Gross-Neveu and deconfined quantum criticalities in an extended Hubbard model on the square lattice

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We investigate the ground state phase diagram of an extended Hubbard model with $\pi$-flux hopping term at half-filling on a square lattice, with unbiased large-scale auxiliary-field quantum Monte Carlo simulations. As a function of interaction strength, there emerge two interaction-driven quantum critical points, with the first between the Dirac semimetal and an insulating phase of weak valence bond solid (VBS) order, and the second separating the VBS order and an antiferromagnetic insulating phase. These intriguing quantum critical points are respectively bestowed with Gross-Neveu and deconfined quantum criticalities. We also investigate the dynamical properties of the spin excitation and find the spin gap open near the first transition and close at the second. The relevance of our findings in realizing deconfined quantum criticality in fermion systems and the implication to lattice models with further extended interactions such as those in quantum Moiré systems, are discussed.

Introduction.— The Landau-Ginzburg-Wilson (LGW) paradigm of phases and their transitions is one of the cornerstones of modern condensed matter physics [1, 2]. In this celebrated paradigm, the phase transition could be understood in terms of symmetry breaking and the establishment of order parameters, which have greatly shaped our thinking about critical phenomena. According to LGW, the transition between two ordered phases with spontaneously broken symmetries should either be first-order or through an intermediate phase. However, new exotic transitions between novel quantum states that are beyond the LGW paradigm have been accumulating in recent years. For example, Senthil et al. proposed that a continuous quantum phase transition between the antiferromagnetic (AF) order and the valence bond solid (VBS) order could exist [3, 4], dubbed as deconfined quantum critical point (DQCP). Strong evidence for such DQCP in spin-1/2 model on square lattice has been first shown in the J-Q model by Sandvik [5, 6]. Subsequently, other numerical examples and new theoretical understanding of DQCP have been developed in quantum spin models [7–19] and have been gradually extended to interacting fermionic systems [20–25]. It is obvious that model design and large-scale quantum Monte Carlo (QMC) simulations played a key role in pushing the frontier of our knowledge on such surprising phenomena.

Except the DQCP discussed above, interaction effects on massless Dirac fermions have also attracted great attentions. Since the linear dispersion is stable against weak interactions, there must be one or more quantum phase transitions separating the Dirac semimetal (SM) phase and various possible Mott insulator states. Depending on the type of interactions, many studies have discovered the insulators being the ferromagnetic and AF states [26–32], VBS state [33–38], nematic phase [39], superconductivity, quantum (spin) Hall states [23–25, 40, 41] and many others. Among these examples, of particular interests is the situation where from Dirac SM to the strong-coupling limit, there exist multiple insulating phases as a function of the interaction strength, and leaves room for interesting intermediate phases such as topological ordered phases and multiple exotic quantum phase transitions, such as Gross-Neveu and DQCP. Previous works have shown, with spin-1/2 electron and SU(2) symmetry, extended interaction on honeycomb lattice, offer a robust VBS with Kekulé pattern exactly as such an intermediate phase between Dirac SM and strong-coupling AF order [37]. However, although the transition between Dirac SM with Kekulé VBS is found to be a Gross-Neveu QCP, the transition between VBS and AF phases is first order. These results motivate us to investigate the interaction effect in $\pi$-flux extended Hubbard model on square lattice, as we show below, in this case, except a Gross-Neveu QCP between Dirac SM and VBS, there indeed further emerges a continuous transition between the VBS and AF phase within the largest system size we can access, consistent with the expected behavior of the DQCP. Our results of the Dirac fermion with an extended interaction, could also shed light on the great on-going efforts in understanding the interaction effects on quantum Moiré material models [38, 42–48] such as twisted bilayer graphene and transition metal dichalcogenides, where the interplay between flat band Dirac cones and the extended Coulomb interactions can be engineered by twisting angles, gating and tailored design of the dielectric environment, and give rise to a plethora of exotic phenomena waiting for explanations.

Model and method.— We study a SU(2) extended Hubbard model with $\pi$-flux hopping term at half-filling on the square lattice, as shown in Fig. 1, with the following Hamiltonian

\[
H = - \sum_{\langle ij \rangle, \sigma} t_{ij} \left( c_{i\sigma}^\dagger c_{j\sigma} + \text{H.c.} \right) + U \sum_{\square} (n_{\square} - 1)^2 ,
\]
where \((ij)\) denotes the nearest neighbor, \(c^\dagger_{ij}\) and \(c_{ij}\) are creation and annihilation operators for fermions on site \(i\) with spin \(\sigma = \uparrow, \downarrow\), \(n_0\) is the extended particle number operator of \(\square\)-plaque defined as \(n_0 \equiv \sum_{\text{sites}} n_i\) with \(n_i = \frac{1}{4} \sum_{\sigma} \epsilon^\dagger_{i\sigma} \epsilon_{i\sigma}\) and at half-filling \(\langle n_\text{eff} \rangle = 1\), \(U\) tunes the interaction strength, which favors AF order in the strong coupling limit [36].

As shown in Fig. 1 (a), we set hopping amplitudes \(t_{i,i+\hat{x}} = t\) and \(t_{i,i+\hat{y}} = (-1)^i t\), where the position of site \(i\) is given as \(r_i = i_x \hat{e}_x + i_y \hat{e}_y\) and \(t = 1\) is the energy unit. Such arrangement bestows a \(\pi\)-flux penetrating each \(\square\)-plaque. As a consequence, two Dirac cones are located at \(K_0 = \left(\frac{\pi}{2}, \pm \frac{\pi}{2}\right)\) in the first Brillouin zone (BZ). We notice although the above hopping amplitudes appear to break translation and rotation symmetries and the BZ of the square lattice is folded in half, as shown in the blue area in Fig. 1 (b). The way of the folding and locations of the Dirac cones all depend on the gauge choice of hopping amplitudes, while the distance between two Dirac cones does not. The model therefore still has full crystalline symmetries of the square lattice (the p4mm wallpaper group), where each crystalline symmetry operation is supplemented by a U(1) gauge transformation. For example, the translation symmetry becomes \(T^\epsilon_x : c_i \rightarrow (-)^i c_{i+\hat{e}_x}\). Consequently, we will still discuss our results in the original square lattice BZ. For the extended Hubbard interaction term, the onsite, first and second nearest neighbor repulsions are all included in one plaquette. As found in the previous studies [37, 38, 42, 43], such extended interactions will require a relative larger \(U\) to gap out the Dirac cones.

One can easily show the Hamiltonian in Eq. (1) is sign problem free for auxiliary field QMC [49] and we implement a projector version of QMC method [50] to solve the model. Details of the algorithm can be found in Supplemental Material (SM) [51], and we only mention the projection length \(\beta t = L\) for equal-time measurements: \(\beta t = L + 10\) for imaginary-time measurements and discrete time slice \(\Delta \tau = 0.1\). We simulate the systems with linear size \(L = 12, 16, 20, 24\) and 28. We have also tested that this setup is enough to achieve controlled errorbars for our model [51].

**QMC results** — The phase diagram obtained from QMC simulations is shown in Fig. 2 (a). We find two quantum critical points when gradually increasing the interaction strength \(U\). They are the phase transition from Dirac SM to VBS phase and that from VBS to AF phase. The first QCP is at \(U_{c1}/t = 23.5(5)\) and of Gross-Neveu type with the VBS acquiring a \(Z_4\) discrete lattice symmetry breaking, and the critical point is expected to have emergent U(1) symmetry as the \(Z_4\) anisotropy is irrelevant [6, 37]. The second QCP is at \(U_{c2}/t = 29.25(25)\), separating two spontaneous symmetry breaking phases, e.g. \(Z_4\) for VBS and SU(2) for AF phases, and shall be explained according to the deconfined quantum criticality [4, 5]. We note the existence of a DQCP in the Hubbard model on the square lattice has not been reported before. In particular, the Hubbard model with only onsite interactions hosts only Dirac SM and AF phases [28, 29], and on the honeycomb lattice with extended interaction, the corresponding transition is first order [37].

FIG. 2. (a) Phase diagram of the extended Hubbard model as function of interaction strength \(U\), obtained from QMC simulations. The SM-VBS transition at \(U_{c1}\) is continuous and of Gross-Neveu universality. The VBS-AF transition at \(U_{c2}\) is also continuous and should be explained according to the deconfined quantum criticality. Correlation ratios of the (b) VBS order and (c) AF order as a function of interaction strength \(U\) are shown.

To quantitatively study the two phase transitions, we define two structure factors

\[
C_{\text{AF}}(k, L) = \frac{1}{L^2} \sum_{i,j} e^{ik \cdot (r_i-r_j)} \langle S_i S_j \rangle \tag{2}
\]

and

\[
C_{\text{VBS}}(k, L) = \frac{1}{L^2} \sum_{i,j} e^{ik \cdot (r_i-r_j)} \langle B_i B_j \rangle \tag{3}
\]

for AF order and VBS order, respectively. In above equations, \(S_i = \frac{1}{2} \sum_{\sigma} \sigma \epsilon^\dagger_{i\sigma} \epsilon_{i\sigma}\), \(\sigma\) denotes the vector of the three Pauli matrices.

\(B_i = \sum_{\sigma} \left( t_{i,i+\hat{e}_x} \epsilon^\dagger_{i+\hat{e}_x,\sigma} + \text{H.c.} \right)\) are gauge invariant bond
operators. For AF order, $C_{AF}(k, L)$ is peaked at momentum $\mathbf{M} = (\pi, \pi)$; for VBS order, $C_{VBS}(k, L)$ is peaked at momentum $\mathbf{X} = (\pi, 0)$. We then use the renormalization-group invariant correlation ratios ($c = AF, VBS$) to perform the data analysis,

$$R_c(U, L) = 1 - \frac{C_c (k = \mathbf{Q} + d\mathbf{q}, L)}{C_c (k = \mathbf{Q}, L)},$$

where $\mathbf{Q}$ is the ordering wave vector, $|d\mathbf{q}| \sim \frac{1}{L}$ denotes the smallest momentum on finite size lattice. By definition, $R_c(U, L) \rightarrow 1 (0)$ for $L \rightarrow \infty$ in the corresponding ordered (disordered) phase. At the QCP, $R_c$ is scale invariant for sufficiently large $L$ and exhibit the scaling behavior for the corresponding universalities [30, 52, 53]. As shown in Fig. 2 (b), when varying $U/t$ from 18 to 33, $R_{VBS}(U, L)$ first increases then decreases. Importantly, $R_{VBS}(U, L) = 20, 24, 28$ for different $L$ cross at two separate regions. These results mean that our model undergoes two phase transitions, and the VBS order is the intermediate phase. Admittedly, the VBS order is very weak and we believe it is attributed to the enhanced quantum fluctuations from the interplay of Dirac fermions and extended interactions within a plaquette. We also find at the transition, the VBS order parameter histogram is consistent with emergent $U(1)$ symmetry [37]. The correlation ratio of AF order is shown in Fig. 2 (c). Interestingly, a clear crossing takes place around $U/t = 29.4$, which further indicates the phase transition between the VBS and the AF order is continuous.

To understand the two intriguing QCPs, we first focus on the more complicated VBS-AF transition, to demonstrate it has the flavor of DQCP. To this end, we collapse the correlation ratio of AF order with finite-size scaling relation $R_{AF}(U, L) = f_1 \left( L^{1/\nu} (U - U_c - 1) \right)$ [30, 54], as shown in Fig. 3 (a), and obtain the position of corresponding QCP at $U_{c2}/t = 29.25(25)$ and the correlation length exponent $\nu = 1.13(5)$. Then at $U_{c2}$ the AF and VBS structure factors, in Eqs. (2) and (3), shall obey the following finite-size scaling ansatz [38, 55]

$$C_c(U, L) = L^{-\zeta - \eta f_2 \left( L^{1/\nu} (U - U_c) / U_c \right)}, \tag{5}$$

where $\eta$ is the anomalous dimension exponent, $z = 1$ is the dynamic exponent. As shown in Fig. 3 (c) and (d), we extract $\eta$ from the slope of log-log plot of $C_c(U, L)$ vs $L$ curves at $U_{c2}/t = 29.25$ and find $\eta_{AF} = 0.58(3)$ for AF order and $\eta_{VBS} = 0.6(1)$ for VBS order, respectively, give rise to good quality linear fits.

According to the scenario of DQCP [3, 4], the closeness of the exponents from the two ordered phase approaching the critical point, i.e. $\eta_{AF} \approx \eta_{VBS}$ in our setting, is considered as an important signature for the associated emergent symmetry [7, 8], and numerical evidences of such closenesses have been seen in the J-Q model [5, 11], 3D loop model [8], as well as those in the recently discovered the fermionic models [23–25]. In the literature [7, 8, 11, 23–25], the precise value of the exponents appears to depend on the detailed implementation of each model, and such discrepancies also complicate the understanding of DQCP. There also exists the conformal field theory (CFT) bound that the emergent symmetry need to satisfy [56], and the more recent entanglement measurements further point out the DQCP, at least in the J-Q model, might not be an unitary CFT at the first place [17, 19] and other possible scenarios such as multicritical point [15], complex CFT [57, 58] and weakly first order transition [16], have been constantly and actively proposed and explored. Despite of these efforts and the enigmatic situation of the DQCP, our observation, that within the size systems simulated, $\eta_{AF} \approx \eta_{VBS}$, is consistent with the deconfinement at $U_{c2}$. It is further worthwhile to point out that a relative large $\eta \approx 0.6$ is also the hallmark of many QCPs associated with the fractionalization of elementary excitations such as the condensation transition of spinon and visons in $Z_N$ topological orders [59–61]. We also collapse the AF structure factor according to the Eq. (5) with $U_{c2}/t = 29.25$, $\nu = 1.13$ and $\eta = 0.58$, as shown in Fig. 3 (b), all data points fall on a smooth curve. Therefore, our numerical data in Fig. 3 certainly reveal an internally consistent description along the line of DQCP for the VBS-AF transition.

Next we move on to the SM-VBS transition. It is known that the massless Dirac fermion is robust at weak interaction, and the single-particle gap will only open at a finite interaction strength [22, 42]. In our model, we indeed find as a function of $U$, the Dirac SM transit to an insulating VBS order continuously through a Gross-Neveu QCP [30, 33–38, 62].
This is also consistent with the similar situation of the extended Hubbard model on honeycomb lattice [37, 38]. To find the accurate position of \( U_{c1} \), we can extract the single-particle gap \( \Delta_{sp}(K_0, L) \) from a fit to the asymptotic long imaginary time behavior of the single-particle Green's function

\[
G(k, \tau, L) = \left(1/L^4\right) \sum_{i} e^{ik \cdot (r_i - r_j)} \langle c_{i,\sigma}(\tau) c_{j,\sigma}(0) \rangle \approx e^{-\Delta_{sp}(k, L)\tau}.
\]

The obtained \( \Delta_{sp}(K_0, L) \) are shown in Fig. 4 (a). It is clear that \( \Delta_{sp}(K_0, L \to \infty) \to 0 \) at \( U/t < 23 \) and \( \Delta_{sp}(K_0, L \to \infty) > 0 \) at \( U/t > 24 \), which indicate the critical point \( U_{c1} \) lies in the range \( U/t \in (23, 24) \). The position of QCP of SM-VBS transition obtained from the extrapolation of \( \Delta_{sp}(K_0) \) is overall consistent yet a bit deviated with the cross point of \( R_{VBS} \) shown in Fig. 2 (b). This again signifies the weakness of the VBS order and the strong fluctuations at this QCP which give rise to strong finite-size effect. To locate the \( U_{c1} \) with more accuracy, as shown in Fig. 4 (b), we try to collapse the VBS structure factor according to the Eq. (5) near the range \( U/t \in (23, 24) \). Although the finite-size effect is strong, the data collapse results behave well, which gives rise to an estimation of position of QCP, \( U_{c1}/t = 23.5(5) \), and critical exponents, \( \nu' = 1.0(1) \) and \( \eta' = 0.89(3) \). These exponents are consistent with previous numerical simulations of similar SM-VBS transitions on the honeycomb lattice [37]. Since the VBS orders on the two lattices are three-fold and four-fold degenerate, respectively, the consistent exponents suggest that, at both QCPs, the three-fold and four-fold lattice rotation symmetries are enhanced to an emergent U(1) symmetry. In fact, the emergent U(1) symmetry at the QCP on the honeycomb lattice [37] implies the three-fold anisotropy of the U(1) order parameter is (dangerously) irrelevant at the QCP, which further implies that the four-fold anisotropy should be even more irrelevant.

At \( U < U_{c1} \), the ground state is in Dirac SM state, and thus there is no spin excitation gap in the thermodynamic limit (TDL). At \( U > U_{c2} \), the ground state is in AF state, the spin excitation gap in the TDL should also vanish because of the existence of Goldstone mode. However, the spin excitation gap will open in the VBS state due to the formation of a spin singlet [38]. To verify the theoretical predic-

![Fig. 4](attachment:image4.png)

**FIG. 4.** (a) The \( 1/L \) extrapolation of the single-particle gap \( \Delta_{sp}(K_0, L) \), the gap opens at an interaction strength locating in a range from \( U/t = 23 \) to \( U/t = 24 \). (b) The data collapse of the structure factor of VBS order, which gives an estimation of \( U_{c1}/t = 23.5(5) \), \( \nu' = 1.0(1) \) and \( \eta' = 0.89(3) \).

With the help of large scale sign-free projector QMC simulation, we investigate the phase diagram of spin-1/2 \( \pi \)-flux extended Hubbard model on square lattice at half-filling. Based on all the numerical results obtained, we conclude that this simple looking model acquire an interesting phase diagram with an intermediate phase with weak VBS order separating the well known Dirac SM and AF phases. More importantly, we find the Gross-Neveu transition from Dirac SM to VBS is continuous, and the transition from VBS to AF is also continuous, consistent with deconfined quantum
critical criticality.

Our results, along with the previous works of extended interaction model on honeycomb lattice [36–38], point out the directions that to realize interesting phase diagrams with (multiple) intermediate phases between the free Dirac SM and strong coupling limit insulators, the extended interactions beyond on-site Hubbard term, will bring more surprises. In fact, the weak VBS order discovered here and the DQCP associated with it towards the AF order, imply further perturbations could give rise to even more exotic interaction-driven phases and transitions. In this context, our results could also have the relevance to the great on-going efforts in understanding the interaction effect on quantum Moiré material models [38, 42–46] such as twisted bilayer graphene and transition metal dichalcogenides, where the interplay between flat band Dirac cones and the extended dielectric environment, and give rise to a plethora of exotic phenomena and interesting phase and phase transitions. It is natural to anticipate, with the technique and analysis presented in this work, once further degrees of freedom such as pseudospin and valley in the Moiré systems are added, interesting phases and transitions and their physics mechanism will be revealed from the lattice model simulations.

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[51] The implementation of the projector quantum Monte Carlo simulation, the analysis of the Trotter error and more data of structure factors as a function of interaction strength, are presented in this Supplemental Material.

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Supplemental Material for "Gross-Neveu and deconfined quantum criticalities in an extended Hubbard model on the square lattice"

I. PROJECTION QMC METHOD AND SIMULATION DETAILS

Since we are interested in the ground state properties, the projection quantum Monte Carlo (PQMC) method is a suitable choice. In PQMC, we calculate the ground state wave function |Ψ_0\rangle through the projection of a trial wave function |Ψ_T\rangle i.e. |Ψ_0\rangle = \lim_{β→∞} e^{-\frac{β}{2}H}|Ψ_T\rangle [S50]. Physical observables can be measured according to

$$\langle \hat{O} \rangle = \frac{\langle Ψ_0 | \hat{O} | Ψ_0 \rangle}{\langle Ψ_0 | Ψ_0 \rangle} = \lim_{β→∞} \frac{\langle Ψ_T | e^{-\frac{β}{2}H} \hat{O} e^{-\frac{β}{2}H} | Ψ_T \rangle}{\langle Ψ_T | e^{-βH} | Ψ_T \rangle}. \quad (S1)$$

where \( \hat{O} \) is the concerned physical observable, \( H \) is the Hamiltonian operator, \( β \) is the projection length. Since \( H \) consists of the non-interacting and interacting parts, \( K \) and \( V \), respectively, that are usually not commute, we should performed Trotter decomposition to discretize \( β \) into \( L_τ \) slices (\( β = L_τ Δτ \)), then

$$\langle Ψ_T | e^{-βH} | Ψ_T \rangle = \langle Ψ_T | \left( e^{-ΔτHU} e^{-ΔτH0} \right)^{L_τ} | Ψ_T \rangle + O(Δτ^2) \quad (S2)$$

where the non-interacting and interacting parts of the Hamiltonian is separated. It is worth to note that above processing will give rise to a small systematic error \( O(Δτ^2) \), thus we should set \( Δτ \) as a small number. In Hubbard model, the interaction parts contain the quartic fermionic operator that can not be measured directly with PQMC method. To calculate the interacting part, one need to employ a SU(2) symmetric Hubbard-Stratonovich (HS) decomposition, then the auxiliary fields will couple to the charge density. For example, in our extended Hubbard model on square lattice at half-filling, the HS decomposition is required to employ a SU(2) symmetric Hubbard-Stratonovich (HS) decomposition, then the auxiliary fields will couple to the charge density. For example, in our extended Hubbard model on square lattice at half-filling, the HS decomposition is

$$\left( \begin{array}{c} e^{-ΔτU(n(0)−1)}^2 \ = \ \frac{1}{4} \sum_{(s_0)} \gamma(s_0) e^{αη(s_0)(n(0)−1)} \end{array} \right) \quad (S3)$$

with \( α = \sqrt{−ΔτU} \), \( γ(±1) = 1 + \sqrt{6}/3 \), \( γ(±2) = 1 − \sqrt{6}/3 \), \( η(±1) = ±\sqrt{2(3−√6)} \), \( η(±2) = ±\sqrt{2(3+√6)} \) and the sum symbol is taken over the auxiliary fields \( s_0 \) on each square plaquette.

In the end, we derive out the following formula with constant factors omitted

$$\langle Ψ_T | e^{-βH} | Ψ_T \rangle = \sum_{\{s_0, τ\}} \left[ \prod_τ \prod_△ \gamma(s_{△, τ}) e^{-4αη(s_{△, τ})} \right] \det \left[ P^α B(β, 0) P \right] \quad (S4)$$

here \( P \) is the coefficient matrix of trial wave function \( |Ψ_T\rangle \); \( B \) matrix is defined as \( B(τ + 1, τ) = e^{V_τ} e^{−ΔτK} \) and has a property \( B(τ_1, τ_2) = B(τ_2, τ_1) B(τ_1, τ_2) \). In the practice, we choose the ground state wavefunction of the half-filled non-interacting parts of Hamiltonian as the trial wave function. The Monte Carlo sampling of auxiliary fields are further performed based on the weight defined in the sum of Eq. (S4). The measurements are performed near \( τ = β/2 \). Single particle observables are measured by Green’s function directly and many body correlation functions are measured from the products of single-particle Green’s function based on their corresponding form after Wick-decomposition. The equal time Green’s function are calculated as

$$G(τ, τ) = 1 − R(τ) (L(τ) R(τ))^{-1} L(τ) \quad (S5)$$

with \( R(τ) = B(τ, 0) P, L(τ) = P^† B(β, τ) \). The imaginary-time displaced Green’s function \( G(τ, 0) \equiv \langle c(\frac{τ}{2}) c^†(−\frac{τ}{2}) \rangle \) are calculated as

$$\langle c(\frac{τ}{2}) c^†(−\frac{τ}{2}) \rangle = \frac{\langle Ψ_T | e^{−(\frac{τ}{2}+\frac{τ}{2})H} c e^{−τH} c^† e^{−(\frac{τ}{2}−\frac{τ}{2})H} | Ψ_T \rangle}{\langle Ψ_T | e^{-βH} | Ψ_T \rangle} \quad (S6)$$

where \( c = \{c_1, c_2, ..., c_N \}, N = L × L \) is the system size. More technique details of PQMC method, please refer to Refs [S50].

As mentioned in our paper, we set projection time \( βt = L \) for equal-time measurement, \( βt = L + 10 \) for imaginary-time measurement and discrete time slice \( Δτ = 0.1 \). As shown in Fig S1, we have confirmed that this setup is enough to achieve convergent and error controllable numerical results for our model.
II. THE SMOOTH CURVES OF STRUCTURE FACTORS

As shown in previous QMC studies [S37, S38], physical observable, like kinetic energy, structure factors, correlation ratios and so on, will present different behavior depending on the different phase transition types, i.e. continuous or first-order. For the first-order phase transition, the curve of physical observable verse interaction strength will show a kink or jump, while for continuous one, smooth. We have shown the smooth behavior of kinetic energy as varying interaction strength in main text. Here, as shown in the Fig S2, we notice that the curves of structure factors are also smooth, which support the SM-VBS and VBS-AF transition all are continuous.

FIG. S2. The structure factors of (a) AF order and (b) VBS order as function of interaction strength $U$. All lines in the figure are guides to the eye, we can notice they show no jump, kink or discontinuous.