Topological gap opening without symmetry breaking from dynamical quantum correlations

F. Paoletti [1,*] L. Fanfarillo [2,1] M. Capone [1,3] and A. Amaricci [3]

[1]*International School for Advanced Studies (SISSA), via Bonomea 265, 34136 Trieste, Italy*

[2]*Istituto dei Sistemi Complessi (ISC-CNR), Via dei Taurini 19, I-00185 Rome, Italy*

[3]*CNR-IOM, Istituto Officina dei Materiali, Consiglio Nazionale delle Ricerche, Via Bonomea 265, 34136 Trieste, Italy*

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Topological phase transitions are typically associated with the formation of gapless states. Spontaneous symmetry breaking can lead to a gap opening, thereby obliterating the topological nature of the system. Here we highlight a completely different destiny for a topological transition in the presence of interaction. Solving a Bernevig-Hughes-Zhang model with local interaction, we show that dynamical quantum fluctuations can lead to the opening of a gap without any symmetry breaking. As we vary the interaction and the bare mass of the model, the continuous gapless topological transition turns into a first-order one, associated with the presence of a massive Dirac fermion at the transition point, showing a Gross-Neveu critical behavior near the quantum critical endpoint. We identify the gap opening as a condensed matter analog of the Coleman-Weinberg mechanism of mass generation.

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I. INTRODUCTION

The discovery of symmetry-protected topological phases of matter [1–8] has enriched the landscape of phase transitions beyond the conventional Landau paradigm of the symmetry breaking [9,10]. In the presence of a given symmetry, the possible electronic band structures of an insulator can be divided into distinct equivalence classes, which can only be connected through the continuous closure of the energy gap from both sides of the transition through a topological quantum phase transition (TQPT). The corresponding formation of symmetry-protected massless Dirac fermions at the transition is a distinctive feature of the topological insulators [11–15].

The presence of interactions can change the above scenario and a gapped state can also appear at the transition. The standard phenomenology for this to occur requires a spontaneous symmetry breaking (SSB) [4,16–20]. Indeed, breaking a continuous symmetry opens a gap in the energy spectrum or, equivalently, gives a finite mass for the Dirac electrons which is understood in terms of the Anderson-Higgs mechanism. Clearly, a SSB can lead to break any of the symmetry protecting the topological state, thus leaving behind a topologically trivial long-range ordered phase. A similar scenario can be described within a static mean-field (MF) picture in the channel where SSB takes place.

A great deal of attention has been recently drawn to different fields [21–23] to possible mechanisms of spontaneous mass generation which preserve the symmetry, beyond the conventional SSB description. Here we show that such a process describes the gap opening for Dirac electrons at the boundary of a topological insulator. More concretely, we address the question of whether or not electron-electron interactions can drive the formation of a spontaneous mass for the otherwise gapless electrons at a topological transition. The lack of a gap closing at the TQPT is expected to change the character of the transition, which becomes necessarily discontinuous despite the symmetries protecting the topological phase being preserved.

For the sake of definiteness, we consider a two-dimensional Bernevig-Hughes-Zhang (BHZ) model augmented via the inclusion of local electron-electron interactions that preserve some symmetries of the model. Without interactions, this model features a TQPT through the formation of a gapless state. The control parameter of the transition is the energy splitting between two electronic orbitals playing the role of a mass term. As a consequence, the difference in the occupation of the orbitals, or orbital polarization, is expected to assume a prominent role. Since the orbital symmetry is broken by the mass term, the concept of SSB does not apply to the TQPT. Within a MF theory, the interactions simply dress the mass term, shifting the topological transition without changing its nature with respect to the noninteracting limit.

In this paper, we go beyond MF using a variational approach including quantum fluctuations not only of the orbital polarization but also in the other particle-hole channels. We demonstrate a scenario in which a gap opens at the TQPT without breaking any of the symmetries of the model. We show explicitly that the quantum fluctuations in the different channels make the TQPT discontinuous for sufficiently large interactions [24,25]. The first-order line ends in a critical endpoint, where we show a Gross-Neveu quantum critical behavior as a function of the relevant coupling strength [26]. As we shall discuss in the following, the mechanism we revealed is reminiscent of the Coleman-Weinberg (CW) theory of mass generation [27,28].

The rest of the paper is organized as follow. In Sec. II, we introduce the model and set up the starting point of our analysis. We address the MF solution of the model in Sec. III, while in Sec. IV we develop a nonperturbative method to investigate the quantum fluctuations near the TQPT. In Sec. V, we show

*francesca.paoletti@sissa.it

the existence of a spontaneous gap formation and characterize its critical behavior. Finally, in Sec. VI we summarize our results.

II. MODEL AND METHOD

We solve an interacting BHZ model on a square lattice [1,2,24,29,30],

$$\mathcal{H} = \sum_{\mathbf{k}} \psi^{\dagger}_{\mathbf{k}} H^{0}_{\mathbf{k}} \psi_{\mathbf{k}} + \sum_{\mathbf{i}} \mathcal{H}^{\mathrm{int}}_{\mathbf{i}}, \tag{1}$$

where $\psi_{\mathbf{k}} = [c_{\mathbf{k}1\uparrow}, c_{\mathbf{k}2\uparrow}, c_{\mathbf{k}1\downarrow}, c_{\mathbf{k}2\downarrow}]^{T}$ and the operators $c_{\mathbf{k}\alpha\sigma}$ annihilate an electron with momentum $\mathbf{k}$, orbital $\alpha = 1, 2$, and spin $\sigma = \uparrow, \downarrow$. If we define $\Gamma_{\mu\nu} = \sigma_{\mu} \otimes \tau_{\nu}$, with σ_{μ} and τ_{ν} the Pauli matrices, respectively, in the spin and orbital subspaces the single-particle Hamiltonian reads $H^{0}_{\mathbf{k}} = [M - 2t(\cos(k_x) + \cos(k_y))]\Gamma_{03} + \lambda \sin(k_x)\Gamma_{31} + \lambda \sin(k_y)\Gamma_{02}$, where $M \geqslant 0$ is the energy separation between the two orbitals which plays the role of the mass term, and t and λ are the intra- and interorbital hopping amplitudes. The model is invariant under time-reversal $\mathcal{T}$ and inversion $\mathcal{P}$ symmetries, $U(1)$ spin rotation around the z axis [20,31–33]. In the following, we set our energy unit so $2t = 1$ and focus on the regime of two electrons per site, i.e., half filling. The noninteracting model has a continuous topological transition between a quantum spin Hall insulator (QSHI) for $M < 2$ and a trivial band insulator (BI) for $M > 2$ through the formation of a gapless Dirac state at $M = 2$ [1].

We assume a generic local interaction which preserves inversion $\mathcal{P}$ and $U(1)$ spin symmetry around the z axis [18,34],

$$\mathcal{H}^{\mathrm{int}}_{\mathbf{i}} = -\frac{g_N}{2}\hat{N}^{2}_{\mathbf{i}} - \frac{g_T}{2}\hat{T}^{2}_{z\mathbf{i}} - \frac{g_S}{2}\hat{S}^{2}_{z\mathbf{i}} - \frac{g_R}{2}\hat{R}^{2}_{z\mathbf{i}}, \tag{2}$$

where $\hat{N}_{\mathbf{i}} = \frac{1}{2}\psi^{+}_{\mathbf{i}}\Gamma_{00}\psi_{\mathbf{i}}$ is half of the total occupation per site, $\hat{T}_{z\mathbf{i}} = \frac{1}{2}\psi^{+}_{\mathbf{i}}\Gamma_{03}\psi_{\mathbf{i}}$ and $\hat{S}_{z\mathbf{i}} = \frac{1}{2}\psi^{+}_{\mathbf{i}}\Gamma_{30}\psi_{\mathbf{i}}$ are, respectively, the z component of the orbital polarization and the spin operators and $\hat{R}_{z\mathbf{i}} = \frac{1}{2}\psi^{+}_{\mathbf{i}}\Gamma_{33}\psi_{\mathbf{i}}$; $\psi_{\mathbf{i}}$ is the Fourier transform of $\psi_{\mathbf{k}}$. In the numerical calculations, we will consider $g_N = -(3U - 5J)$, $g_T = U - 5J$, $g_S = U + J$, and $g_R = U - J$ to recover the density-density version of the popular Kanamori-Hubbard [35] model used in a variety of works to study the interplay between the Hubbard U and the Hund's exchange J and their effect on TQPTs [24,36,37]. For the sake of simplicity, here we disregard the contributions of spin-flip and pair-hopping terms, which do not qualitatively modify the conclusions of this paper. We consider nonmagnetic solutions to study $\mathcal{T}$ symmetry-preserving transitions [38,39]. To simplify the notation, in the following we define $\Gamma_{a=N,T,S,R}$ as the set $\frac{1}{2}\Gamma_{\mu\nu=00,03,30,33}$ to highlight the different channels.

The starting point of our analysis is to rewrite the partition function of the interacting model Eq. (1) in terms of an effective problem coupled to space- and time-dependent (real) bosonic fields $\mathbf{\Delta}_q = \Delta^{a=N,T,S,R}_q$, by performing a Hubbard–Stratonovich (HS) transformation [40]. We emphasize that our results are general with respect to the change of signs of g_T occurring for values of $J/U > 1/5$. Although a negative coupling to the HS decoupling would result in a purely imaginary field, we can formally reduce to the case of real bosonic fields taking into account the negative part of the coupling, as for instance implicitly done for the charge channel term.

The partition function reads $\mathcal{Z} \equiv e^{-\beta\mathcal{F}} = \int \mathcal{D}\mathbf{\Delta}\, e^{-\beta\mathcal{N}F[\mathbf{\Delta}]}$ in terms of the free-energy functional

$$F[\mathbf{\Delta}] = \sum_{aq} \frac{|\Delta^{a}_{q}|^{2}}{2g_a} - \frac{1}{\beta\mathcal{N}}\mathrm{Tr}\ln(-\mathcal{G}^{-1}_{kq}), \tag{3}$$

where $q = (\mathbf{q}, i\nu_m)$, $k = (\mathbf{k}, i\omega_n)$, with $\mathbf{q}$, $\mathbf{k}$ wave vectors in the first Brillouin zone (BZ) and ν_m, ω_n the bosonic and fermionic Matsubara frequencies, respectively, β is the inverse temperature and $\mathcal{N}$ the total number of sites. Tr indicates the trace over momentum, frequency, orbital, and spin. Finally, $\mathcal{G}_{kq} = (i\omega_n + \mu - H^{0}_{\mathbf{k}}\delta_{\mathbf{k},\mathbf{k}-\mathbf{q}} - V_q)^{-1}$ is the interacting one-body Green's function, where $V_q = -\sum_a \Delta^{a}_{q}\Gamma_a$ is an effective time-dependent potential depending on $\mathbf{\Delta}_q$.

III. MEAN FIELD

The natural lowest-order approximation of Eq. (3) is a static MF solution, where the bosonic fields Δ^{a}_{q} are approximated with time-independent and spatially uniform quantities. The presence of the mass term explicitly breaks the symmetry between the orbitals, leading to a finite values of the orbital polarization T_z already in the noninteracting model. The nonmagnetic solution ($\Delta^{S}_{\mathrm{MF}} = \Delta^{R}_{\mathrm{MF}} = 0$) at half filling ($\Delta^{N}_{\mathrm{MF}} = 1$), reduces to the single self-consistency equation $\Delta^{T} = \frac{g_T}{\beta\mathcal{N}}\mathrm{Tr}(G^{\mathrm{MF}}_{k}\Gamma_T)$ where $G^{\mathrm{MF}}_{k} = (i\omega_n - H^{0}_{\mathbf{k}} + \Delta^{T}\Gamma_T)^{-1}$ is the MF Green's function. Thus, the MF solution simply corrects the mass term M so the model describes a continuous TQPT occurring at the critical line $M - \frac{1}{2}\Delta^{T}_{\mathrm{MF}} = 2$ [41] as reported in the phase diagram of Fig. 1(b). All along this line, the energy gap closes through the formation of a gapless Dirac node at the Γ point as in the noninteracting model. As we show in Fig. 1(a) (dotted line), the orbital polarization smoothly evolves across the topological transition. Coherently with the above scenario, a direct inspection of the MF free energy shows only one minimum for every value of g_T and M.

IV. CORRECTIONS FROM QUANTUM FLUCTUATIONS

In this paper, we overcome the limitations of the MF by approximating the exact free-energy functional with a second-order expansion in the fluctuating fields, whose coefficients are variationally chosen [42]. This approach enables us to derive a closed expression for the covariance matrix as the random potential V_q has a Gaussian statistics. We underline that our methodology significantly differs from conventional perturbation theory in which the expansion is performed around the MF solution [43]. As a matter of fact, our approach thoroughly incorporates all higher-order fluctuation corrections by considering a variationally determined dressed propagator. This corresponds to systematically correct the expectation values of the fields around which we perform the expansion: $\bar{\mathbf{\Delta}} \neq \bar{\mathbf{\Delta}}_{\mathrm{MF}}$. This latter point turns out to be crucial. Assuming $\mathbf{\Delta}_q \to \bar{\mathbf{\Delta}} + \mathbf{\Delta}_q$, we can write

$$F[\mathbf{\Delta}] \simeq F^{(2)}[\mathbf{\Delta}] = F[\bar{\mathbf{\Delta}}] + \frac{1}{2}\sum_{abq} \Delta^{a}_{q} A^{ab}_{q} \Delta^{b}_{-q}, \tag{4}$$

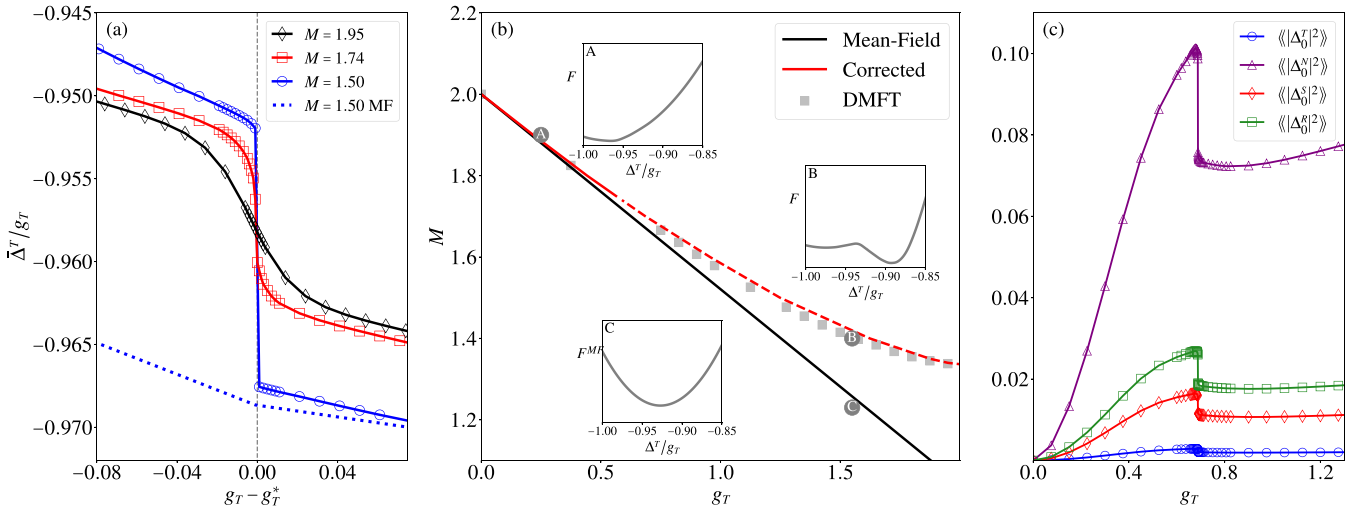

FIG. 1. (a) Orbital polarization $\bar{\Delta}^T/g_T$ as a function of g_T measured with respect to the TQPT point g_T^*. The open symbols correspond to the fluctuation-corrected results. The dotted line is the MF solution for $M = 1.50$. (b) Phase diagram in the M-g_T plane comparing the topological transition line in the two approximations. The solid lines (black and red) denote a continuous TQPT, while the dashed line (red) marks a discontinuous one. Data from DMFT are indicated with filled symbols (gray). Insets A, B, and C show the free energy F as a function of the orbital polarization for the three points marked on the curves. (c) Static and homogeneous components $\langle\langle|\Delta^a_{q=0}|^2\rangle\rangle$ for $a = T, N, S, R$ of the potential fluctuations as a function of g_T across the topological transition for $M = 1.70$.

where the variational principle $\mathcal{F} \leqslant \mathcal{F}^{(2)} + \langle\langle F[\mathbf{\Delta}] - F^{(2)}[\mathbf{\Delta}]\rangle\rangle$, with $\mathcal{F}^{(2)} = -\frac{1}{\beta}\ln\int \mathcal{D}\mathbf{\Delta}\, e^{-\beta\mathcal{N}F^{(2)}[\mathbf{\Delta}]}$, leads to the new stationary condition $\langle\langle\partial_{\Delta^a_q} F[\mathbf{\Delta}]\rangle\rangle = 0$ and $A^{ab}_q = \langle\langle\partial_{\Delta^a_q}\partial_{\Delta^b_{-q}} F[\mathbf{\Delta}]\rangle\rangle$. The symbol $\langle\langle\cdot\rangle\rangle$ indicates that the averages over the possible configurations are calculated using the second-order probability density of the fluctuating field.

The coefficients A^{ab} in the free-energy expansion Eq. (4) depend on the averaged dressed Green's function $\langle\langle\mathcal{G}_{kq}[\mathbf{\Delta}]\rangle\rangle$, which cannot be calculated exactly. To circumvent this problem, we can introduce an auxiliary potential Σ_k implicitly determined by the condition [44,45] $\frac{\partial\mathcal{F}^{(2)}}{\partial\Sigma_k} = 0$, which in turn implies that $\langle\langle\mathcal{G}_{kq}[\mathbf{\Delta}]\rangle\rangle$ coincides with an interacting Green's function in which Σ_k plays the role of a self-energy $G_k = [i\omega_n + \mu - H_0(\mathbf{k}) - \Sigma_k]^{-1}$.

The stationary condition and the expression for A^{ab}_q become

$$\frac{\bar{\Delta}^a}{g_a} = \frac{1}{\beta\mathcal{N}}\mathrm{Tr}(G_k\Gamma_a); \quad A^{ab}_q = \frac{\delta_{ab}}{g_a} - \chi_{ab}(q). \tag{5}$$

The first expression contains the fluctuation-corrected Green's function, hence it leads to corrected values of the $\bar{\mathbf{\Delta}}$, while the second can be seen as an optimized version of the random-phase approximation (RPA), as we shall discuss in the following [42]. Indeed, $\chi_{ab}(q) = -\frac{1}{\beta\mathcal{N}}\sum_k \mathrm{Tr}[G_k\Gamma_a G_{k+q}\Gamma_b]$ is the susceptibility matrix in the space of the different channels which, in the presence of odd hybridization between the orbitals, has diagonal structure $\chi_{aa}\delta_{ab}$. Moreover, the symmetries of the interaction ensure that $\chi_{SS} = \chi_{NN}$ and $\chi_{RR} = \chi_{TT}$. Σ_k can be written explicitly up to second order as

$$\Sigma_k = \bar{V} + \sum_{qa} G_{k-q}\left\langle\!\left\langle\left|\Delta^a_q\right|^2\right\rangle\!\right\rangle, \tag{6}$$

where

$$\left\langle\!\left\langle\left|\Delta^a_q\right|^2\right\rangle\!\right\rangle = \frac{1}{\beta\mathcal{N}}\left[\frac{1}{g_a^{-1} - \chi_{aa}(q)} - g_a\right]. \tag{7}$$

The diagonal form in the channel index allows us to analyze the contribution of each fluctuating term of the interaction to the potential Σ_k [46]. We emphasize that, formally, the second term in Eq. (6) plays the same role of the one-loop quantum correction of the effective CW potential [27].

Equations (5) and (6) provide a closed system of nonlinear equations for Σ_k and the bosonic fields $\bar{\mathbf{\Delta}}$. Previous studies suggest that the interaction effects on the TQPT are mainly local and that nonlocal fluctuations play a minor role [47]. Thus, to further simplify the treatment, in the following we will assume a local $\Sigma_k \simeq \Sigma(i\omega_n)$. We solve this system iteratively using a linear mixing algorithm which typically converges in $10 - 20$ steps. The BZ is discretized with a linear grid of 20×20 points and the Matsubara axis with $L = 8192$ frequencies using an effective inverse temperature $\beta = 500$. The convolution in Eq. (6) is evaluated using a fast Fourier transform algorithm. We discuss the results obtained for $J/U = 1/8$ and $\lambda = 0.3$.

In Fig. 1(a), we show the evolution of the orbital polarization, obtained from the self-consistent value of the bosonic field $\bar{T}_z = \bar{\Delta}^T/g_T$. The behavior at the transition point $g_T^*(M)$ changes qualitatively according to the value of the bare mass M. For a value close to $M = 2$, i.e., the noninteracting transition point, the orbital polarization is continuous with respect to the increasing interaction g_T. This corresponds to a smooth modification of the BI into a nontrivial insulator through the formation of a gapless state at the TQPT. Starting from a farther point, the orbital polarization displays a critical behavior at the transition characterized by a divergent susceptibility $\partial_M\bar{T}_z$. Beyond this point, for any value of M, the orbital polarization is characterized by a discontinuous evolution across the topological transition. This is in stark contrast with the continuous behavior obtained in MF for the same value of the mass, see Fig. 1(a), and it agrees with previous results obtained via dynamical mean-field theory (DMFT) [24]. The

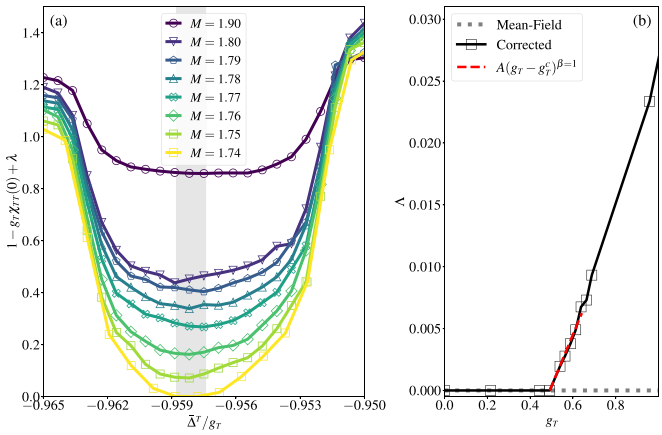

FIG. 2. (a) The denominator of Eq. (8) as a function of the orbital polarization $\bar{\Delta}^T/g_T$ across the TQPT in the fluctuation-corrected approximation and for different values of the bare mass M. The narrow grey stripe indicates the minimum with its numerical uncertainty. (b) The gap Λ along the TQPT line as a function of the interaction g_T. MF is the dotted grey line, while fluctuation-corrected results are indicated by open symbols and solid line. The (red) dashed line is a linear fit $A(g_T - g_T^c)^{\beta=1}$ ($A \simeq 0.042$) of the critical behavior.

agreement with DMFT is indeed even quantitative, as shown by the data reported in Fig. 1(b).

These results can be summarized in a phase diagram in the plane g_T-M [see Fig. 1(b)], where we compare the MF and fluctuation-corrected results for the TQPT. The two transition lines remain close for small values of the interaction g_T. Accordingly, the free-energy functional F displays a single minimum as a function of Δ^T [inset A in Fig. 1(b)]. However, upon increasing the interaction strength, the two curves start deviating significantly, signaling a crucial impact of the fluctuations.

A direct information about the contribution of the fluctuations in the different channels is reported in Fig. 1(c). The fast increasing behavior with the interaction g_T in all the channels stops at the topological transition towards the QSHI, where these quantities display a discontinuous drop and a successive slow increase. The terms $\langle\langle|\Delta_q^a|^2\rangle\rangle$ enter, through Eq. (7), in $\Sigma(i\omega_n)$ giving it a dynamical nature which significantly deviates from its static MF form $\bar{V} = -\bar{\Delta}^T\Gamma_T$. This results in a crucial shift of the self-consistent saddle point value of the bosonic fields. Moreover, as discussed above, while the MF always describes a continuous transition, in the corrected theory the boundary line is continuous up to a critical value g_T^c of the interaction beyond which it becomes of first order. This reflects in the behavior of the free energy near the TQPT point in the intermediate to strong coupling regime, i.e., $g_T > g_T^c$. In insets B and C of Fig. 1(b), we compare the free energies of a QSHI state near the topological transition for, respectively, the MF and the fluctuation-corrected approximation, where two minima are found: a stable QSHI and a metastable BI.

V. CRITICAL BEHAVIOR

We thus find that a quantum critical point (QCP) separates the continuous from the discontinuous regime on the topological transition line, where we also found a divergent orbital susceptibility. Indeed, the uniform orbital susceptibility $\mathcal{X}_T = \partial_M \bar{\Delta}^T_{q=0}$ reads, using Eq. (5),

$$\mathcal{X}_T = \frac{-2g_T\chi_{TT}(0)}{1 - g_T\chi_{TT}(0) + \lambda}, \tag{8}$$

which reminds us of the RPA result with a correction λ that stems from the implicit dependence of $\Sigma(i\omega_n)$ on the orbital polarization. This quantity also accounts for the contributions of all the other channels of the interaction through the expression of $\Sigma(i\omega_n)$, see Eq. (6) [48]. At the MF level, we find $\Sigma(i\omega_n) = -\bar{\Delta}^T\Gamma_T$, so $\lambda = 0$ and $\mathcal{X}_T$ reduces to the RPA form. For a Hubbard-Kanamori interaction, the RPA $\mathcal{X}_T$ diverges only either for negative U or negative M (which lead to different physics), in agreement with the continuous TQPT we always find. Note that this result also holds true when all the coupling constants but g_T vanish and the interaction reduces to $\frac{g_T}{2}\hat{T}_z^2$ [49].

To compute λ, we consider the zero-frequency limit of $\Sigma(i\omega_n)$ where we obtain $\lambda \simeq \chi_{TT}\partial_{\bar{T}_z}\Lambda$. Since $\partial_{\bar{T}_z}\Lambda < 0$, this correction is negative and tends to enhance $\mathcal{X}_T$. The significance of the negative λ factor, as well as the crucial role of the dynamical effects from all interaction channels on the orbital polarization, becomes even more apparent in the repulsive scenario, i.e., $g_T < 0$. In fact, in this case, without including λ as a form of coupling renormalization, it would be impossible for the denominator in Eq. (8) to vanish altogether. This further clarifies why the discontinuity was unexpectedly first observed in the $J/U > 1/5$ regime [24].

For a fixed value of M, the TQPT corresponds to the maximum of the response function which connects the continuous transition to a Widom line [50–52]. The minimum of the denominator in Eq. (8) approaches zero when we reach the QCP, as shown in Fig. 2(a). As we discussed above, in a noninteracting TQPT, the spectral gap closes at the transition. We now compute the gap in our scheme from the zero-frequency limit of the self-energy $\Lambda = \mathrm{Re}\Sigma(i\omega_n \to 0) - \bar{V}$. In Fig. 2(b), we report the behavior of Λ at the TQPT as a function of the interaction strength g_T. While in MF the gap is always zero, including the fluctuations we find a finite gap above the QCP ($g_T > g_T^c$). A finite value of the gap Λ corresponds to give a mass to the Dirac fermions at the boundary line. This is consistent with a spontaneous symmetric mass generation process [21–23]. The presence of such finite gap (or mass) makes it impossible to continuously connect the trivial with the nontrivial phase and leads to a first-order TQPT. In addition, we find numerically that the critical behavior near the QCP falls in the Gross-Neveu universality class [22,26], with an estimated critical exponent $\beta \simeq 1$.

VI. CONCLUSIONS

In this paper, using a nonperturbative analytical approach to include interactions in the BHZ model, we have demonstrated the crucial role of fluctuations in the different local particle-hole channels to qualitatively change the nature of the TQPT. Within a MF, the interactions only lead to a renormalization of the bare mass of the model (coupled to the

orbital polarization) so the TQPT has the same character of the noninteracting model. Within our approach, the fluctuation contributions change the MF parameters and lead to a discontinuous TQPT for large interactions with a QCP separating the continuous and discontinuous branches. This effect is intrinsically related to a spontaneous gap opening (mass formation) for the otherwise gapless Dirac nodes at TQPT point without any symmetry breaking. The gap follows a Gross-Neveu critical behavior. This process of spontaneous mass generation takes place through a condensed matter analog of the CW mechanism in which one-loop quantum fluctuations lead to a mass without symmetry breaking [27]. We expect that our mechanism can be applied to other models for topological phase transitions but also to a wider class of phenomena. A natural example is that of Lifshitz transitions in interacting electronic systems [53], where the continuous deformation of the Fermi surface topology characteristic of noninteracting systems is expected to share the same destiny of the TQPT, i.e., to become discontinuous for large interactions.

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