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Finite correlation length scaling with infinite projected entangled pair states at finite temperature

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We study second-order finite-temperature phase transitions of the two-dimensional quantum Ising and interacting honeycomb fermion models using infinite projected entangled pair states (iPEPS). We obtain an iPEPS thermal state representation by variational tensor network renormalization. We find that at the critical temperature $T_c$ the iPEPS correlation length is finite for the computationally accessible values of the iPEPS bond dimension $D$. Motivated by this observation, we investigate the application of finite correlation length scaling (FCLS), which has been previously used for iPEPS simulations of quantum critical points at $T = 0$, to obtain precise values of $T_c$ and the universal critical exponents. We find that in the vicinity of $T_c$ the behavior of observables follows well the one predicted by FCLS. Using FCLS we obtain $T_c$ and the critical exponents in agreement with quantum Monte Carlo results except for couplings close to the quantum critical points where larger bond dimensions are required.

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I. INTRODUCTION

Tensor networks [1–5] are representations of weakly entangled states obeying an area law of entanglement [6–8]. They are a basis for variational numerical methods for strongly correlated quantum many-body systems, enabling simulations of fermionic, bosonic, and spin models with the same leading computational complexity [9–14]. The powerful density matrix renormalization group (DMRG) [15,16] approximates a state of a system by a one-dimensional (1D) tensor network called matrix product state (MPS) [17–19]. Two-dimensional (2D) projected entangled pair states (PEPS) [20], called also tensor product states [21,22], were initially applied as a variational Ansatz for 2D ground states [12,14,23–25] bringing new insights into paradigmatic models of strongly correlated systems (see, e.g., [26–32]). Recent years brought new applications of PEPS [33–37] and further progress in the fields of numerical optimization [38–41] and contraction [42,43] of PEPS.

Thermal states of 2D local Hamiltonians obey an area law for mutual information, which is reproduced by projected entangled pair operators (iPEPO) representing thermal states and infinite projected entangled pair states (iPEPS) representing purifications of thermal states [44], giving motivation to use iPEPS for thermal state simulations [45,46]. Recently, new methods for simulation of thermal states, based on iPEPS and iPEPO, were proposed [36,47–50]. Recent years brought also developments in the field of the closely related direct contraction methods for three-dimensional (3D) tensor networks representing partition functions of 2D quantum models [51–55] and MPS/MPO based simulations of thermal states of finite-width cylinders [56,57]. Some of those methods were already applied to challenging problems [58–62].

Among demanding problems in the field of 2D strongly correlated systems are finite-temperature critical phenomena and in particular finite-temperature second-order phase transitions. Some of the methods mentioned above were already applied to investigate 2D critical phenomena based on the assumption that large enough $D$ can be obtained to provide results which are converged in $D$ [36,58,59,61], however, for more challenging cases reaching convergence in $D$ will in general be difficult.

Here, we demonstrate that even in the case when convergence in $D$ cannot be obtained, it is possible to take finite-$D$ effects systematically into account using a finite correlation length scaling (FCLS) [63–66]. Furthermore, we show that FCLS can be used to obtain critical data for a finite-temperature phase transition, i.e., the critical temperature $T_c$ and the universal critical exponents.

FCLS, originally called finite entanglement scaling (FES), was first proposed to investigate 1D quantum critical points by infinite MPS (iMPS) [63,64,67]. These critical points violate the area law of entanglement [68] and as such cannot be represented by finite-$D$ iMPS, which have a finite correlation length $\xi_D$. It was shown that in the case of the optimal iMPS finite-$D$ ground-state approximation, the finite $D$ modifies observables of the critical state as if the system was finite with the size proportional to $\xi_D$ [63,64]. It was also shown that the scaling of the observables with increasing $\xi_D$ can be used to determine the critical exponents and the precise location of the critical point similarly as in standard finite-size scaling for quantum Monte Carlo (QMC) simulations [63,64,67].

A similar idea was applied earlier in corner transfer matrix renormalization group (CTMRG) simulations of 2D critical thermal states of classical models. CTMRG approximately contracts a 2D tensor network representing a partition function of a 2D classical system. The approximation introduces an effective correlation length controlled by a refinement parameter $\chi$ of the method [69]. A scaling Ansatz assuming that this correlation length is proportional to the effective...
finite-
how to determine the critical temperature
In Sec.VII we introduce FCLS and in Sec.VIII we describe
determine the correlation length of an iPEPS using CTMRG.
contract a 2D tensor network. In Sec.VI we explain how to
describe the CTMRG method which allows us to efficiently
apply to determine the critical data.
benchmark results demonstrating that indeed FCLS can be
applied to determine the critical data.
This paper is organized as follows: In Sec. II we introduce
a thermal state’s representation by a purification and in Sec. III
we describe how to represent such a purification using an
iPEPS. In Sec. IV we introduce VTNR which we use to obtain
the purification’s iPEPS representation. In Sec. V we
describe the CTMRG method which allows us to efficiently
contract a 2D tensor network. In Sec. VI we explain how to
determine the correlation length of an iPEPS using CTMRG.
In Sec. VII we introduce FCLS and in Sec. VIII we describe
how to determine the critical temperature \( T_c \) for a second-
order phase transition using FCLS. In Secs. IX and X we
present the benchmark results for the application of FCLS to
simulations of thermal second-order phase transitions in the
quantum Ising and interacting honeycomb fermions models,
respectively. Finally, we provide our conclusions in Sec. XI.

II. PURIFICATIONS OF THERMAL STATES

A thermal state of a Hamiltonian \( H \) for temperature \( T \)
given by its thermal density matrix
\[
\rho(T) = \frac{1}{Z(T)} e^{-H/T}, \quad Z(T) = \text{Tr} e^{-H/T}.
\]
Here, we consider lattice models for which the Hilbert space \( \mathcal{H} \)
is a tensor product of Hilbert spaces of individual lattice sites \( \mathcal{H}_i \) spanned by states \( \{|s_i\}, \ s_i = 1 \ldots d \}.\n
We represent \( \rho(T) \) by its purification \( |\Psi(T)\rangle \) which is a
pure state in an enlarged Hilbert space \( \tilde{\mathcal{H}} \) created by intro-
ducing ancillary degrees of freedom. \( \tilde{\mathcal{H}} \) is a tensor product
of enlarged Hilbert spaces of individual sites \( \tilde{\mathcal{H}}_i \), which are
spanned by states \( \{|s_i, a_i\}, \ s_i = 1 \ldots d, \ a_i = 1 \ldots d \} \) with an
index \( a_i \) numbering the ancillary degrees of freedom. To obtain \( \rho(T) \) from \( |\Psi(T)\rangle \) one needs to trace out the ancillary
degrees of freedom
\[
\rho(T) = \text{Tr}_{a_i} |\Psi(T)\rangle \langle \Psi(T)|.
\]
For \( T = \infty \) we have
\[
|\Psi(T = \infty)\rangle \propto \bigotimes_i \left( \sum_{s_i=1 \ldots d} |s_i\rangle \right),
\]
and for finite \( T \), \( |\Psi(T)\rangle \) is obtained by an action of \( e^{-H/(2T)} \)
on the physical degrees of freedom of \( |\Psi(T = \infty)\rangle \):
\[
|\Psi(T)\rangle \propto e^{-H/(2T)} |\Psi(T = \infty)\rangle.
\]

III. IPEPS REPRESENTATION OF THERMAL STATES

A projected entangled pair state (PEPS) [20], also
called a tensor product state [21,22], is a 2D tensor network
representing a state obeying the area law of entanglement. In
the simplest case, a PEPS represents a pure state and is built
from a network of rank-5 tensors on a square lattice, with one
tensor per lattice site. Each tensor has a physical index
representing the local Hilbert space of a site. The other four
indices of a tensor, called the virtual indices, are contracted
with the virtual indices of the neighboring tensors. Their
dimension is called the bond dimension \( D \) which controls the
accuracy of the Ansatz. With growing \( D \) states with stronger
tanglement can be represented by the PEPS. An infinite
projected entangled pair state (iPEPS) is a PEPS representing
a state on an infinite lattice. To obtain an iPEPS we introduce
a unit cell of tensors which is periodically repeated on the
lattice, i.e., the iPEPS is translational invariant by shifts of
the unit-cell size. With each site in the unit cell we associate
a different PEPS tensor. In this work we use a unit cell
with two tensors \( A \) and \( B \) arranged in a checkerboard pattern
(all states studied in this work are compatible with this unit
cell).

As \( |\Psi(T)\rangle \) is a pure state in the extended Hilbert space,
it can be represented by an iPEPS with the tensors having an
additional index for the ancillary degrees of freedom [see Figs. 1(a) and 1(b)]. Then, \( \rho(T) \) (2) can be obtained by a
contraction of the iPEPS representation of \( |\Psi(T)\rangle \) and its
Hermitian conjugate [see Fig. 1(c)].

The state at infinite temperature, \( |\Psi(T = \infty)\rangle \), can be repres-
ented exactly by an iPEPS with the bond dimension \( D = 1 \), i.e.,
a product state (3). For finite \( T \), \( |\Psi(T)\rangle \) is obtained by an
action of an operator \( e^{-H/(2T)} \) on the physical degrees of
freedom of \( |\Psi(T = \infty)\rangle \) [Eq. (4)]. As \( e^{-H/(2T)} \) in general does
not have a numerically tractable exact tensor network repre-
sentation, we use its Suzuki-Trotter decomposition [70–72]
and VTNR [47,48] to find an iPEPS approximating \( |\Psi(T)\rangle \)
for a given bond dimension \( D \).

IV. VTNR

We treat an operator \( e^{-iH/T} \) as an imaginary-time evolu-
tion operator \( e^{-i\tau H} \) with the imaginary time \( \tau = 1/2T \). We de-
compose \( H \) into a sum of classical Hamiltonians \( H^{j\ell}_{cl} \), i.e.,
Hamiltonians which are sums of commuting terms,
\[
H = \sum_{j=1, \ldots ,m} H^j_{cl}.
\]
We use a second-order Suzuki-Trotter decomposition \([70–72]\) to approximate \(e^{-\tau H}\):

\[
e^{-\tau H} = \left( e^{-\delta \tau / 2 H_{11}^I} \cdots e^{-\delta \tau / 2 H_{11}^m} e^{-\delta \tau / 2 H_{11}^n} \right) \left( e^{-\delta \tau / 2 H_{22}^I} \cdots e^{-\delta \tau / 2 H_{22}^m} e^{-\delta \tau / 2 H_{22}^n} \right)^{1/\delta \tau} + O(\delta \tau^2). \tag{5}
\]

The accuracy of the decomposition (5) is controlled by the size of the small time step \(\delta \tau\). An exact iPEPO representation of an exponential of a classical Hamiltonian with a finite range of interaction can be found analytically and has \(D = d^2\) at most (see Ref. [44] and simple examples in Refs. [47,48]). Using the Suzuki-Trotter decomposition (5) and iPEPO representations of the exponentials of \(H_{11}^J\), we approximate \(|\Psi(T)\rangle\) [Eq. (4)] by the 3D tensor network shown in Fig. 2(a).

We use VTNR to approximate the 3D tensor network by an iPEPS with a numerically tractable bond dimension \(D\), which yields an approximate representation of \(|\Psi(T)\rangle\). In VTNR the iPEPS tensors are obtained by acting with tree tensor networks consisting of isometries \([73]\) on the virtual indices of the 3D tensor network, as shown in Fig. 2(b). The isometries are found by a variational update to minimally distort the partition function \(Z(T)\) as described in detail in Refs. [47,48]. The accuracy of the final iPEPS is controlled systematically by the bond dimension of the isometries, which here equals the bond dimension of the final iPEPS.

V. CORNER TRANSFER MATRIX RENORMALIZATION GROUP (CTMRG)

To compute expectation values of observables and the correlation length \(\xi\), we use CTMRG \([26,74–76]\). CTMRG approximates contractions of an infinite number of copies of the PEPS tensors by contractions of a finite number of environment tensors \(C, E\), where the accuracy is systematically controlled by the bond dimension \(\chi\) of the environment tensors. An example of such an approximation is shown in Fig. 2(a).
we define, analogously to (d), an operator

\[ E_\chi \approx \Delta^2 \]

which we use to compute \( \rho(T) \). By grouping the three left and the three right indices of the transfer matrix into single indices, we obtain a matrix which we use to compute \( \xi \) (see Sec. VI). In (e) we define, analogously to (d), an operator \( o \) transfer matrix \( T_o \), which is used in the \( \xi \) extrapolation procedure (see Sec. VI). In (c)–(e) for simplicity we show CTMRG for the translationally invariant (\( \Psi(T) \)). Furthermore, we assume a translational-invariant state with tensor \( A \) being rotational and mirror symmetric.

Figs. 3(a)–3(c). Details of the algorithm can be found in Refs. [26,48].

VI. \( \xi \) EXTRAPOLATION

The correlation length \( \xi \) converges very slowly with increasing \( \chi \), unlike local observables, e.g., the energy or the magnetization [65,77]. Therefore, to determine the correlation length \( \xi \) of \( \Psi(T) \) we use an extrapolation procedure from Ref. [77] which we summarize briefly below in the simplest case of a translationally invariant PEPS. The procedure uses the eigenvalues of the CTMRG transfer matrix \( T \) [see Fig. 3(d)].

To set up the extrapolation we define \( \epsilon_j \) for each eigenvalue \( \lambda_j \) of \( T \),

\[ \lambda_j/\lambda_0 = e^{-\epsilon_j+i\phi_j}, \]

where \( j \) numbers the eigenvalues of \( T \) ordered by the absolute value |\( \lambda_0 | \geq |\lambda_1 | \geq |\lambda_2 | \cdots \) and \( \phi_j \in (-\pi, \pi) \) determines the phase of \( \lambda_j \). A connected correlation function of a one-site operator \( o \) at a distance \( R \) is then expressed as

\[ C_o(R) = \langle o_o o_R \rangle = \sum_{j=0}^{N} f_j^{oo} e^{-\epsilon_j+i\phi_j} R^j. \]

Here, the form factors \( f_j^{oo} \) are defined as

\[ f_j^{oo} = \langle \phi_0 | T_{oo} | \phi_j \rangle (\phi_j | T_{oo} | \phi_0). \]

\( T_o \) is a transfer matrix of the operator \( o \) shown in Fig. 3(e), and \( | \phi_j \rangle \), \( ( \phi_j ) \) are left and right eigenvectors of \( T \) normalized as \( \langle \phi_0 | \phi_j \rangle = \delta_j \). Therefore, the correlation length obtained from \( T \) equals

\[ \xi_T = 1/\epsilon_1. \]

As the spectrum of \( T \) is continuous in the limit of \( \chi \to \infty \), the extrapolation uses its deviation from continuity \( \delta \) as a measure of finite-\( \chi \) effects. To extrapolate we use eigenvalues of \( T \) contributing to a connected correlation function of the phase transition’s order parameter \( C_{mm}(R) \), i.e., the eigenvalues with nonzero form factors of the order parameter \( f_j^{mm} \). We denote \( \epsilon \)'s of such eigenvalues by \( \{ \epsilon_k^{mm}, k = 1, 2, \ldots \} \) with \( \epsilon_1^{mm} \leq \epsilon_2^{mm} \leq \cdots \). We remark that in the case of a second-order phase transition, the diverging \( \xi \) is associated with the symmetry breaking, so we expect that the leading eigenvalue of the transfer matrix determines the asymptotics of the order parameter correlation function, i.e., \( \epsilon_1^{mm} = \epsilon_1 \). We observe that this is indeed the case for the transitions investigated below. We note that \( \lambda_0 \) does not contribute to \( C^{mm}(R) \) by definition. We define \( \delta \) as a distance in-between two dominant \( \epsilon_k^{mm} \), i.e.,

\[ \delta = \epsilon_2^{mm} - \epsilon_1^{mm}. \]

We note that this choice of \( \delta \) was proposed and benchmarked in Ref. [77]. Using the spectra from different values of \( \chi \), we extrapolate \( \epsilon_1^{mm} \) as a function of \( \delta \) by fitting

\[ \epsilon_1^{mm} = \epsilon_1^f + a \delta^b, \]

where \( \epsilon_1^f \) is an extrapolated value of \( \epsilon_1^{mm} \) and \( a, b \) are parameters of the fit. The extrapolation gives us

\[ \xi = 1/\epsilon_1^f. \]

In Fig. 4 we present examples of the extrapolation for the quantum Ising model with parameters which are investigated later in Sec. IX.

VII. FINITE CORRELATION LENGTH SCALING

Finite correlation length scaling (FCLS) was introduced for simulations of 1D quantum critical phenomena [63,64,67] with infinite MPS. In this case, a finite iMPS bond dimension \( D \) introduces a finite correlation length \( \xi_D \) at a critical point. \( \xi_D \)
This observation motivates us to consider an application of FCLS to obtain $T_c$ and the critical exponents from the VTNR results.

We obtain the FCLS Ansatz for observables from the standard finite-size scaling Ansatz by replacing the finite system size by $\xi_D$, e.g., for the order parameter $m(t, D)$ we use

$$m(t, D) = \xi_D^{-\beta/v} / \tilde{\mathcal{M}} \left( t \xi_D^{1/v} \right),$$

where $\beta, v$ are the critical exponents, and $\tilde{\mathcal{M}}$ is a nonuniversal function.

To compute observables we contract the iPEPS using CTMRG. The finite $\chi$ introduces an effective length scale $\xi_D$ [69], i.e., finite-$\chi$ values of the observables are given by a more complicated scaling Ansatz depending on both $\xi_D$ and $\bar{\chi}$, e.g.,

$$m(t, D, \chi) = \bar{\chi} \xi_D^{-\beta/v} / \tilde{\mathcal{M}} \left( t \xi_D^{1/v}, \xi_D / \bar{\chi} \right),$$

where $\tilde{\mathcal{M}}$ is another nonuniversal function. To avoid working with the more complicated Ansatz, we work in the limit of $\chi \to \infty$ as proposed in Refs. [65,66]. We observe that there is no need to extrapolate the order parameter in $\chi$ as it converges quickly. On the other hand, to obtain a good estimate of $\xi(t, D) \equiv \xi(t, D, \chi \to \infty)$ we use the $\xi$ extrapolation described in Sec. VI.

VIII. $T_c$ ESTIMATION

Finite-size scaling usually makes use of the Binder cumulant to locate the critical point without prior knowledge of its critical exponents, but in the case of iPEPS computation of the Binder cumulant is challenging because the fourth-order moment of the order parameter would need to be computed. Instead, we apply the $m'/m$ collapse introduced in [65], which makes use of the derivative of the order parameter $m'(t, D) = dm(t, D)/dt$, to find $T_c$. A FCLS scaling Ansatz for $m'$ is

$$m'(t, D) = \xi_D^{-\beta/(v-1)} / \tilde{\mathcal{M}}' \left( t \xi_D^{1/(v-1)} \right),$$

where $\tilde{\mathcal{M}}'$ is a nonuniversal function [65]. Using (13) and (15) we obtain

$$\bar{\mathcal{M}}' \left( t \xi_D^{1/(v-1)} \right),$$

$$\xi_D^{1/(v-1)} \propto \frac{m'(t = 0, D)}{m(t = 0, D)},$$

where $\bar{\mathcal{M}}$ is another nonuniversal function. Equations (16) and (17) give us the $m'/m$ collapse

$$\frac{m'(t, D)}{m(t, D)} = \bar{\mathcal{M}} \left( \frac{m'(t = 0, D)}{m(t = 0, D)} \right).$$

We estimate $T_c$ by plotting $y = t \left( \frac{m'(t, D)}{m(t, D)} \right)$ versus $x = t \left( \frac{m'(t = 0, D)}{m(t = 0, D)} \right)$ for different choices of $T_c$ and finding the one for which $y(x)$ data points obtained with different $D$ collapse best onto a single curve.
FIG. 5. The order parameter \( m \) in the vicinity of the critical temperature for the 2D quantum Ising model with \( \Gamma = 2.5 \) and \( D = 5–7 \). The critical behavior is smoothed out by finite-\( D \) effects. The black lines indicate the QMC estimate of the critical temperature \( T_c = 1.2737(6) \) [79].

IX. QUANTUM ISING MODEL: BENCHMARK RESULTS

The quantum Ising model is given by the Hamiltonian

\[
H = -\sum_{\langle i,j \rangle} \sigma_i^x \sigma_j^x + \Gamma \sum_i \sigma_i^z,
\]

where \( \sigma_x, \sigma_z \) are Pauli matrices. For \( \Gamma = 0 \) the model reduces to the classical Ising model with \( T_c = 2/\ln(1 + \sqrt{2}) \approx 2.269 \) and for \( \Gamma = 3.04438(2) \) [78] it has a quantum critical point. For \( 0 < \Gamma < \Gamma_c \) it exhibits a low-temperature ferromagnetic phase with the order parameter \( m = \langle \sigma_z \rangle \), which is separated from the paramagnetic phase by a line of finite-temperature second-order phase transitions belonging to a 2D classical Ising universality class.

As \( \Gamma \) is approaching \( \Gamma_c \), quantum fluctuations are becoming stronger and \( T_c \) gets suppressed with respect to the classical case of \( \Gamma = 0 \). Therefore, we expect that with increasing \( \Gamma \) the accurate simulation of the finite-temperature transition using tensor networks is becoming more challenging since a larger \( D \) is necessary to correctly capture the stronger quantum fluctuations. To examine this more closely, we investigate in the following \( \Gamma = 2.5 \) as well as a point close to the quantum critical point \( \Gamma = 2.9 \), for which quantum Monte Carlo (QMC) [79] gives \( T_c = 1.2737(6) \) and \( T_c = 0.6085(8) \), respectively, corresponding to a reduction in \( T_c \) with respect to \( \Gamma = 0 \) by a factor of 1.8 and 3.7, respectively.

A. \( \Gamma = 2.5 \)

We first consider a case well away from the quantum critical point \( \Gamma = 2.5 \) to provide a proof of principle of the applicability of FCLS to finite-temperature VTNR simulations. In Fig. 5 we present data for the order parameter as a function of temperature in the vicinity of the critical temperature, for bond dimensions \( D = 5–7 \). As expected, we do not obtain a sharp phase transition but we see that the order parameter is systematically reduced with increasing \( D \) similarly to the case of finite-size effects.

We first attempt to estimate \( T_c \) by using the known critical exponents of the 2D classical Ising universality class, i.e., \( \beta = 1/8 \) and \( \nu = 1 \). To do that we plot \( m(T) D^{\beta/\nu} \) which, according to FCLS, should not depend on \( D \) at \( T = T_c \):

\[
m(t,D) D^{\beta/\nu} = \mathcal{A}(t^{1/\nu}_D),\quad (20)
\]

\[
m(t = 0, D) D^{\beta/\nu} = a,\quad (21)
\]

where \( a \) does not depend on \( D \). Indeed, the \( m(t,D) D^{\beta/\nu} \) curves for different \( D \)’s cross as predicted by FCLS (see Fig. 6). We identify \( T_c \) as the temperature for which the variance of \( m(t,D) D^{\beta/\nu} \) is smallest, obtaining \( T_c = 1.2737(2) \) in agreement with the QMC estimate \( T_c = 1.2737(6) \) [79]. The \( T_c \) uncertainty is obtained by varying the range of \( D \) and the range of \( \chi \) used for \( \xi_D \) estimation by extrapolation.

To provide further evidence of FCLS we determine the critical exponents, using \( T_c = 1.2737(2) \) found earlier. We first estimate \( \beta/\nu \) using the data obtained at \( T = T_c \). From Eq. (21) we obtain

\[
\log m(t = 0, D) = -\beta/\nu \log \xi_D + \log a.\quad (22)
\]

A linear fit to the data on a log-log scale shown in Fig. 7(a) yields \( \beta/\nu = 0.123(15) \), in agreement with the exact \( \beta/\nu = 1/8 \). The error bar takes into account the \( T_c \) uncertainty and the statistical error of the fit.

Next, we estimate \( \beta \) and \( \nu \) by performing data collapses based on data in the vicinity of \( T_c \), using the scaling Ansätze:

\[
m(t,D) D^{\beta/\nu} = \mathcal{A}(t^{1/\nu}_D),\quad (23)
\]

\[
m(t,D)^{-\beta} = \mathcal{A}(t^{1/\nu}_D).\quad (24)
\]

Using the Ansätze (23) we obtain \( \beta = 0.127(2), \nu = 1.04(4) \), in agreement with the exact universality class [see Fig. 7(b)]. Ansätze (24) yields \( \beta = 0.126(5), \nu = 1.04(5) \), again in agreement with the exact exponents [see Fig. 7(c)]. In both cases, the uncertainties are obtained by taking into account the \( T_c \) uncertainty and by varying the data range.

Finally, we show that we can estimate \( T_c \) and \( \beta \) without prior knowledge of the universality class nor the value of \( T_c \).
respectively. The \( \beta \) is estimated \( \bar{\beta}_0 \) by performing data collapses using Eqs. (23) and (24) we obtain \( \beta/\nu = 0.123(15) \), in agreement with the exact \( \beta = 0.126(5) \), \( \nu = 1.045(4) \), respectively. The \( \beta \) and \( \nu \) estimates agree with the exact \( \beta = 1/8 \), \( \nu = 1 \).

First, we estimate \( T_c \) by performing a data collapse using the \( m'/m \) Ansatz (18), which yields \( T_c = 1.273(1) \) [see Fig. 8(a)]. We remark that while the quality of the obtained results is good, we see some deviations from perfect scaling which may be caused either by corrections to finite-size scaling or limitations of VTNR in getting optimal tensors. First, we note that VTNR is not guaranteed to give the best iPEPS approximation of the thermal state's purification for a given \( D \), as it does not search directly for the best iPEPS tensor representing thermal state purification. Instead, it optimizes a tree tensor network (TTN) of isometries which, applied to virtual indices of the tensor network representing a Suzuki-Trotter decomposition of the purification, gives the iPEPS approximating the purification. While this approach makes the variational optimization of the iPEPS efficient, it is not equivalent to the most general iPEPS variational optimization procedure. Still, our results demonstrate that the accuracy of the optimized tensors is high enough to extract the critical coupling and critical exponents with a good accuracy.

Second, we note that for \( \Gamma = 2.5 \) CTMRG convergence is challenging, as for \( D = 5–7 \) we obtain \( \xi_{D=7} \sim 350 \). For an iPEPS with such a large \( \xi \) many iterations of the CTMRG procedure are necessary to converge \( m \). Good convergence

\begin{equation}
\langle T_c = 1.273 \rangle
\end{equation}

\begin{equation}
\langle m'(t) / m(t) \rangle
\end{equation}

\begin{equation}
\langle m(t) \rangle
\end{equation}

Taking into account the \( T_c \) uncertainty and varying the data range.

FIG. 8. \( T_c \) and \( \beta \) estimation for \( \Gamma = 2.5 \) without knowledge of the universality class nor the value of \( T_c \). In (a) we use the \( m'/m \) Ansatz (18) to perform a data collapse, obtaining \( T_c = 1.273(1) \). In (b) we estimate \( \beta \) by performing a data collapse based on Eq. (25), using the value of \( T_c \) found in (a), which yields \( \beta = 0.12(1) \). The obtained \( T_c \) and \( \beta \) agree with the QMC estimate \( T_c = 1.2737(6) \) [79] and the exact \( \beta = 1/8 \).
of the CTMRG environment is important for the variational optimization since VTNR uses the CTMRG environmental tensors to find the best iPEPS.\footnote{To ensure a good CTMRG convergence, we require the change of $m$ per CTMRG iteration to be smaller than $10^{-8}$.} Here, to optimize the iPEPS we use $\chi = 8D$. We check that using $\chi = 6D$ and $\chi = 7D$ we obtain results [$T_c$ and $\beta$ obtained by the $m'/m$ collapse and collapse (25)] in agreement with the ones obtained with $\chi = 8D$. Nevertheless, we cannot fully exclude the possibility that finite-$\chi$ effects contribute to the observed small deviations from the perfect collapse as simulations with larger $\chi$ would be computationally very expensive.

To obtain $\xi_D$ we contract the final iPEPS with $8D \leq \chi \leq 24D$ and use the extrapolation procedure described in Sec. VI. We remark that in the case of VTNR simulations obtaining convergence in the small Trotter time step $\delta \tau$ is relatively easy as the computational cost of the simulations scales at most as $O(\log(1/\delta \tau))$. For $\Gamma = 2.5$ we use a second-order Trotter decomposition with $\delta \tau = \tau/2^{10} \leq 0.001$, which is small enough to give results converged in $\delta \tau$.

\section{B. $\Gamma = 2.9$}

For the more challenging $\Gamma = 2.9$ case we analyze VTNR results for $D = 7–10$ [see Fig. 9(a)]. Using the $m'/m$ collapse we obtain $T_c = 0.609(4)$ [see Fig. 9(b)]. Using this result for $T_c$ and performing a collapse with Eq. (25) we obtain $\beta = 0.152(8)$ [see Fig. 9(c)]. While the obtained $T_c$ estimate agrees with the QMC estimate $T_c = 0.6085(8)$ \cite{79}, the $\beta$ estimate deviates by about 20% from the exact $\beta = \frac{1}{3}$.

Comparing the results obtained with $D = 8–10$ and $D = 7–10$ suggests that the $\beta$ estimate still depends significantly on the $D$ range as we obtain $\beta = 0.145(5)$ for $D = 8–10$.\footnote{Contrary to the case of $\beta$, the $T_c$ estimate for $D = 8–10$, $T_c = 0.610(4)$, is similar to the $D = 7–10$ estimate.} Furthermore, we expect that the necessary $\xi_D$ to obtain the asymptotic scaling is larger for $\Gamma = 2.9$ than for $\Gamma = 2.5$, as $\Gamma = 2.9$ is closer to the quantum critical point. Despite larger $D$, the $\xi_D$ obtained for $\Gamma = 2.9$, although quite large ($\xi_D \sim 20–50$), is smaller than $\xi_D$ for $\Gamma = 2.5$ ($\xi_D \sim 40–350$). Therefore, we expect that the quality of the results can still be improved by increasing $D$, although it would be computationally very expensive.

The $T_c$ and $\beta$ uncertainties are estimated in the same way as in the $\Gamma = 2.5$ case. Similar values for $T_c$ and $\beta$ are obtained with VTNR using an optimization with $\chi = 5D, 6D, 7D$. A second-order Trotter decomposition with a time step $\delta \tau = \tau/2^{10} \leq 0.002$ is used.

\section{X. INTERACTING HONEYCOMB FERMIONS: BENCHMARK RESULTS}

We consider a model of interacting spinless fermions on a honeycomb lattice \cite{80}, given by the Hamiltonian

\begin{equation}
H = -t \sum_{\langle i,j \rangle} (c_i^\dagger c_j + c_j^\dagger c_i) + V \sum_{\langle i,j \rangle} n_i n_j. \tag{26}
\end{equation}

Here, $c_i$ ($c_i^\dagger$) is a fermionic annihilation (creation) operator at site $i$ and $n_i = c_i^\dagger c_i$ is a fermion number operator. We set $t = 1$ in the following. Furthermore, for the purpose of the benchmark we restrict ourselves to the case of half-filling, $n = \sum_{i=1}^N \langle n_i \rangle / N = \frac{1}{2}$, for which sign-problem free QMC results are available \cite{79,81}. The model has a quantum critical point at $V_c = 1.356(1)$ \cite{81}. For $V > V_c$ there is a low-temperature phase with a charge density wave (CDW) order,
The CDW order parameter is defined as
\[
\langle n_\alpha \rangle - \langle n_\beta \rangle,
\]
(27)
where \(\langle n_\alpha \rangle\) and \(\langle n_\beta \rangle\) are the fermion densities on sub-lattices \(A\) and \(B\), respectively. In the limit of \(V \to \infty\) the model becomes equivalent to the 2D classical antiferromagnetic Ising model. Here, we simulate the model for \(V = 3\) and the more challenging case of \(V = 2\), which is closer to \(V_c\).

A. \(V = 3\)

For \(V = 3\) we analyze \(D = 12, 14, 16\) VTNR results in the vicinity of \(T_c\) [see Fig. 10(a)]. Here, QMC predicts \(T_c = 0.936(10)\) [79]. We determine \(T_c\) using the \(m'/m\) collapse obtaining \(T_c = 0.939(4)\) in agreement with QMC [see Fig. 10(b)]. Furthermore, we obtain \(\beta = 0.132(8)\) in agreement with the exact \(\beta = \frac{1}{8}\) by performing a data collapse using Eq. (25) and by taking \(T_c\) obtained from the \(m'/m\) collapse [see Fig. 10(c)].

The \(T_c\) and \(\beta\) uncertainties are obtained similarly as for the quantum Ising model. We use \(\chi = 5D\) to perform the VTNR optimization obtaining results which are consistent with the ones obtained with \(\chi = 3D\) and \(\chi = 4D\). We use a second-order Suzuki-Trotter decomposition with a time step \(\delta \tau = \tau/2^{11} < 0.001\).

B. \(V = 2\)

Next, we analyze \(D = 12, 14, 16\) VTNR results for the more challenging \(V = 2\) case [see Fig. 11(a)]. Using the \(m'/m\) collapse we obtain \(T_c = 0.480(5)\) in agreement with the QMC
estimate $T_c = 0.47(1)$ [79] [see Fig. 11(b)]. Using Ansatz (25) and taking $T_c$ found by the $m'/m$ collapse, we obtain $\beta = 0.108(4)$ [see Fig. 11(c)]. The $\beta$ estimate deviates by about 10% from the exact $\beta = \frac{1}{3}$. We see that the obtained $\xi_D$ for $V = 2$ ($\xi_D \sim 2-4$) is much smaller than $\xi_D$ for $V = 3$ ($\xi_D \sim 10-20$). Furthermore, as $V = 2$ is closer to the quantum critical point than $V = 3$ we expect that for $V = 2$, a larger $\xi_D$ is necessary to be in the asymptotic scaling regime and we expect that the accuracy of $\beta$ can be improved by increasing $D$.

The $T_c$ and $\beta$ uncertainties are obtained similarly as for the quantum Ising model. Here, we use $\chi = 5D$ to perform the VTNR optimization. Consistent results are obtained also with $\chi = 3D$ and $4D$. We use a second-order Suzuki-Trotter decomposition with a time step $\delta t = \tau/2^8 < 0.01$.

**XI. Conclusions**

In this paper we have studied second-order finite-temperature phase transitions in the 2D quantum Ising [Eq. (19)] and interacting honeycomb fermion [Eq. (26)] models using infinite projected entangled pair states (iPEPS) to represent thermal states. The iPEPS were obtained by variational tensor network renormalization (VTNR). We found that at the critical temperature $T_c$, the iPEPS correlation length $\xi_D$ is finite for the computationally accessible values of the iPEPS bond dimension $D$. Motivated by this observation, we investigated the application of finite correlation length scaling (FCLS) to obtain precise values of $T_c$ and universal critical exponents. We found that in the vicinity of $T_c$, the order parameter obeys well the expected behavior predicted by FCLS.

The two models studied in this work exhibit second-order finite-temperature phase transitions for the transverse fields $\Gamma < \Gamma_c \approx 3.04438$ and the interaction strengths $V > V_c \approx 1.356$, respectively. At $\Gamma_c$ and $V_c$, second-order quantum phase transitions occur at $T = 0$. Using FCLS we obtained estimates of $T_c$ and the critical exponents in agreement with the QMC results for $\Gamma = 2.5$ and $V = 3$ which are sufficiently far from the quantum critical points. For $\Gamma$ and $V$ approaching the quantum critical points, we observed that the magnitude of $\xi_D$ and the accuracy of the critical data become lower for the same values of $D$. Nevertheless, we were still able to obtain $T_c$ in agreement with the QMC results for the challenging $\Gamma = 2.9$ and $V = 2$ cases. For these couplings, the values of the critical exponents exhibit a dependence on the range of $D$ values used in the scaling analysis, suggesting that larger $D$’s are needed in order to obtain more accurate estimates of the critical exponents.

In summary, our results further demonstrate the usefulness of tensor network simulations for quantum many-body systems at finite temperature, even for the challenging case of a finite-temperature continuous phase transition, for which convergence in $D$ can typically not be reached, but which can be systematically studied using FCLS.

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