromagnetic phase of the three-dimensional Hubbard model ~~in the~~ on a cubic lattice. ~~We compared our results with DiagMC [13] showing an~~ Our results show excellent quantitative agreement ~~for~~ with Diagrammatic Monte Carlo (DiagMC) [13] across a wide range of interaction values. We ~~show that the degree of correlation is reduced by decreasing temperature~~ demonstrate that as the temperature decreases from the critical value, ~~extending the range of applicability of the theoryto higher values of the interactions~~ the degree of correlation is reduced, which extends the theory's applicability to higher interaction values deep in the ~~BSP. We demonstrate how symmetry breaking implies~~ broken symmetry phase. Additionally, we show that symmetry breaking leads to a differentiation of vertex corrections in ~~different scattering channels, which play~~ various scattering channels. This differentiation plays a central role in lowering the Higgs resonance ~~with respect~~ relative to the quasi-particle excitation gap in the spin-longitudinal susceptibility, ~~yielding a well visible Higgs-mode. The stage is yours. Write your article here.The bulk of the paper should be clearly divided into sections with short descriptive titles, including an introduction and a conclusion.~~resulting in a clearly distinguishable Higgs mode.

The manuscript is organised as follows: in Sec. 2 we introduce the Hubbard model and establish the notation; Sec.3 describes the method and explains how two-particle self-consistency can be achieved in magnetic broken symmetry phases while preserving the Goldstone modes; in Sec.4 we show the numerical data of the order parameter and double occupancies comparing them with DiagMC, and we also show how TPSC is able to capture the elusive amplitude (Higgs) mode in the susceptibility spectra; in Sec.5 we provide our conclusions and outlook; in Appendix A we discuss some technical details relative to the derivation of the effective irreducible vertices; in Appendix B we present the derivation of the Bethe-Salpeter equations; in Appendix C we show the steps needed to obtain the corrected one-loop self-energy.

2 The model

In this work we will explicitly consider the single band Hubbard model in the cubic lattice,

$$H = -t \sum_{\langle ij \rangle \sigma} c^{\dagger}_{i\sigma} c_{j\sigma} + U \sum_{i} \hat{n}_{i\uparrow} \hat{n}_{i\downarrow}, \tag{1}$$

where t is the electronic hopping amplitude between ~~nearest-neighbors~~ nearest-neighbours and U is the local Coulomb repulsion. In the case of the AF phase, the system loses the full translational symmetry of the original cubic lattice and it is useful to introduce the sub-lattice index $a = A, B$ for specifying whether the fermionic field $c^{\dagger}_{ia\sigma}$ is evaluated at one site belonging to the sub-lattice A or B. Therefore, it is useful to introduce the generalised multi-flavor indices α, which for example coincide with $\alpha = (a, \sigma)$ containing both sub-lattice (a) and spin (σ) indices in the AF or to spin indices in the FM case. Then, we can rewrite the Hubbard Hamiltonian in the following form:

$$H = \sum_{\langle ij \rangle} \sum_{\alpha\beta} c^{\dagger}_{i\alpha} \mathcal{H}^{\alpha\beta} c_{j\beta} + \frac{1}{2} \sum_{i} \sum_{\alpha\beta} U_{\alpha\beta}\, \hat{n}_{i\alpha} \hat{n}_{i\beta}. \tag{2}$$

In the case of FM, we have that $\mathcal{H}^{\alpha\beta} = -t\delta_{\alpha\beta}$ and $U_{\alpha\beta} = \delta_{\alpha\bar{\beta}} U$, whereas for the AF case we have $\mathcal{H}^{\alpha\beta} = -t\delta_{\sigma\sigma'}\delta_{a\bar{b}}$ and $U_{\alpha\beta} = \delta_{\sigma\bar{\sigma}'}\delta_{ab} U$, where $\bar{\ell}$ denote the opposite of index ℓ, referring to the complementary spin or sub-lattice index (e.g., if ℓ is spin-up or sub-lattice A, then $\bar{\ell}$ is spin-down or sub-lattice B).

3 The method

The Two-Particle Self-Consistent (TPSC) method requires relatively low computational power and achieves its efficiency through a series of approximations, which we will examine in detail in this section. In practice, the self-energy is approximated in a form similar to that used in Hartree-Fock (HF) (see Figure 1). This assumption simplifies the expressions for the Green's function and two-particle susceptibilities, which can then be analytically obtained using a formula akin to the Random Phase Approximation (RPA).

However, unlike in HF and RPA, the vertex in the self-energy and susceptibility diagrams is represented not by the bare interaction but by an effective vertex. This effective vertex includes a renormalisation factor that depends on the double occupations, which are determined by imposing an exact sum rule on the susceptibility in the spin-transverse channel. This sum rule complements the 'usual' gap equation for the order parameter by coupling it to the double occupancies, which, unlike in HF, are determined self-consistently.

While we will provide an explicit derivation of all the equations, readers primarily interested in the results may refer to Figure 2, which presents a flow diagram summarizing the main steps and equations of the TPSC method.

3.1 The TPSC ansatz

The ~~starting point~~ core of TPSC consists in ~~assuming~~ finding an approximate form ~~of~~ for the electron self-energy ~~. In particular, we assume the following ansatz for the self-energy:~~ from

which one can construct a conserving approximation in the Baym-Kadanoff sense [39, 40]. In order to do so, one can start from the equation of motion that reads:

$$\Sigma^{ab}_{\sigma\sigma'}{}^{\alpha\gamma}(x,y') = \delta G^{\gamma,\beta}(x-y y', y)\delta_{ab}\,\lambda = U\delta_{\sigma\sigma'}n_{a\bar{\sigma}} - \delta_{\sigma\bar{\sigma}'}s^{\sigma\bar{\sigma}}_{a}{}_{\alpha\gamma}G^{(2)\beta\alpha}_{\gamma\gamma}(y, x+0^-, x+0^+, x), \quad (3)$$

where ~~$x = (\mathbf{R}, \tau)$ is a quadrivector~~ $G^{\alpha\beta}(x,y) = -T_t\left\langle c_\alpha(x)c^\dagger_\beta(y)\right\rangle$ is the Green's function, with $x = (R_i, \tau_i)$ being a four-vector containing the lattice ~~coordinates ($\mathbf{R}$)~~coordinate R_i and the imaginary time τ_i, $c_\alpha(x) = e^{H\tau_i}c_{i\alpha}e^{-H\tau_i}$, $\Sigma^{\alpha\beta}(x,y)$ is the electronic self-energy, and $G^{(2)\alpha\beta}_{\gamma\delta}(x_1,x_2,x_3,x_4) = T_\tau\left\langle c^\dagger_\alpha(x_1)c_\beta(x_2)c^\dagger_\gamma(x_3)c_\delta(x_4)\right\rangle$ represents the two-particle Green's function. In Eq.(3), a summation is intended for the repeated indices γ and y'. Due to the presence of $G^{(2)}$, Eq.(3) is not closed for the self-energy and ~~the imaginary time (τ), $n_{a\sigma} = \left\langle c^\dagger_{ia\sigma}c_{ia\sigma}\right\rangle$, $s^{\sigma\sigma'}_a = \left\langle c^\dagger_{ia\sigma'}c_{ia\sigma}\right\rangle$, and :~~ single-particle Green's function, and in order to obtain an explicit expression for Σ further approximations must be carried on. In mean-field theory the two-particle Green's function is replaced by its disconnected part, that is a valid approximation only at weak coupling. In TPSC [26, 27, 29, 30], in order to take into account of correlation effects, and at the same time to reduce the complexity of Eq.(3), the following assumption is considered:

$$\Sigma^{\alpha\gamma}(x,y')G^{\gamma,\beta}(y',y) \sim \lambda^{\alpha\gamma}U_{\alpha\gamma}\left[G^{\alpha\beta}(x,y)n_\gamma - s^{\alpha\gamma}G^{\gamma\beta}(x,y)\right], \quad (4)$$

where $n_\alpha = \langle\hat{n}_{i\alpha}\rangle$, $s^{\alpha\beta} = \left\langle c^\dagger_{i\beta}c_{i\alpha}\right\rangle$, and $\lambda^{\alpha\beta}$ is an extra-coefficient that must be determined self-consistently and contains correlation effects. When $\lambda^{\alpha\beta} = 1$ mean-field theory is recovered. The parameter λ can be determined by requiring that the equal-time/position limit of Eq.(3), i.e. $y = x + 0^{++}$, is preserved exactly when $\beta = \alpha$ [1], by imposing:

$$\lambda^{\alpha\gamma} = \frac{\langle\hat{n}_\alpha\,\hat{n}_\gamma\rangle}{n_\alpha n_\gamma - s^{\alpha\gamma}s^{\alpha\gamma}}\,. \quad (5)$$

From Eq.(4), we can isolate the self-energy that reads:

$$\Sigma^{\alpha\beta}(x-y) = \delta(x-y)\left(\delta_{\alpha\beta}\lambda^{\alpha\gamma}U_{\alpha\gamma}n_\gamma - \lambda^{\alpha\beta}U_{\alpha\beta}s^{\alpha\beta}\right). \quad (6)$$

In the FM/AF phase of the Hubbard model, the expression for the λ parameter simplifies as follows:

$$\lambda = \frac{\langle\hat{n}_\uparrow\hat{n}_\downarrow\rangle}{n_\uparrow n_\downarrow}, \quad (7)$$

which is ~~the same prefactor appearing~~ identical to the one in the paramagnetic case ~~[27, 29, 41]. In the next section, we will see how the parameter~~ [27, 29]. Since in the FM/AF phase Eq.(7)

[1] In the case of the AF phase that we address in this work, spin conservation implies that $\langle c^\dagger_\alpha\hat{n}_\gamma c_\beta\rangle = 0$ at zero field, when $\alpha \neq \beta$, and therefore we shall introduce the λ-correction only for the two-particle Green's functions that do not vanish in the limit of zero external field.

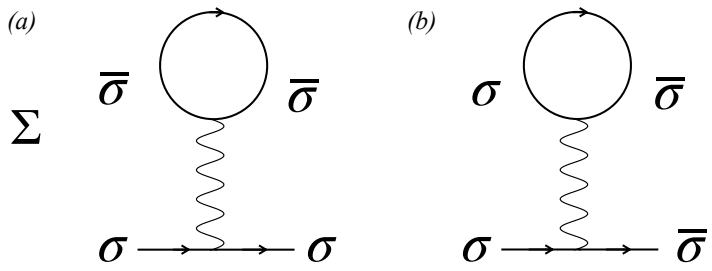

Figure 1: Diagrammatic representation of the diagonal (a) and off-diagonal (b) components of the self-energy, as analytically expressed in Eq.(22). The wiggly line represents the effective vertex $\Gamma_{\overline{\uparrow\downarrow}}$, which must be obtained self-consistently along with the order parameter (encoded in the Green's function, represented by the thick continuous line). Both are determined through the simultaneous solution of Eqs.(13,14).

does not depend on the spin/sub-lattice indices, we omitted those indices in the expression of λ, ~~or equivalently the double occupancies, can be evaluated self-consistently together with the order parameter, by imposing specific sum rules of physical susceptibilities.~~ and therefore we need to optimize only one parameter even within the broken-symmetry phases under scrutiny. Hence, in the AF phase the expression of the self-energy can be written as:

$$\Sigma^{ab}_{\sigma\sigma'}(x,y) = \delta(x-y)\delta_{ab}U_{\overline{\uparrow\downarrow}}\left(\delta_{\sigma\sigma'}n_{a\bar{\sigma}} - \delta_{\sigma\bar{\sigma}'}s^{\sigma\bar{\sigma}}_{a}\right), \tag{8}$$

where $U_{\overline{\uparrow\downarrow}} = \lambda U$, a, b are sub-lattice indices. We notice that in ~~the~~ a AF phase the ~~off-diagonal~~ components of the self-energy off-diagonal in the spin indices should vanish, i.e. $s^{\sigma\bar{\sigma}} = 0$. However, it is useful to keep those terms in the expression of the self-energy for the derivation of the Bethe-Salpeter equation in the spin-transverse channel. Therefore, we will consider the presence of an external field that breaks spin-conservation ~~to~~ and eventually compute the functional derivatives of Σ with respect to the off-diagonal component of the propagator ~~evaluated~~ in the limit of a vanishing field.

~~In Appendix ?? we show that the expression for Σ in Eq.(??) can be derived starting from the equation of motion and by performing an approximation at the two-particle level that preserves an exact constraint in the limit of equal times and positions.~~

In order to obtain self-consistency at the two-particle level, we have to calculate physical susceptibilities and therefore we need the knowledge of the irreducible vertex function Γ, which is obtained by carrying the functional derivative of Σ with respect to G, i.e. $\Gamma(1,2,3,4) = \frac{\delta\Sigma(2,1)}{\delta G(3,4)}$ [42].

In the FM/AF phases the original SU(2) symmetry of the Hubbard Hamiltonian is spontaneously broken and the two independent scattering channels to be considered are the spin-transverse and spin-longitudinal channels [22].

3.1.1 Spin-transverse channel

The vertex function in the spin-transverse channel is defined as:

$$\Gamma^{abcd}_{\overline{\uparrow\downarrow}}(x_1,x_2,x_3,x_4) = \frac{\delta\Sigma^{ba}_{\downarrow\uparrow}(x_2,x_1)}{\delta G^{cd}_{\downarrow\uparrow}(x_3,x_4)} = -\lambda U\delta_{ab}\delta_{ac}\delta_{ad}\delta(x_1-x_2)\delta(x_1-x_3)\delta(x_1-x_4), \tag{9}$$

where we used Eq.(??22) and the fact that $s_a^{\sigma\bar{\sigma}} = G^{aa}_{\sigma\bar{\sigma}}(x, x+0^-)$ [2].

Let us now define the physical susceptibility in the transverse-spin channel:

$$\chi^{ab}_{\sigma\bar{\sigma}}(x_1, x_2) = T_\tau \left\langle S_a^{\sigma\bar{\sigma}}(x_1) S_b^{\bar{\sigma}\sigma}(x_2) \right\rangle , \tag{10}$$

where $S_a^{\sigma\sigma'}(x) = e^{H\tau} c^\dagger_{ia\sigma} c_{ia\sigma'} e^{-H\tau}$, with $x = (R_i, \tau)$. Since the vertex function in Eq.(??22) is local and static, the ~~BSE (~~Bethe-Salpeter equation (BSE) [see Appendix B for the derivation~~)~~] for the physical susceptibilities is similar to the one obtained in RPA [22] and reads:

$$\bar{\bar{\chi}}^{-1}_{\sigma\bar{\sigma}}(q) = \bar{\bar{\chi}}^{-1}_{0,\sigma\bar{\sigma}}(q) + \bar{\bar{\Gamma}}_{\sigma\bar{\sigma}}, \tag{11}$$

where we used the double bar to indicate 2×2 matrices, $q = (i\omega_n, \mathbf{q})$ with $\omega_n = 2\pi n/\beta$ and $\mathbf{q}$ being respectively the bosonic Matsubara frequency and crystalline exchanged momentum, $\bar{\bar{\chi}}_{\sigma\bar{\sigma}}(q)$ is given by the Fourier transform of the susceptibility defined in Eq.(10), $\bar{\bar{\Gamma}}_{\sigma\bar{\sigma}} = -\lambda U \, \mathbb{I}_{2\times 2}$ and $\chi^{ab}_{0,\sigma\bar{\sigma}} = -\frac{1}{V\beta} \sum_k G^{ab}_\sigma(k) G^{ab}_{\bar{\sigma}}(k+q)$. The Green's function is obtained using the Dyson equation and reads:

$$\bar{\bar{G}}^{-1}_\sigma(k) = \epsilon_{\mathbf{k}} \, \sigma^x + [i\nu + \mu - \frac{\Gamma_{\overline{\uparrow\downarrow}}}{2}(n + \sigma m)] \mathbb{I}_{2\times 2}, \tag{12}$$

where n is the electron density and $m = n_{A\uparrow} - n_{A\downarrow}$ is the staggered magnetisation.

In order to univocally determine single-particle and two-particle properties, we have to solve a set of self-consistent equations that will allow us to find the chemical potential, staggered ~~magnetization~~ magnetisation and double occupancies (μ, m, $\langle \hat{n}_\uparrow \hat{n}_\downarrow \rangle$) as a function of the electron density, on-site interaction and temperature. In this work we will specialize in the case of the three-dimensional cubic lattice at half-filling , i.e. $n = 1$, that corresponds to fixing the chemical potential to $\mu = \frac{\Gamma_{\overline{\uparrow\downarrow}}}{2}$.

Since the self-energy is static and local, the gap equation for the order parameter is similar to the one obtained in mean-field theory and is given by following expression:

$$\frac{1}{(2\pi)^3} \int_{BZ} d\mathbf{k} \, \frac{|\Gamma_{\overline{\uparrow\downarrow}}|}{2E_{\mathbf{k}}} \tanh\left(\frac{\beta E_{\mathbf{k}}}{2}\right) = 1, \tag{13}$$

where $E_{\mathbf{k}} = \sqrt{\epsilon^2_{\mathbf{k}} + \left(\frac{m\,\Gamma_{\overline{\uparrow\downarrow}}}{2}\right)^2}$, with $\epsilon_{\mathbf{k}} = -2t\,[\cos(k_x) + \cos(k_y) + \cos(k_z)]$. Differently from mean-field theory however, the order parameter is not univocally determined by the gap equation, because the double occupancies, appearing in Eq.(13), are still unknown.

As a direct consequence of its definition in Eq.(10), the susceptibility in the transverse channel assumes the following limiting value $\sum_\sigma \chi^{aa}_{\sigma\bar{\sigma}}(x, x+0^-) = n - 2\langle \hat{n}_\uparrow \hat{n}_\downarrow \rangle$, which implies the following sum rule for its Fourier transform:

$$\frac{1}{\beta(2\pi)^3} \sum_{\omega_n \sigma} \int_{\mathrm{BZ}} d\mathbf{q} \, \chi^{aa}_{\sigma\bar{\sigma}}(q) = n - 2\langle \hat{n}_\uparrow \hat{n}_\downarrow \rangle . \tag{14}$$

Hence, Eqs.(13,14) provide with a closed set of equations that must be solved self-consistently in order to determine the order parameter and the double occupancies.

[2] We used the overline symbol, i.e. $\overline{\uparrow\downarrow}$, to distinguish this vertex component from those belonging to the longitudinal spin channel, that are defined in the next paragraphs.

3.2 Spin-longitudinal channel

The irreducible vertex function in the longitudinal-spin channel reads:

$$\Gamma^{abcd}_{\sigma\sigma'}(x_1,x_2,x_3,x_4) = \frac{\delta\Sigma^{ba}_{\sigma\sigma}(x_2,x_1)}{\delta G^{cd}_{\sigma'\sigma'}(x_3,x_4)} \sim U_{\sigma\sigma'}\delta_{\sigma\bar{\sigma}'}\delta_{ab}\delta_{ac}\delta_{ad}\,\delta(x_1-x_2)\delta(x_1-x_3)\delta(x_1-x_4). \tag{15}$$

Differently from Eq.(9) which is an exact equality, a further approximation, similar to the one performed in the charge channel in the paramagnetic phase [26, 29], is needed to write Eq.(15) in its final form (see Appendix A).

Let us define the susceptibilities in the spin-longitudinal channel:

$$\chi^{ab}_{\sigma\sigma'}(x_1,x_2) = T_\tau \langle n_{a\sigma}(x_1) n_{b\sigma'}(x_2)\rangle - \langle n_{a\sigma}\rangle \langle n_{b\sigma'}\rangle \tag{16}$$

Given the local and static form of the vertex function in Eq.(15), the expression of the susceptibilities in the charge and spin-longitudinal channel, in presence of particle-hole symmetry [3], can be written as ~~following~~follows:

$$\chi_z(q) = \frac{\chi_{0,\parallel}(q)}{1-\Gamma_z\chi_{0,\parallel}(q)} \tag{17}$$

$$\chi_\rho(q) = \frac{\chi_{0,\parallel}(q)}{1+\Gamma_\rho\chi_{0,\parallel}(q)}, \tag{18}$$

where $\chi_z = \frac{1}{2}\sum_{ab\sigma\sigma'}(-1)^{a+b+\sigma+\sigma'}\chi^{ab}_{\sigma\sigma'}$, $\chi_\rho = \frac{1}{2}\sum_{ab\sigma\sigma'}\chi^{ab}_{\sigma\sigma'}$, $\Gamma_z = \frac{1}{2}\sum_{\sigma\sigma'}(-1)^{\sigma+\sigma'}\Gamma_{\sigma\sigma'}$, $\Gamma_\rho = \frac{1}{2}\sum_{\sigma\sigma'}\Gamma_{\sigma\sigma'}$, $\chi_{0,\parallel} = -\frac{1}{2V\beta}\sum_{k\sigma b} G^{Ab}_\sigma(k)G^{bA}_\sigma(k+q)$. Analogously for the spin-transverse channel, we can determine the value of the vertices Γ_z and Γ_ρ by imposing the following sum rule for the longitudinal channel susceptibilities:

$$\frac{2}{\beta(2\pi)^3}\sum_{\omega_n}\int_{\mathrm{BZ}} d\mathbf{q}\,\chi_z(q) = n - 2\langle\hat{n}_\uparrow\hat{n}_\downarrow\rangle - m^2 \tag{19}$$

$$\frac{2}{\beta(2\pi)^3}\sum_{\omega_n}\int_{\mathrm{BZ}} d\mathbf{q}\,\chi_\rho(q) = n + 2\langle\hat{n}_\uparrow\hat{n}_\downarrow\rangle - n^2. \tag{20}$$

Since Eqs.(13,14) are a set of closed equations, Eqs.(19,20) can be solved separately once the values of m and $\langle\hat{n}_\uparrow\hat{n}_\downarrow\rangle$ have been self-consitently obtained from the spin-transverse channel.

3.3 Improved one-loop self-energy

In TPSC it is possible to obtain an improved self-energy that, differently from the one appearing in Eq.(~~??~~22), depends on both momenta and frequency. This can be achieved by computing the TPSC vertices and susceptbilities and using them as input for the equation of motion [27, 45]. Extending this procedure to the broken symmetry phase we obtain the following expression for the improved self-energy:

$$\Sigma^{ab}_\sigma(k) - Un_{a\bar{\sigma}} = -\frac{U}{2V\beta}\sum_q G^{ab}_{\bar{\sigma}}(k+q)\Gamma^a_{\bar{\sigma}\bar{\sigma}}\chi^{ab}_{\bar{\sigma}\bar{\sigma}}(q) + \frac{U}{2V\beta}\sum_{q\sigma_1} G^{ab}_\sigma(k+q)\Gamma^a_{\sigma\sigma_1}\chi^{ab}_{\sigma_1\bar{\sigma}}(q), \tag{21}$$

where $G^{ab}_\sigma(k)$ is given by Eq.(12). In appendix C we show the derivation of Eq.(21).

[3] In general, the charge and longitudinal-spin channels interact via a mixed terms $\chi_{\bar{z}\rho}$ [43, 44] that vanishes only in presence of particle-hole symmetry [22]

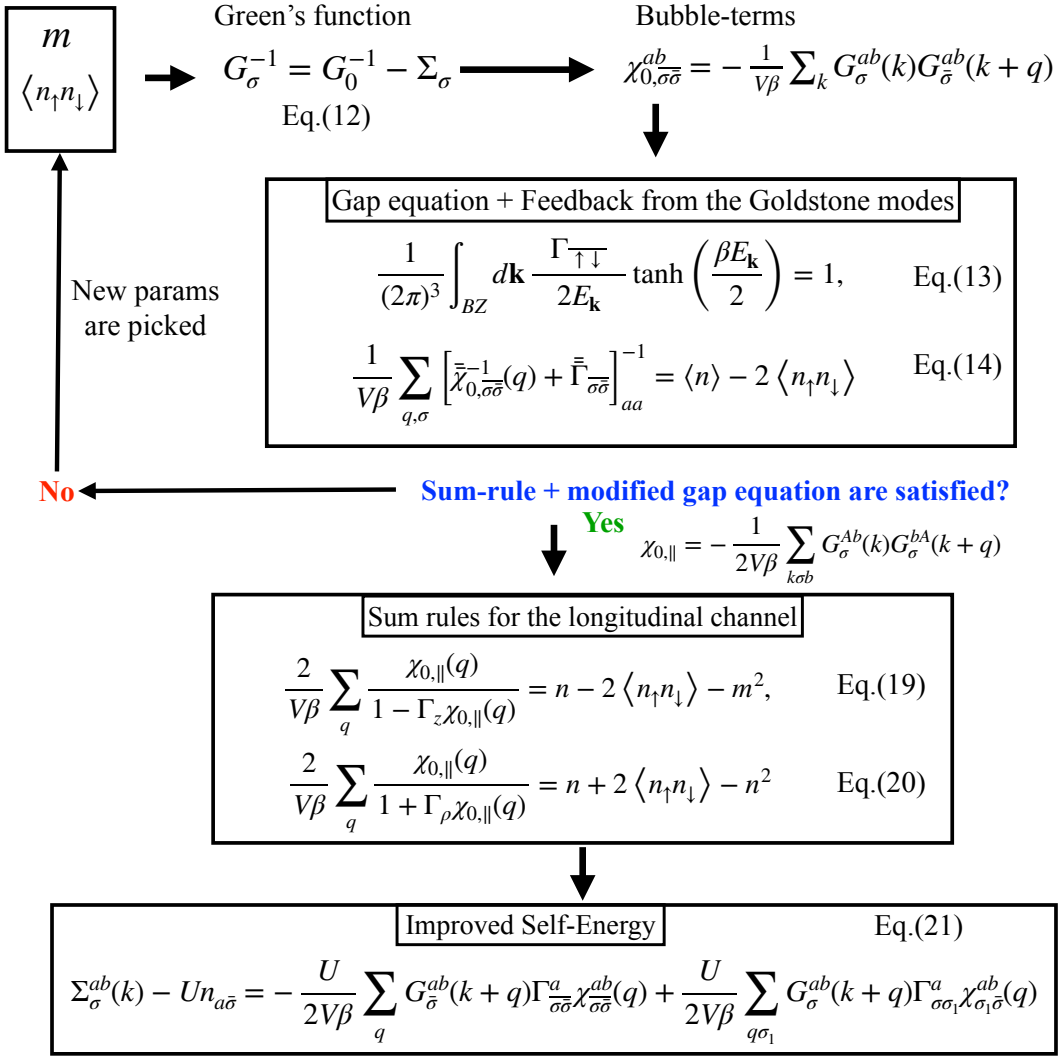

Figure 2: **Work-flow of the Two-Particle Self-Consistent (TPSC) method for antiferromagnetic phases of the Hubbard model.** The first box at the top shows how the staggered magnetisation $m = n_{A\uparrow} - n_{A\downarrow}$ and double occupancies are obtained self-consistently by solving the gap equation [Eq. (13)] and the sum rule [Eq. (14)] for the spin-transverse channel, where Goldstone modes appear. An initial guess for m and $\langle n_\uparrow n_\downarrow \rangle$ is used to calculate the Green's function and susceptibility. If Eqs. (13) and (14) are not satisfied, a minimisation routine adjusts the values. Once satisfied, the next step is to find the renormalized vertices in the spin-longitudinal channel by enforcing Eqs. (19) and (20), shown in the middle box. With all renormalized interactions in the different channels, the improved electron self-energy can finally be computed using Eq. (21) displayed in the box at the bottom.

4 Numerical results

~~(a) Staggered magnetisation m as a function of T for three different values of $U/t = 3, 4, 5$. Dashed lines are best fits of the function $\alpha|T - T_c|^{1/2}$ close to T_c. (b) λ parameter as a function of the reduced temperature $\theta = \frac{T-T_c}{T_c}$ for the three different values of the on-site interaction. (c) Double occupancies $D = \langle \hat{n}_\uparrow \hat{n}_\downarrow \rangle$ as a function of θ for the three different U values. (d) Magnetisation and double occupancies as a function of U for $T/t = 1/10$. TPSC data (open symbols) are compared to the DiagMC results (filled symbols) adapted from Ref. [13]. The black dashed line is the magnetisation curve obtained using Hartree-Fock.~~

Fig. (3-a) shows the order parameter as a function of temperature for different values of the on-site interaction. The order parameter decreases as a function of increasing temperature until it vanishes at the critical temperature. Close to the phase transition, the order parameter behaves like $m = \alpha|T - T_c|^\beta$ with critical exponent $\beta = 1/2$, which is ~~consistent with the universality class of the spherical model [28, 46]. In~~ different from the exact one belonging to the $O(3)$ (Heisenberg) universality class $\beta \sim 0.369$ [47]. The exponent value for the order parameter $\beta = \frac{1}{2}$ might suggest that TPSC is a mean-field theory. However, the method actually belongs to a different universality class. The critical exponents for TPSC, as well as those for other theories based on λ-Moriya [34, 35] corrections [see, for example, [36, 37, 46], fall within the $O(N)$ universality class in the limit $N \to \infty$ [28]. This class is distinct from the mean-field, which is obtained in the limit of infinite spatial dimensions and is characterized by the exponents $\nu = \frac{1}{2}$, $\gamma = 1$, and $\beta = \frac{1}{2}$. In contrast, in three dimensions, the critical exponents for $O(\infty)$ are $\nu = 1$, $\gamma = 2$, and $\beta = \frac{1}{2}$. In two dimensions, the critical temperature vanishes, as predicted by the Mermin-Wagner theorem, which also holds for the $O(3)$ case. This can be understood by considering the divergence of the sum in Eq.(14) at finite temperature in 2D within the broken symmetry phase, while it remains finite at $T = 0$, where the discrete sum over Matsubara frequencies is replaced by an integral over a continuous variable. Our results for the critical exponent $\beta = \frac{1}{2}$ is therefore consistent with previous calculations showing that TPSC belongs to the $O(\infty)$ universality class.

In Fig.(3-b), we show the value of the vertex renormalisation $\lambda = |\Gamma_{\overline{\uparrow\downarrow}}|/U$ as a function of temperature for different values of U. We observe that λ decreases as a function of increasing interactions, as expected, since the system get more correlated when U increases. On the other hand, λ increases by decreasing the temperature from the critical one, which can be rationalised in the following way: when symmetry breaking is allowed, the system can reduce the number of double occupancies $\langle \hat{n}_\uparrow \hat{n}_\downarrow \rangle = \frac{\lambda}{4}(n^2 - m^2)$, shown in Fig.(3-c), (and therefore minimize the potential energy) by increasing the order parameter, rather than by decreasing λ. Hence, our results show that the degree of correlation of the system is reduced deep in the broken symmetry phase far away from the the critical temperature.

In Fig.(3-d), we show the order parameter and double occupancies as a function of U by fixing the temperature to $T/t = 1/10$. As expected we observe that the order parameter (double occupancies) increases (decrease) as a function of U. It is worth to highlight that the introduction of quantum fluctuations leads to a significant decrease in the staggered ~~magnetization~~ magnetisation compared to its mean-field predicted value [black curve in Fig. (3-d)]. We compared our results to the ones obtained using Monte Carlo in Ref. [13] and we observe an excellent quantitative agreement.

After solving Eqs.(13,14) we can use the values of double occupations and staggered magnetisation as input for Eqs.(19,20) in order to obtain the renormalised vertices in the

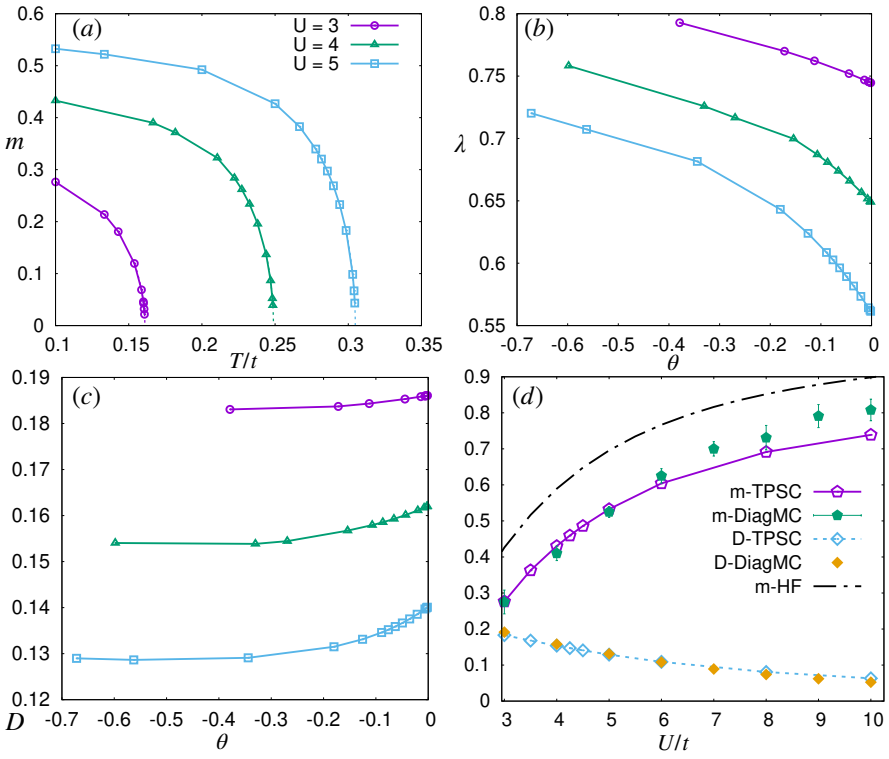

Figure 3: (a) Staggered magnetisation m as a function of T for three different values of $U/t = 3, 4, 5$. Dashed lines are best fits of the function $\alpha|T - T_c|^{1/2}$ close to T_c. (b) λ parameter as a function of the reduced temperature $\theta = \frac{T-T_c}{T_c}$ for the three different values of the on-site interaction. (c) Double occupancies $D = \langle \hat{n}_\uparrow \hat{n}_\downarrow \rangle$ as a function of θ for the three different U values. (d) Magnetisation and double occupancies as a function of U for $T/t = 1/10$. TPSC data (open symbols) are compared to the DiagMC results (filled symbols) adapted from Ref. [13]. The black dashed line is the magnetisation curve obtained using Hartree-Fock.

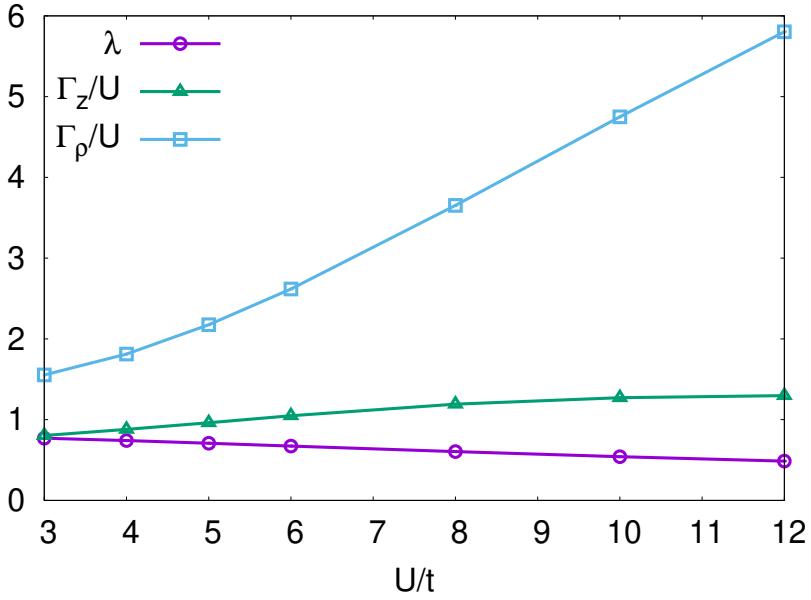

Figure 4: Vertex renormalisations in the density (Γ_ρ/U), spin-longitudinal (Γ_z/U) and spin transverse ($\lambda = |\Gamma_{\overline{\uparrow\downarrow}}|/U$) channels as a function of the bare interaction for $T/t = 1/7.5$.

longitudinal channel. In Fig.(4), we show the renormalisation factors of the vertices, i.e. Γ_ρ/U, Γ_z/U and λ as a function of U for $T/t = 1/7.5$. We observe that Γ_ρ is highly enhanced with respect to the bare vertex which is similar to what has been already observed in the paramagnetic phase of the Hubbard model using TPSC [27]. Differently from the symmetric case, in the AF phase $\Gamma_z \neq |\Gamma_{\overline{\uparrow\downarrow}}|$, and our results show that $\Gamma_z > |\Gamma_{\overline{\uparrow\downarrow}}|$ for all values of U and that the difference between the two vertices increases as a function of the on-site interaction. Interestingly, while $|\Gamma_{\overline{\uparrow\downarrow}}|$ is always lower than the bare vertex (as U_s in the paramagnetic phase [27]), this is not true anymore for Γ_z/U, which is also an increasing function of U and crosses the unity at $U/t \sim 5.4$ for $T/t = 1/7.5$ [see Figure 4].

~~In Appendix ?? we describe how the integral appearing in Eq. (14) has been numerically evaluated~~Integrals in the Brillouin zone were numerically calculated using the trapezoidal rule in three dimensions, employing grids of $N_k \times N_k \times N_k$ points with N_k values up to 32. For the numerical integration of the spin-transverse susceptibility evaluated at zero frequency, i.e., $\int d\mathbf{q} \sum_\sigma \chi_{\sigma\bar{\sigma}}(\mathbf{q}, 0)$, a specific strategy was applied. Since this function diverges at $\mathbf{q} = \mathbf{\Pi}$, that point was excluded from the integration grid. We evaluated the integral for different N_k values (21, 24, 28, 32) and then extrapolated the integral value by fitting the function $I + h/N_k$, where I represents the extrapolated value.

4.1 Dynamical Susceptibilities

We can use the solution of the self-consistent equations to evaluate spectral properties of two-particle propagators. Regarding the spin-transverse channel, we observe that self-energy and vertex corrections are both controlled by the same quantity, i.e. $\Gamma_{\overline{\uparrow\downarrow}}$, which substitutes *de facto* the bare vertex appearing in RPA. Therefore, the spin-transverse dynamical susceptibility defined in Eq.(11), which contains the information about the Goldstone modes, calculated at

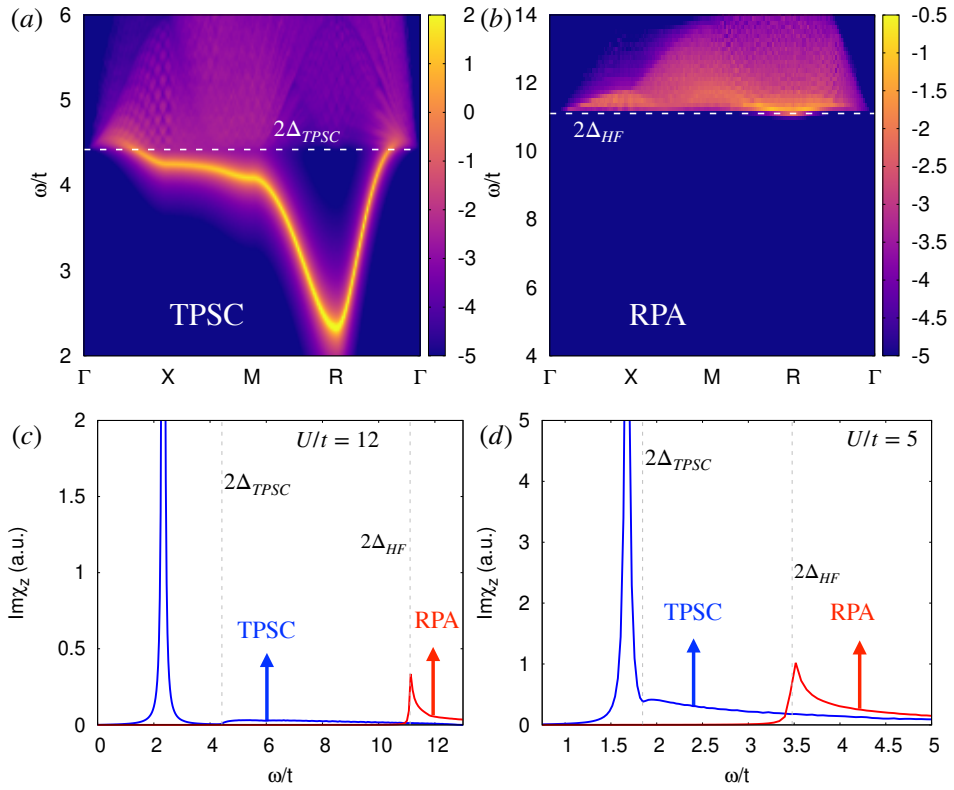

Figure 5: (a) Imaginary part of $\chi_z(\omega + i\eta, \mathbf{q})$ (in log scale) defined in Eq.(17) evaluated along the BZ high-symmetry path and for a wide range of real frequencies, for $U/t = 12$, $T/t = 1/7.5$, and $\eta/t = 0.03$. (b) Imaginary part of $\chi_z(\omega + i\eta, \mathbf{q})$ (in log scale) calculated using RPA for ~~for~~ $U/t = 12$, $T/t = 1/7.5$, and $\eta/t = 0.02$. (c) $\mathrm{Im}\chi_z(\omega + i\eta, \mathbf{q})$ evaluated using TPSC and RPA at fixed momentum $\mathbf{q} = (\pi, \pi, \pi + 0.1)$ at $U/t = 12$, $T/t = 1/7.5$ and for $\eta = 0.03$. (d) Same as (c) but for $U/t = 5$.

a given U corresponds to the RPA one evaluated at a lower value of the interaction, namely $|\Gamma_{\overline{\uparrow\downarrow}}(U)|$.

Conversely, the vertex in the spin-longitudinal susceptibility Γ_z assumes different values than $\Gamma_{\overline{\uparrow\downarrow}}$ because of symmetry breaking, and $\Gamma_z > |\Gamma_{\overline{\uparrow\downarrow}}|$ as shown in Figure 4. This implies that the spin-longitudinal susceptibility evaluated in TPSC does not correspond to any RPA one evaluated at different effective parameters, and consequently the two methods yield qualitatively different results for the spin-longitudinal susceptibility. In particular, since $\Gamma_z > |\Gamma_{\overline{\uparrow\downarrow}}|$ the gap in the χ_z spectrum is reduced with respect to the quasi-particle gap predicted by TPSC, i.e. $2\Delta_{\rm TPSC} = |\Gamma_{\overline{\uparrow\downarrow}}|\, m$, which is controlled by self-energy corrections. In Fig.(5-a) we show a color plot of $\mathrm{Im}\chi_z(q)$ that has been evaluated in the high-symmetry path of the BZ and for a wide range of frequencies at $U/t = 12$ and $T/t = 1/7.5$. We observe that a well visible Higgs mode appears well below the quasi-particle continuum starting at $2\Delta_{\rm TPSC}$, it has a minimum at $R = (\pi, \pi, \pi)$, and presents a substantial dispersion along the M-R and R-Γ directions. This is in stark contrast with the RPA predicted spectrum [shown in Fig.(5-b)], where the Higgs resonance occurs at $\omega/t = 2\Delta_{\rm HF}$ and therefore is overdamped by the particle-hole continuum [48, 49]. Our findings agree qualitatively with recent numerical results based on a time-dependent Gutzwiller approach showing that the Higgs resonance is shifted below the edge of the particle-hole continuum upon increasing the interaction [50]. In Figs.(5-c/d) we show $\mathrm{Im}\chi_z$ evluated using TPSC and RPA as a function of the real frequencies for a fixed momentum close to R and two values of the interactions $U/t = 12, 5$ and at $T/t = 1/7.5$. It is apparent that for both values of the interaction the Higgs resonance predicted by TPSC is well separated from the particle-hole continuum and occurs at lower energies, while RPA does not yield any true isolated pole.

5 Conclusions

We extended the formalism of TPSC to ~~the case of~~ account for spontaneous symmetry breaking and ~~employed~~ applied the new method to the AF phase of the ~~single band Hubbard model in the~~ single-band Hubbard model on a cubic lattice at half-filling. Our comparison with DiagMC ~~shows an~~ reveals excellent quantitative agreement between the two methods for the order parameter and double occupancies.

We show that the differentiation of vertex corrections in the different scattering channels due to symmetry breaking ($\Gamma_z \neq |\Gamma_{\overline{\uparrow\downarrow}}|$) has remarkable effects in the spin-longitudinal channel. In particular, the Higgs resonance occurs at energies lower than the quasi-particle continuum leading to a well visible Higgs mode for a wide range of parameters.

Since our data demonstrate that the level of correlation decreases by decreasing temperature deep in the BSP, one could argue that TPSC is particularly suited to the study of BSP where correlation are not negligible but less pronounced.

~~The formalism that we developed is generic and could be adapted to more complicated multi-band models and open the possibility for efficient treatment of correlation effects in realistic materials in BSP.~~ Additionally, TPSC has already been successfully integrated with ab-initio methods, though only for symmetric phases [31]. This opens up exciting possibilities for extending our method to broken symmetry phases in combination with DFT (Density Functional Theory) for realistic electronic structure calculations.

Also, since TPSC already has been used as a benchmark for cold atomic simulators [51, 52],

its generalisation will provide further guidance to cold-atom experiments exploring broken symmetry phases [53].

Generalising improved version of TPSC such as TPSC+ and TPSC+SFM [33] to the BSP case could lead to the partial inclusion of dynamical effects, which have been shown to be particularly important close to the Neél temperature [54, 55], and is left to future work.

The potential for applying TPSC to understand complex magnetic phases in novel materials is vast. For example, the approach we present here can be applied to models hosting altermagnetism [56, 57, 58, 59], a recently identified category of broken-symmetry phases. Group theory predictions suggest that many such materials might exist in three dimensions [60], providing an ideal scenario where our method can be readily applied. Investigations of these novel magnetic phases in candidate compounds [61, 62, 63, 64, 65, 66] are underway, and we anticipate that new magnetic materials will soon be proposed theoretically and realized experimentally. We also demonstrated that TPSC is an effective tool for studying the amplitude (Higgs) mode, which is often elusive in most mean-field theories. This paves the way for theoretical calculations of amplitude collective modes in altermagnets, providing a reference for future experimental investigations and offering insights into fundamental questions–such as how the topological properties of altermagnets electronic structures [67] are reflected in their collective modes.

Acknowledgment

~~I thank Walter Metzner, Alessandro Toschi, Georg Rohringer, Thomas~~ I thank Walter Metzner, Alessandro Toschi, Georg Rohringer, Thomas Schäfer ~~and Lara Benfatto for valuable discussions. I also thank Renaud Garioud for providing the DiagMC data.~~ and Lara Benfatto for valuable discussions. I also thank Renaud Garioud for providing the DiagMC data.

A ~~Equation of motion and TPSC ansatz~~

~~Let us introduce the generalised multi-flavor indices α, which for example coincide with $\alpha = (a, \sigma)$ containing both sub-lattice (a) and spin (σ) indices in the AF or to spin indices in the FM case. Then, we can rewrite the Hubbard hamiltonian in the following form:~~

$$\underline{H} = \sum_{\langle ij \rangle} \sum_{\alpha\beta} c^{\dagger}_{i\alpha} \mathcal{H}^{\alpha\beta} c_{j\beta} + \frac{1}{2} \sum_{i} \sum_{\alpha\beta} U_{\alpha\beta}\, \hat{n}_{i\alpha} \hat{n}_{i\beta}.$$

~~In the case of FM, we have that $\mathcal{H}^{\alpha\beta} = -t\delta_{\alpha\beta}$ and $U_{\alpha\beta} = \delta_{\alpha\bar{\beta}} U$, whereas for the AF case we have $\mathcal{H}^{\alpha\beta} = -t\delta_{\sigma\sigma'}\delta_{a\bar{b}}$ and $U_{\alpha\beta} = \delta_{\sigma\bar{\sigma}'}\delta_{ab} U$.~~

~~Let us define the electronic Green's function as $G^{\alpha\beta}(x, y) = -T_t \left\langle c_\alpha(x) c^{\dagger}_{\beta}(y) \right\rangle$, where $x = (R_i, \tau_i)$ is a four-vector containing the lattice coordinate R_i and the imaginary time τ_i, and $c_\alpha(x) = e^{H\tau_i} c_{i\alpha} e^{-H\tau_i}$. Then, the equation of motion of the Green's function reads:~~

$$\underline{\partial_{\tau_i} G^{\alpha\beta}(x, y) = \delta_{\alpha\beta}\delta(x - y) - \mathcal{H}^{\alpha\gamma}(x, \overline{y}) G^{\gamma\beta}(\overline{y}, y)}$$

$$\underline{- U_{\alpha\gamma} G^{(2)\beta\alpha}_{\gamma\gamma}(y, x + 0^-, x + 0^+, x),}$$

where a summation over repeated indices is intended, $\mathcal{H}^{\alpha\gamma}(x,y) = \mathcal{H}_{ij}^{\alpha\beta}\delta(\tau_i - \tau_j)$, and $G^{(2)}$ is the two-particle Green's function, that is defined as following:

$$G^{(2)\alpha\beta}_{\gamma\delta}(x_1,x_2,x_3,x_4) = T_\tau \left\langle c^\dagger_\alpha(x_1)c_\beta(x_2)c^\dagger_\gamma(x_3)c_\delta(x_4)\right\rangle .$$

Let us introduce the self-energy Σ through the Dyson's equation:

$$\Sigma^{\alpha\gamma}(x,\overline{y})G^{\gamma,\beta}(\overline{y},y) = G_0^{-1\,\alpha\gamma}(x,\overline{y})G^{\gamma\beta}(\overline{y},y) + \delta_{\alpha\beta}\delta(x-y).$$

If we substitute the last equation into Eq. (??), we obtain the equation of motion in the following form:

$$\Sigma^{\alpha\gamma}(x,\overline{y})G^{\gamma,\beta}(\overline{y},y) = U_{\alpha\gamma}G^{(2)\beta\alpha}_{\gamma\gamma}(y,x+0^-,x+0^+,x).$$

Due to the presence of $G^{(2)}$, Eq.(3) is not closed for the self-energy and single-particle Green's, and in order to obtain an explicit expression for Σ further approximations must be carried on. In mean-field theory the two-particle Green's function is replaced by its disconnected part, that is a valid approximation only at weak coupling. In TPSC [26, 27, 29, 30] , in order to take into account of correlation effects, and at the same time to reduce the complexity of Eq.(3), the following assumption is considered:

$$\Sigma^{\alpha\gamma}(x,\overline{y})G^{\gamma,\beta}(\overline{y},y) \simeq \lambda^{\alpha\gamma}U_{\alpha\gamma}\left[G^{\alpha\beta}(x,y)n_\gamma - s^{\alpha\gamma}G^{\gamma\beta}(x,y)\right],$$

where $n_\alpha = \langle \hat{n}_{i\alpha}\rangle$, $s^{\alpha\beta} = \left\langle c^\dagger_{i\beta}c_{i\alpha}\right\rangle$, and $\lambda^{\alpha\beta}$ is an extra-coefficient that must be determined self-consistently and contains correlation effects. When $\lambda^{\alpha\beta} = 1$ mean-field theory is recovered. The parameter λ can be determined by requiring that the equal-time/position limit of Eq.(3) , i.e. $y = x + 0^{++}$, is preserved exactly when $\beta = \alpha$ [4], by imposing:

$$\lambda^{\alpha\gamma} = \frac{\langle \hat{n}_\alpha\, \hat{n}_\gamma\rangle}{n_\alpha n_\gamma - s^{\alpha\gamma}s^{\alpha\gamma}} .$$

From Eq.(3), we can isolate the self-energy that reads:

$$\Sigma^{\alpha\beta}(x-y) = \delta(x-y)\left(\delta_{\alpha\beta}\lambda^{\alpha\gamma}U_{\alpha\gamma}n_\gamma - \lambda^{\alpha\beta}U_{\alpha\beta}s^{\alpha\beta}\right).$$

In the FM/AF phase of the Hubbard model, the expression for the λ parameter simplifies as following:

$$\lambda = \frac{\langle \hat{n}_\uparrow\hat{n}_\downarrow\rangle}{n_\uparrow n_\downarrow},$$

[4] In the case of the AF phase that we address in this work, spin conservation implies that $\langle c^\dagger_\alpha \hat{n}_\gamma c_\beta\rangle = 0$ at zero field. when $\alpha \neq \beta$, and therefore we shall introduce the λ-correction only for the two-particle Green's functions that do not vanish in the limit of zero external field.

which is identical to the one in the paramagnetic case [27, 29]. Since in the FM/AF phase Eq.(7) does not depend on the spin/sub-lattice indices, we omitted those indices in the expression of λ, and therefore we need to optimize only one parameter even within the broken-symmetry phases under scrutiny. Hence, in the AF phase the expression of the self-energy can be written as:

$$\Sigma^{ab}_{\sigma\sigma'}(x,y) = \delta(x-y)\delta_{ab}U_{\uparrow\downarrow}\left(\delta_{\sigma\sigma'}n_{a\bar{\sigma}} - \delta_{\sigma\bar{\sigma}'}s^{\sigma\bar{\sigma}}_a\right),$$

where $U_{\uparrow\downarrow} = \lambda U$, a, b are sub-lattice indices. We notice that in a AF phase the components of the self-energy off-diagonal in the spin indices should vanish, i.e. $s^{\sigma\bar{\sigma}} = 0$. However, it is useful to keep those terms in the expression of the self-energy for the derivation of the Bethe-Salpeter equation in the spin-transverse channel. Therefore, we will consider the presence of an external field that breaks spin-conservation and eventually compute the functional derivatives of Σ with respect to the off-diagonal component of the propagator in the limit of a vanishing field.

A Irreducible vertices

In this section we shall give some details about the derivation of the expression for the irreducible vertices in the spin-transverse and spin-longitudinal channels.

A.1 Spin-transverse channel

It is worth to note that the expression for the irreducible vertex in the spin-transverse channel presented in the main text is an exact equality. In fact, even if λ is a functional of the Green's function, it does not appear in the expression of the irreducible vertex function because its functional derivative with respect to the off-diagonal propagator vanishes, i.e.

$$\frac{\delta\lambda}{\delta G^{cd}_{\downarrow\uparrow}(x_3,x_4)} = 0. \tag{22}$$

In fact, from Eq.(3) we can derive the following formula for the double occupancies:

$$\langle \hat{n}_{a\sigma}\hat{n}_{a\bar{\sigma}}\rangle = \frac{1}{2U}\Sigma^{aa'}_{\sigma\sigma'}(x,y')G^{a'a}_{\sigma'\sigma}(y',x). \tag{23}$$

Let us now compute the functional derivative of the double occupancies:

$$\frac{\delta\langle \hat{n}_{a\uparrow}\hat{n}_{a\downarrow}\rangle}{\delta G^{\downarrow\uparrow}_{cd}(x_3,x_4)} \propto \delta(x-x_4)\delta_{ad}\Sigma^{\uparrow\downarrow}_{dc}(x_4,x_3) + \frac{\delta\Sigma^{\uparrow\sigma'}_{aa'}(x,\overline{y})}{\delta G^{\downarrow\uparrow}_{cd}(x_3,x_4)}\frac{\delta\Sigma^{\uparrow\sigma'}_{aa'}(x,y')}{\delta G^{\downarrow\uparrow}_{cd}(x_3,x_4)}G^{\sigma'\uparrow}_{a'a}(y',x), \tag{24}$$

where we can now easily see that the LHS does not conserve the spin along the z-axis and therefore vanishes at zero external field.

A.2 Spin-longitudinal channel

On the other hand the expression for the irreducible vertex in the spin-longitudinal channel given in the main text is not an exact equality. Here we shall clarify where the extra approximation comes from. The irreducible vertex function in the longitudinal-spin channel reads:

$$\begin{aligned}\Gamma^{abcd}_{\sigma\sigma'}(x_1,x_2,x_3,x_4) &= \frac{\delta\Sigma^{ab}_{\sigma\sigma}(x_2,x_1)}{\delta G^{cd}_{\sigma'\sigma'}(x_3,x_4)} \\ &= U_{\overline{\uparrow\downarrow}}\delta_{\sigma\bar{\sigma}'}\delta_{ab}\delta_{ac}\delta_{ad}\delta(x_1-x_2)\delta(x_1-x_3)\delta(x_1-x_4) \\ &+ U\, n_{a\bar{\sigma}}\,\delta(x_1-x_2)\delta_{ab}\frac{\delta\lambda}{\delta G^{cd}_{\sigma'\sigma'}(x_3,x_4)}.\end{aligned} \tag{25}$$

Therefore, the irreducible vertex in the spin-longitudinal channel acquires non-local and dynamical corrections, which would complicate the expression of the Bethe-Salpeter equations and further approximations are needed. In practice, one approximates the extra dynamical term to a constant deviation from the value obtained in the spin-transeverse channel, i.e. $\Gamma_{\rho/z} \sim -\Gamma_{\overline{\uparrow\downarrow}} + \delta U_{\rho/z}$.

B ~~Numerical evaluation of integrals~~

~~Integrals in the Brillouin zone have been numerically calculated employing the trapezoidal rule in three dimensions and using grids of $N_k \times N_k \times N_k$ points with N_k up to 32. A particular numerical strategy has been employed for the numerical integration of the spin-transverse susceptibility evaluated at zero frequency, i.e. $\int d\mathbf{q} \sum_\sigma \chi_{\sigma\sigma}(\mathbf{q},0)$. In fact, since this function diverges at $\mathbf{q}=\mathbf{\Pi}$, we have excluded that point from the grid of integration. We evaluated the integral for different values of $N_k = 21, 24, 28, 32$ and then we extrapolated the integral value by fitting the function $I + h/N_k$, where I is the extrapolated value.~~

B Bethe-Salpeter Equations

Let us define the generalized susceptibility as:

$$\chi_{1234} = \frac{\delta G(21)}{\delta h(34)}, \tag{26}$$

where $G(12) = -T_\tau \left\langle c_\alpha(x_1) c^\dagger_\beta(x_2) \right\rangle$ is the propagator, $x = (R,\tau)$, $1 = (\alpha, x_1)$ and $h(12)$ is the perturbing field whose action reads:

$$S_{\text{ext}} = -\int d1 d2\, h(1,2)\, \bar{c}(1) c(2), \tag{27}$$

where in the last equations c and $\bar{c}$ are Grassmann variables, and $\int d1 = \sum_\alpha \sum_R \int_0^\beta d\tau$, with $\beta = 1/k_B T$. Given the form of the external perturbation, the inverse of the non-interacting propagator reads:

$$\mathcal{G}_0^{-1}(12) = [\partial_\tau + \mu - H_0]_{12} + h(12). \tag{28}$$

We now want to obtain a closed equation for χ_{1234} by explicitly performing the functional derivative in Eq.(26). For doing so we first note that:

$$\frac{\delta G(21)}{\delta h(34)} = -\int\int d1'd2'\, G(2,2')\frac{\delta G^{-1}(2'1')}{\delta h(34)}\, G(1',1). \tag{29}$$

We can further develop Eq.(29) by making use of the Dyson equation, that reads:

$$G^{-1}(12) = \mathcal{G}_0^{-1}(12) - \Sigma(12). \tag{30}$$

In fact, by substituting Eq.(30) into Eq.(29) and using Eq.(28), we obtain the following identity:

$$\chi_{1234} = -G(2,3)G(4,1) + \int \prod_{i=1}^{4} di' G(2,2')G(1',1)\Gamma_{1'2'3'4'}\chi_{4'3'34}, \tag{31}$$

where we defined the two-particle irreducible (2PI) vertex function $\Gamma_{1234} = \frac{\delta\Sigma(2,1)}{\delta G(3,4)}$. Let us express the last equation in Fourier space. For this purpose let us expand the propagators and vertices in terms of their Fourier components, i.e.:

$$\begin{aligned} f_{1234} &= \frac{1}{(V\beta)^3}\sum_{kk'q} e^{i[kx_1-(k+q)x_2+(k'+q)x_3-k'x_4]} f^{\alpha\beta}_{\gamma\delta}(k,k',q), \\ G(1,2) &= \frac{1}{V\beta}\sum_{k} e^{-ik(x_1-x_2)} G^{\alpha\beta}_{k}\,. \end{aligned} \tag{32}$$

We first note that:

$$-G(2,3)G(4,1) = \frac{1}{(V\beta)^3}\sum_{kk'q} e^{i[kx_1-(k+q)x_2+(k'+q)x_3-k'x_4]} \times \chi^{\alpha\beta}_{0,\gamma\delta}(k,k',q), \tag{33}$$

where we defined the bubble terms as:

$$\chi^{\alpha\beta}_{0,\gamma\delta}(k,k',q) = -(V\beta)\,\delta_{kk'}\, G^{\delta\alpha}_{k}\, G^{\beta\gamma}_{k+q}. \tag{34}$$

The final equation in Fourier space reads:

$$\chi^{\alpha\beta}_{\gamma\delta}(kk'q) = \chi^{\alpha\beta}_{0,\gamma\delta}(k,k',q) - \frac{1}{(V\beta)^2}\sum_{k_1k_2}\sum_{\alpha'\beta'\gamma'\delta'} \chi^{\alpha\ \beta}_{0,\beta'\alpha'}(k,k_1,q)\,\Gamma^{\alpha'\beta'}_{\gamma'\delta'}(k_1,k_2,q)\,\chi^{\delta'\gamma'}_{\gamma\ \delta}(k_2,k',q)\,. \tag{35}$$

C Improved one-loop self-energy

Let us note that from its definition the generalised susceptiblity is related to the two-particle Green's function in the following way: $\chi^{\alpha\beta}_{\gamma\delta}(x_1,x_2,x_3,x_4) = G^{(2)\alpha\beta}_{\gamma\delta}(x_1,x_2,x_3,x_4) - G^{\beta\alpha}(x_2,x_1)G^{\delta\gamma}(x_4,x_3)$. Hence, we can rewrite the RHS of Eq.(??3) in the following way:

$$\frac{1}{V\beta}\sum_{k\gamma} e^{-ik(x-y)}U_{\alpha\gamma}n_\gamma G^{\alpha\beta}_{k} + \frac{1}{(\beta V)^3}\sum_{kk'q}\sum_{\gamma} U_{\alpha\gamma}e^{-ik(x-y)}\chi^{\alpha\beta}_{\gamma\gamma}(kk'q).$$

α	(a,σ)
β	(b,σ')
γ	(c,σ'')
α'	(a_1,σ_1)
β'	(a_2,σ_2)
γ'	(a_3,σ_3)
δ'	(a_4,σ_4)

Table 1: Relation between indices expressed in the compact and extended notations.

If we substitute Eq.(35) into the second term of last equation we obtain the following expression:

$$-\frac{1}{(V\beta)^2}\sum_{kk'q}\sum_{\gamma}e^{ik(x-y)}U_{\alpha\gamma}G^{\gamma\alpha}_{k}G^{\beta\gamma}_{k+q}$$
$$+\frac{1}{(V\beta)^4}\sum_{kk'qk_1}\sum_{\gamma\alpha'\beta'\gamma'\delta'}U_{\alpha\gamma}G^{\alpha'\alpha}_{k}G^{\beta\beta'}_{k+q}\Gamma^{\alpha'\beta'}_{\gamma'\delta'}(kk_1q)\times\chi^{\gamma'\beta'}_{\gamma\gamma}(k_1k'q), \tag{36}$$

~~The last equation is quite generic and valid for the exact case.~~which is a generic and exact expression of the RHS of Eq.(3). Now we shall specialize to the antiferromagnetic phase of the Hubbard model, and approximate the vertex function to a local quantity that does not depend on the crystalline momenta. In order to do so it is useful to explicitly express the spin-orbital indices in sub-lattice and spin indices as shown in Table 1.

~~Furthermore~~Furthermore, if we ~~assum~~ assume spin-conservation we can express the irreducible vertex function as ~~following:~~ follows:

$$\Gamma^{a_1a_2}_{a_3a_4}|^{\sigma_1\sigma_2}_{\sigma_3\sigma_4}\sim\delta_{a_1a_2}\delta_{a_1a_3}\delta_{a_1a_4}(\Gamma^{a_1}_{\sigma_1\sigma_2}\delta_{\sigma_1\sigma_2}\delta_{\sigma_3\sigma_4}+\Gamma^{a_1}_{\sigma_1\bar{\sigma}_1}\delta_{\sigma_1\bar{\sigma}_2}\delta_{\sigma_3\bar{\sigma}_4}\delta_{\sigma_1\sigma_3}), \tag{37}$$

where we used the following notation $\Gamma^{a}_{\sigma\sigma'}=\Gamma^{aa}_{aa}|^{\sigma\sigma}_{\sigma'\sigma'}$ and $\Gamma^{a}_{\sigma\bar{\sigma}}=\Gamma^{aa}_{aa}|^{\sigma\bar{\sigma}}_{\bar{\sigma}\sigma}$. Substituting Eq.(37) in Eq.(36) we obtain the following expression for the equation of motion in momentum space:

$$\Sigma^{ab}_{\sigma}(k)-Un_{a\bar{\sigma}}=\frac{U}{(V\beta)^3}\sum_{k_1k'q\sigma_1}G^{ab}_{\sigma}(k+q)\Gamma^{b}_{\sigma\sigma_1}(\nu\nu'\omega)\chi^{ba}_{\sigma_1\bar{\sigma}}(k_1k'q). \tag{38}$$

We notice that in this representation the self-energy is expressed in terms of the longitudinal scattering channel only. It is possible to obtain an equivalent expression where the transverse vertex and susceptibility appear by using the following crossing relation:

$$G^{(2)\beta\alpha}_{\gamma\gamma}(y,x+0^-,x+0^+,x)=-G^{(2)\beta\gamma}_{\gamma\alpha}(y,x,x+0^+,x+0^-). \tag{39}$$

Plugging the last equation into the equation of motion in Eq.(~~??~~3) and following similar passages to the ones we did for obtaining Eq.(38), we obtain the following expression for the self-energy:

$$\Sigma^{ab}_{\sigma}(k)-Un_{a\bar{\sigma}}=-\frac{U}{(V\beta)^3}\sum_{k_1k'q}G^{ab}_{\bar{\sigma}}(k+q)\Gamma^{a}_{\bar{\sigma}\bar{\sigma}}(\nu\nu'\omega)\chi^{ab}_{\bar{\sigma}\bar{\sigma}}(k_1k'q). \tag{40}$$

In TPSC the irreducible vertices are local and static, i.e. they do not depend on the Mastubara frequencies and further simplification arise. In particular, if we assume static and local

vertex functions, if we average Eqs.(38,40) we obtain the following expression for the one-loop improved self-energy:

$$\Sigma^{ab}_{\sigma}(k) - Un_{a\bar{\sigma}} = -\frac{U}{2V\beta}\sum_{q} G^{ab}_{\bar{\sigma}}(k+q)\Gamma^{a}_{\bar{\sigma}\bar{\sigma}}\chi^{ab}_{\bar{\sigma}\bar{\sigma}}(q) + \frac{U}{2V\beta}\sum_{q\sigma_1} G^{ab}_{\sigma}(k+q)\Gamma^{a}_{\sigma\sigma_1}\chi^{ab}_{\sigma_1\bar{\sigma}}(q). \tag{41}$$

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