

Dehybridization transition in multiband systems: A renormalization group approach

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In strongly correlated multiband systems, such as intermetallics, heavy fermions, and Kondo insulators, electron-electron interactions lead to the damping of electronic quasi-particles at finite temperatures. This phenomenon results in an effective dehybridization between the electrons in the large conduction band and those in the narrow, correlated band. The dehybridization effect has been employed to explain the temperature dependence of several physical properties of intermetallics, heavy fermions, and actinide compounds. Recent insights from the theory of non-Hermitian systems provide a novel perspective on this problem. In this work, we adopt this viewpoint through a renormalization group (RG) approach. We derive RG equations for a one-dimensional two-band model to identify the exceptional points of this system and to characterize its topological properties. Our findings reveal that for a topological, non-Hermitian particle-hole-symmetric chain, dehybridization can be interpreted as a topological phase transition, characterized by critical exponents that we calculate. We discuss the nature of the energy scale for dehybridization in the context of the phase diagram of heavy fermions.

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I. INTRODUCTION

At low temperatures, in the Fermi liquid regime of heavy-fermion systems, the f states of the rare-earth or actinide elements hybridize with the electrons in the conduction bands, yielding a coherent state of heavy quasiparticles. In Kondo insulators, this hybridization gives rise to an insulating or semimetallic ground state. However, as the temperature increases, many-body effects and other temperature-dependent scattering mechanisms, such as electron-phonon scattering, cause damping of the electronic quasiparticles and destroy the coherent state. This occurs due to a dehybridization of the f -electronic states and those in the conduction band [1,2] caused by finite lifetime effects. In the incoherent high-temperature regime, the f states behave independently, decoupled from the conduction band.

The notion that electron-electron or electron-phonon scattering in multiband systems may alter the hybridization between these bands was discussed in Refs. [1–3]. These works laid the ground for understanding the nature of the lifetime effects [2] and how at finite temperatures, by making the hybridization less effective, they eventually give rise to a *dehybridization transition*. This approach has been used to understand transport properties of intermetallic systems [3] and f -band metals [4–17].

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Additionally, the emergence of a gap with decreasing temperature observed in angle-resolved photoelectron spectroscopy (ARPES) experiments on the heavy-fermion compound Ce_2RhIn_8 has been interpreted within this scenario [18–20].

Quite generally, disorder or interactions introduce a lifetime for the quasiparticles and a general criterion whether this is still a useful concept compares their energy with their damping. In a region where the damping, or inverse lifetime, of these modes is small compared with their energy, they are still well defined and the idea of a quasiparticle remains valid. This is the reason why long wavelength phonons exist in glasses and spin waves in amorphous ferromagnets [21]. The damping of these excitations due to disorder has a power-law dependence on the momentum larger than that of their energy. Therefore, in the hydrodynamic limit $k \rightarrow 0$, they remain well defined. In electronic systems, interactions give rise to a damping of the quasiparticles that goes to zero faster than their energy, as the Fermi surface is approached [21,22]. The high-frequency counterpart of this behavior occurs in liquids [23], where well-defined transverse phonon excitations appear only at high frequencies or large momenta [24], giving rise to *gapped momentum states* [25].

More recently, new insights in the phenomenon of dehybridization have been possible due to the progress in the study of non-Hermitian systems, exceptional points, and their topological properties [26–34]. In this paper, we discuss dehybridization in strongly interacting multiband systems, from this perspective, using a renormalization group (RG) approach. The effect of interactions that give rise to dissipation is included in a semiphenomenological way [3], validated by microscopic calculations [2,31]. Our results emphasize an intrinsic connection between exceptional points and dehybridization.

A parametrization of the self-energy of the strongly correlated electrons in a heavy fermion, calculated using the local density approximation and dynamical mean-field theory (LDA + DMFT) [35], is used to obtain the temperature-dependent lifetime of the quasiparticles, which satisfies constraints in the low- and high-temperature limits.

The renormalization group is especially useful to study dehybridization in a non-Hermitian particle-hole-symmetric chain. Since this chain has a topologically nontrivial phase in the Hermitian case, it allows for a full investigation of the interplay between topology and non-Hermiticity in multiband materials. The RG approach associates dehybridization with a topological phase transition and yields its critical exponents. The identification of a characteristic length that diverges at the dehybridization transition enables the numerical calculation of the correlation length exponent in finite chains. A crossover exponent associated with the relevance of interactions at the fixed point of the noninteracting Hermitian chain is directly calculated. We find that the edge modes of the noninteracting, particle-hole-symmetric chain are robust up to a critical value of the interactions.

Finally, we discuss the nature of the characteristic temperature for dehybridization to occur and how it relates to the Kondo and coherence temperatures in the phase diagram of heavy fermions.

II. THE ZERO-TEMPERATURE COHERENT GROUND STATE

In a two-band Fermi system, at zero temperature, the dispersion relations of the quasiparticles in the presence of hybridization between these bands are given by [19]

$$E_k^{12} = \frac{1}{2}[(\epsilon_k^s + \epsilon_k^f) \pm \sqrt{(\epsilon_k^s - \epsilon_k^f)^2 + 4|V_k|^2}], \quad (2.1)$$

where ϵ_k^s is the dispersion of the particles in the large conduction band, ϵ_k^f that in the narrow band, and V_k their k -dependent hybridization. Here, we regard these quantities as effective parameters, renormalized by interactions in the narrow band, and in this sense they incorporate the many-body effects [36]. However, in order to fully include the many-body interactions and finite-temperature effects, it is essential to take into account the damping of the quasiparticles that arises from the imaginary part of the self-energy [2]. The real part of the self-energy is included in the effective

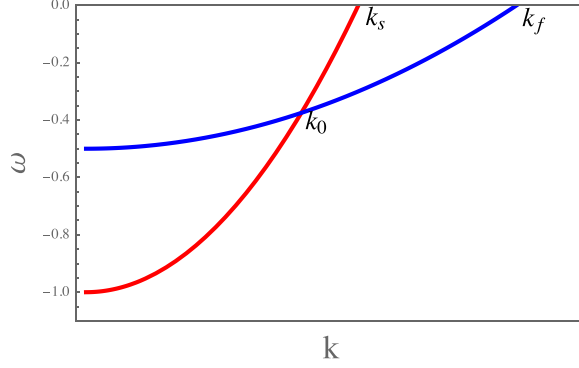


FIG. 1. Dispersion relations of the two-band system in the absence of hybridization. The Fermi level is at $\omega = 0$. We emphasize three special wave vectors, k_0 where the bands cross and the Fermi wave vectors k_s and k_f of the conduction and narrow bands, respectively. We set $\epsilon_f/\mu = 0.5$, $\alpha = 0.2$, and $\mu = 1$.

dispersion relations through a shift in the energy and an effective mass. The damping will be considered below using a semiphenomenological approach.

Initially, we consider a Fermi system where the dispersions of the bands, without hybridization, are given by

$$\epsilon_k^s = k^2 - \mu, \quad \epsilon_k^f = \alpha k^2 - \epsilon_f, \quad (2.2)$$

where $\alpha = m_s/m_f < 1$ is the ratio of the effective masses, μ the chemical potential, and ϵ_f the bottom of the narrow band. The Fermi level is taken at $E_F = 0$. In Fig. 1, we plot these dispersion relations for $\epsilon_f/\mu = 0.5$, $\alpha = 0.2$, and $\mu = 1$. We highlight three special wave vectors in Fig. 1: first, k_0 , where the bands cross, and k_s and k_f , the Fermi wave vectors of the unhybridized s and f bands, respectively.

For the hybridization, we take an antisymmetric $V(\mathbf{k})$, appropriate for the mixing between orbitals of different parities, like s - f , p - d , and d - f [37,38]. In two dimensions for a square lattice, we obtain $V(\mathbf{k}) = iv(\sin k_x + \sin k_y)$. This form of $V(\mathbf{k})$ arises since $V_{ji} = -V_{ij}$ in real space. This form of $V(\mathbf{k})$ breaks inversion symmetry, but since it is purely imaginary, it preserves the time-reversal invariance of the system. We consider a small k expansion, such that, $V(\mathbf{k}) = iv(k_x + k_y)$. For symmetric hybridization, which preserves inversion and time-reversal symmetry, $V(\mathbf{k})$ is real [27]. We focus here on the antisymmetric case since it is associated with nontrivial topological properties, as discussed below.

The dispersion relations of the hybridized bands for $v/\sqrt{\mu} = 0.5$ are shown in Fig. 2. All energies are renormalized by $\mu = 1$. Notice a pair of nodal points, at the line $k_y = -k_x$, that appears due to the antisymmetric hybridization.

III. THE DEHYBRIDIZATION TRANSITION

In order to fully incorporate the many-body and temperature effects in Eq. (2.1), we have to include the damping of the quasiparticles. Here, we consider only that of the heavy quasiparticles. We adopt the usual ansatz [2,3] and introduce an imaginary part in the energy dispersion of these modes, i.e., $\epsilon_k^f \rightarrow \epsilon_k^f - i\tau_f^{-1}$. This approach has been justified in several papers, including Refs. [2] and [31]. We also assume that τ_f is k independent, but has a temperature dependence.

The energy of the quasiparticles is now given by

$$E_k^{12} = \frac{1}{2}[(\epsilon_k^s + \epsilon_k^f - i\tau_f^{-1}) \pm \sqrt{(\epsilon_k^s - \epsilon_k^f + i\tau_f^{-1})^2 + 4|V_k|^2}], \quad (3.1)$$

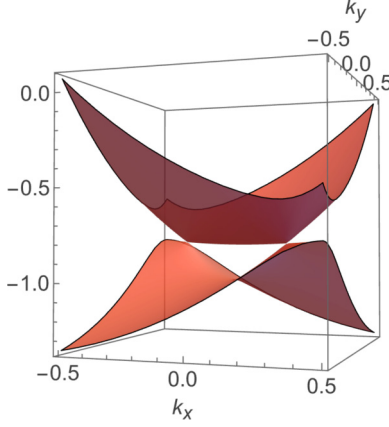


FIG. 2. Dispersion relation of the hybridized bands, for $\alpha = 0.05$, $e_f/\mu = 0.8$, $v/\sqrt{\mu} = 0.7$, with $(\mu = 1)$, showing two nodal points along the line $k_y = -k_x$.

with

$$|V_k|^2 = v^2 k^2 (1 + \sin 2\theta), \quad (3.2)$$

where we used $k_x = k \cos \theta$ and $k_y = k \sin \theta$.

As pointed out in Refs. [2,31], the lifetime of the f modes due to many-body effects is related to the imaginary part of their self-energy,

$$\frac{\hbar}{\tau_f} = -\text{Im}\Sigma. \quad (3.3)$$

This quantity has been calculated for the heavy-fermion system CeIrIn₅ using an LDA + DMFT approach [35]. It is shown in Fig. 3, together with a parametrization of the LDA + DMFT results [18] ($\hbar = 1$) given by

$$\frac{1}{\tau_f} = -\text{Im}\Sigma = \frac{1}{\tau_{0f}} \frac{(T/T^*)^2}{1 + (T/T^*)^2}. \quad (3.4)$$

This function gives an excellent description of the LDA + DMFT results, as can be seen in Fig. 3. It yields $\hbar\tau_{0f}^{-1} = 2$ eV and $T^* = 50$ K, for this compound. When $T \rightarrow 0$, the lifetime of the f quasiparticles becomes very large, $\tau_f \rightarrow \infty$, and the system is in a coherent ground state, with fully hybridized bands. At low temperatures $T \ll T^*$, Eq. (3.4) yields the well-known T^2 temperature dependence of the inverse lifetime of a Fermi liquid. The characteristic temperature T^* marks the onset of the Fermi liquid regime with decreasing temperature. At high temperatures $T \gg T^*$, the inverse lifetime saturates at a finite value that is essentially the bandwidth δe_f of the heavy quasiparticles. This saturation at $\hbar/\tau_{0f} = \delta e_f$ is an important feature and can be viewed as a consequence of the uncertainty principle [3]. From the point of view of transport properties [3,4], at these high temperatures, the resistivity due to s - f scattering saturates since the conduction and f electrons become decoupled and the transport properties are dominated by other scattering mechanisms. The most common is electron-phonon scattering [3] that gives rise to a linear temperature-dependent resistivity.

We now consider the dehybridization process, in more detail, for two cases.

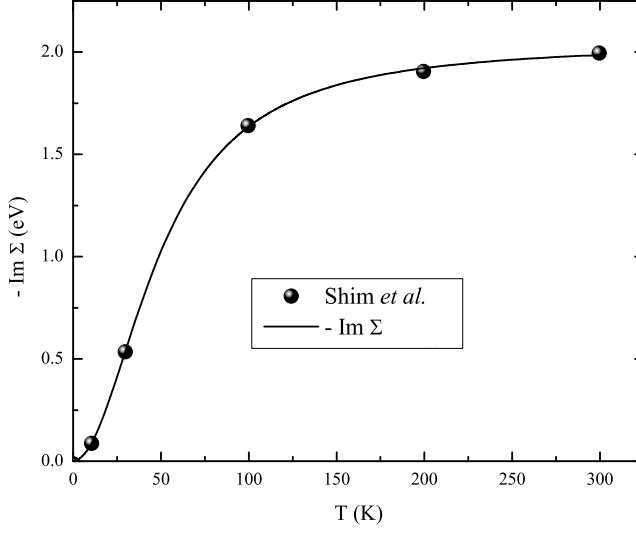


FIG. 3. The imaginary part of the self-energy of CeIrIn₅ as a function of temperature obtained by LDA + DMFT calculations in Ref. [35] (dots). The line is a parametrization of the calculated points given by $-\text{Im}\Sigma = a(T/T^*)^2/(1 + (T/T^*)^2)$, with $a = 2$ eV and $T^* = 50$ K [18].

A. Nodal semimetal

First, consider the case where the bands cross exactly at the Fermi level, i.e., $\epsilon_{k_0}^s = \epsilon_{k_0}^f = 0$ (see Fig. 1). This implies that $k_0^2 = \mu = \epsilon_f/\alpha$ or $\epsilon_f = \alpha\mu$. The energies of the system are given by

$$E_{k_0}^{12} = \frac{1}{2} \left[-i\tau_f^{-1} \pm \sqrt{4|V_{k_0}|^2 - \tau_f^{-2}} \right], \quad (3.5)$$

where $|V_{k_0}|^2 = v^2 k_0^2 (1 + \sin 2\theta)$, with $0 \leq |V_{k_0}|^2 \leq 2v^2 k_0^2$. At $T = 0$, with $\tau_f = \infty$, the coherent ground state has two nodes at the Fermi surface along the line $k_y = -k_x$, as shown in Fig. 2. The ground state is a nodal semimetal [39]. As temperature increases, τ_f^{-1} increases and the Fermi points develop into Fermi arcs, as shown in Fig. 4, where we used the parametrization for τ_f^{-1} given by Eq. (3.4).

As pointed out in Ref. [31], dehybridization in this case occurs gradually, with a different dehybridization temperature for each angle θ in the Fermi surface. As temperature increases, the

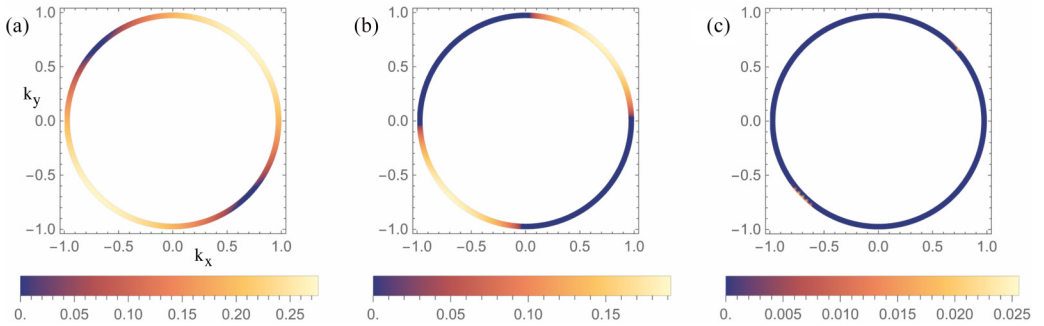


FIG. 4. Gapless regions in the Fermi surface, the Fermi arcs, for different temperatures ($k_0^2 = \mu$). (a) $T = T_D/8$, (b) $T = 3T_D/8$, and (c) $T = T_D$, with T_D given by Eq. (3.6).

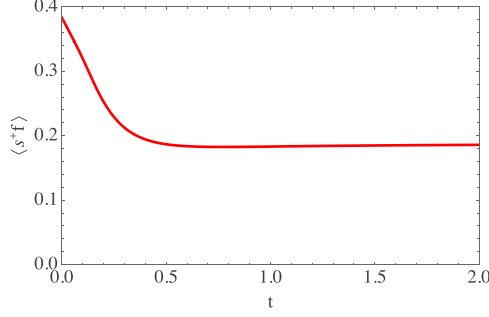


FIG. 5. The correlation function $\langle c_{k_0}^+ f_{k_0} \rangle$ at $k_0 = \sqrt{\mu}$, averaged over the Fermi surface, as a function of scaled temperature ($t = T/T^*$). The parameters are $v/\sqrt{\mu} = 0.5$, $\alpha = 0.1$, and $\epsilon_f/\mu = \alpha = 0.1$, with $\mu = 1$. The bands cross at the Fermi level, $E_F = 0$ and $k = k_0$. The temperature dependence of the relaxation time is given by $\tau^{-1} = \epsilon_f(t^2/(1+t^2))$. The value of the correlation function at $T = 0$ is obtained in Appendix A.

Fermi arcs increase until the whole Fermi surface becomes gapless at a *dehybridization temperature* T_D , when $\sin 2\theta = 1$, as shown in Fig. 4(c). The appearance of the Fermi arcs is connected with the existence of exceptional points where the energies $E^1(k) = E^2(k)$ become degenerate [31]. They are also a direct consequence of the angular or momentum dependence of the hybridization, which determines the points of the Fermi surface where they appear. Thus, although Fig. 4 shows results for a specific form of the hybridization, it illustrates the general feature of nodal points at the zeros of V_k that expand into arcs as dehybridization occurs with increasing temperature. Dehybridization in this case is a gradual process, a crossover phenomenon, with an energy scale that can be characterized by a dehybridization temperature T_D , at which the whole Fermi surface becomes gapless.

For small values of hybridization, i.e., $v/(\alpha\sqrt{\mu}) < 1/(2\sqrt{2})$, this occurs for $\tau_f(T_D)^{-2} = 8v^2k_0^2$. Using $k_0^2 = \mu$, $\tau_{0f}^{-1} = \epsilon_f = \alpha\mu$, we get

$$T_D = T^* \sqrt{\frac{v^*}{1-v^*}}, \quad (3.6)$$

with $v^* = 2\sqrt{2}(v/(\alpha\sqrt{\mu}))$. For $v/(\alpha\sqrt{\mu}) = 0.1$, we find $T_D \approx 1.14T^*$. Notice that dehybridization occurs at exceptional points of the dispersion of the non-Hermitian Hamiltonian [31]. As shown in Eq. (3.6), the crossover temperature T_D is proportional to T^* , a point that we discuss in more detail below.

In general, the correlation function $\langle c_{k_0}^+ f_{k_0} \rangle$ is useful to characterize dehybridization. For the two-band, two-dimensional (2D) electron system, this quantity, averaged over the Fermi surface, is shown as a function of scaled temperature in Fig. 5. This correlation function has been associated with the Kondo effect in lattice systems [40]. Since dehybridization gives rise to a decoupling of the electrons in the narrow and large bands, it implies a breakdown of the Kondo effect. For $t = 1$, a significant decrease of the hybridization correlation function has occurred. Besides, the behavior of the correlation function in Fig. 5 is characteristic of a crossover effect.

B. Metal

Let us consider the case where the ground state is metallic. This can be described by assuming that the line where the dispersion relations cross is not at the Fermi surface, but close to it. We take $\epsilon_{k_0}^s = \epsilon_{k_0}^f = \epsilon = (\alpha\mu - \epsilon_f)/(1 - \alpha)$, at the crossing line given by $\tilde{k}_0^2 = (\mu - \epsilon_f)/(1 - \alpha)$. The energies of the quasiparticles of the hybridized system are given by

$$E_k^{12} = \frac{1}{2} [2\epsilon - i\tau_f^{-1} \pm \sqrt{4|V_{k_0}|^2 - \tau_f^{-2}}], \quad (3.7)$$

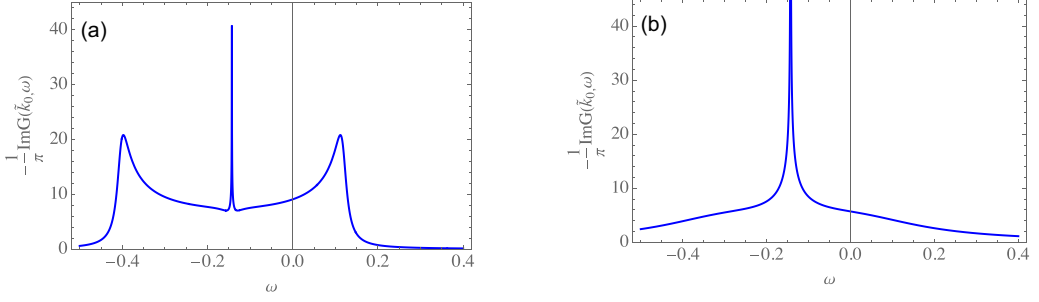


FIG. 6. (a) The average spectral density at $\tilde{k}_0$ for $T = 0.2 T^*$. (b) The average spectral density at $\tilde{k}_0$ for $T = 2 T^*$. The parameters used in both figures are $\alpha = 0.3$, $v/\sqrt{\mu} = 0.2$, and $\epsilon_f/\mu = 0.4$, with $\mu = 1$.

where $|V_{\tilde{k}_0}|^2 = v^2 \tilde{k}_0^2 (1 + \sin 2\theta)$. The frequency-dependent spectral density at the wave vector $\tilde{k}_0$ where the bands cross is given by the imaginary part of the Green's functions [21] (see Appendix A),

$$-\frac{1}{\pi} \text{Im} G(\tilde{k}_0, \omega) = -\frac{1}{\pi} \sum_{i=1,2} \text{Im} \left(\frac{1}{\omega - E_{\tilde{k}_0}^i} \right). \quad (3.8)$$

The average of this quantity over the Fermi surface is shown in Figs. 6(a) and 6(b) for low temperatures, $T = 0.2 T^*$, and high temperatures, $T = 2 T^*$, respectively. The parameters we use are $\mu = 1$, $\alpha = 0.3$, $v = 0.2$, and $\epsilon_f = 0.4$. At low temperatures ($T = 0.2 T^*$), the spectral density at $\tilde{k}_0$ has a peak close to the energy $\epsilon \approx -0.14$ and two peaks separated by repulsion due to hybridization. As temperature increases, dehybridization occurs. The contribution of the f states has a tendency to accumulate close to ϵ , while the conduction states remain spread in energy, as shown in Fig. 6(b) for $T = 2 T^*$. The crossover between the two regimes is set by the characteristic temperature $T_D \sim T^*$, at which significant changes in the spectral weight occur.

When $\alpha \rightarrow 0$, i.e., for a nearly localized f level, $\epsilon = (\alpha\mu - \epsilon_f)/(1 - \alpha) \approx -\epsilon_f$. In this case, the spectral density at $\tilde{k}_0$, which contributes to the density of states, can become very large as ϵ gets closer to the Fermi level, even at $T = 0$. This can be related to a tendency of the system toward a magnetic instability.

IV. THE ONE-DIMENSIONAL PROBLEM

In this section, we study the one-dimensional case of a periodic chain with two orbitals per site and electronic correlations in one of the bands. The antisymmetric hybridization takes into account the different parities of these orbitals. In the noninteracting case, this system realizes probably the simplest example of a topological insulator [38], presenting topological transitions and zero-energy end states [38]. Here, we generalize these chains to include interactions and consequently non-Hermiticity, to study its effect on the topological properties and the dehybridization transition. The Hamiltonian can be written in matrix form as

$$H = \begin{pmatrix} \epsilon_k^s & V_k \\ V_k^* & \epsilon_k^f + i\tau^{-1} \end{pmatrix}, \quad (4.1)$$

where

$$\epsilon_k^s = \epsilon_s - 2t \cos k, \quad (4.2)$$

$$\epsilon_k^f = \epsilon_f - 2\alpha t \cos k, \quad (4.3)$$

$$V_k = -2iV \sin k. \quad (4.4)$$

We apply a rotation using

$$S = \begin{pmatrix} 1 & 1 \\ -1 & 1 \end{pmatrix}$$

to obtain

$$SHS^{-1} = h_0 \mathbb{1} + h_x \sigma^x + h_y \sigma^y, \quad (4.5)$$

with

$$h_0(k) = 1/2(i\tau^{-1} + \epsilon_f + \epsilon_s - 2(1 + \alpha)t \cos k), \quad (4.6)$$

$$h_x(k) = 1/2(i\tau^{-1} + \epsilon_f - \epsilon_s + 2(1 - \alpha)t \cos k), \quad (4.7)$$

$$h_y(k) = 2V \sin k. \quad (4.8)$$

The term proportional to the identity, h_0 , will be neglected in the study of the topological properties. The remaining matrix,

$$\mathcal{H} = h_x \sigma^x + h_y \sigma^y, \quad (4.9)$$

has eigenvalues given by

$$E = \pm \sqrt{h_x^2 + h_y^2}, \quad (4.10)$$

which can be written as

$$E = \pm \sqrt{A(k)B(k)}, \quad (4.11)$$

with

$$A(k) = \epsilon_f - 2t(1 - \alpha) \cos k - i(\tau^{-1} - 4V \sin k), \quad (4.12)$$

$$B(k) = \epsilon_f - 2t(1 - \alpha) \cos k - i(\tau^{-1} + 4V \sin k). \quad (4.13)$$

Notice that $B(k) \neq A^*(k)$. The exceptional points of the non-Hermitian Hamiltonian [26] are given by the zeros of $A(k)$ and $B(k)$,

$$\cos k = \epsilon_f/2t(1 - \alpha), \quad \sin k = \tau^{-1}/4V \quad (4.14)$$

and

$$\cos k = \epsilon_f/2t(1 - \alpha), \quad \sin k = -\tau^{-1}/4V, \quad (4.15)$$

respectively. Squaring the terms in Eqs. (4.14) and (4.15), we find

$$\left(\frac{\epsilon_f}{2t(1 - \alpha)} \right)^2 + \left(\frac{\tau^{-1}}{4V} \right)^2 = 1, \quad (4.16)$$

so that the exceptional points lie on a circle of radius 1. Equation (4.16) separates two phases with different topological properties, characterized by distinct values of the *vorticity* [27,28], as shown in Fig. 7. Vorticity is a topological invariant for non-Hermitian systems defined by $\vartheta = 1/2(\vartheta_A - \vartheta_B)$ [27,28,41,42], where

$$\vartheta_A = \frac{1}{4\pi} \oint dk \frac{\partial}{\partial k} \tan^{-1} \left[\frac{\text{Im}A(k)}{\text{Re}A(k)} \right], \quad (4.17)$$

and a similar definition for ϑ_B , with $B(k)$ in Eq. (4.12), replacing $A(k)$ in Eq. (4.17). The integral is performed over the whole Brillouin zone [28]. It can be calculated numerically and we find that for points inside the circle in Eq. (4.16), i.e., $(\epsilon_f/2t(1 - \alpha))^2 + (\tau^{-1}/4V)^2 < 1$, $\vartheta_A = 1/2$

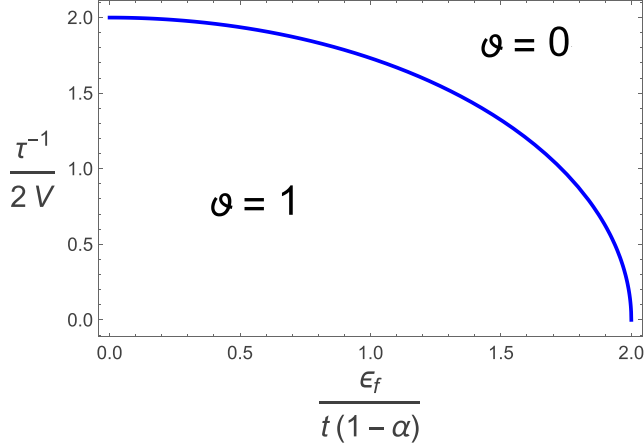


FIG. 7. Phase diagram (first quadrant) of a linear chain with hybridized non-Hermitian bands, with two phases characterized by trivial and nontrivial values of the vorticity ϑ .

and $\vartheta_B = -1/2$; otherwise, both $\vartheta_{A,B} = 0$. Vorticity is quantized as $\mathbb{Z}/2$ [28,41], but in the non-Hermitian two-band chain, pairs of points are either inside or outside the unit circle [Eq. (4.16)], such that $\vartheta_A - \vartheta_B$ is either 1 or 0, in the nontrivial and trivial topological phases, respectively. The phase diagram of the system is shown in Fig. 7. The line $\tau^{-1} = 0$ corresponds to a Hermitian chain with antisymmetric hybridization.

In this Sec. IV, we neglect the temperature dependence of the relaxation time τ^{-1} , which should be viewed as a control parameter driving a phase transition.

A. Particle-hole-symmetric chains

The particle-hole-symmetric case of the two-band problem is particularly interesting. It provides results that allow for a detailed investigation of the interplay between interactions, non-Hermiticity, and topological properties in the dehybridization transition. In the equations of the previous section, this corresponds formally to taking $\alpha = -1$ and $\epsilon_s = -\epsilon_f = \epsilon$. In the Hermitian, noninteracting case, this system has a topological insulating phase, with zero-energy surface states, for $-1 < \epsilon/2t < 1$, and a topological quantum phase transition for a trivial insulator at $|\epsilon/2t| = 1$ [38]. This transition is characterized by the divergence of a characteristic length identified as the penetration depth of the end states in the topological phase. The relevant critical exponents are the correlation length exponent $\nu = 1$ and the dynamic exponent $z = 1$ [38].

The dispersion relations of the bands in the non-Hermitian case are

$$\epsilon_k^s = -\epsilon + 2t \cos k + i\tau^{-1}, \quad \epsilon_k^f = \epsilon - 2t \cos k - i\tau^{-1}. \quad (4.18)$$

The exceptional points satisfy

$$\cos k = \epsilon/2t, \quad \sin k = \pm \tau^{-1}/2V. \quad (4.19)$$

and lie on a circle of radius 1, $(\epsilon/2t)^2 + (\tau^{-1}/2V)^2 = 1$.

B. Renormalization group analysis

We perform a decimation of the chain, eliminating every other site, such that the lattice parameter a is rescaled by a factor $a' = a/b$, with $b = 2$ [21]. For the wave vector, this corresponds to taking $k' = bk$, where k' is the rescaled momentum. Using Eq. (4.19) together with the trigonometric relations, $\sin 2k = 2 \sin k \cos k$ and $\cos 2k = \cos^2 k - \sin^2 k$, we obtain the following recursions

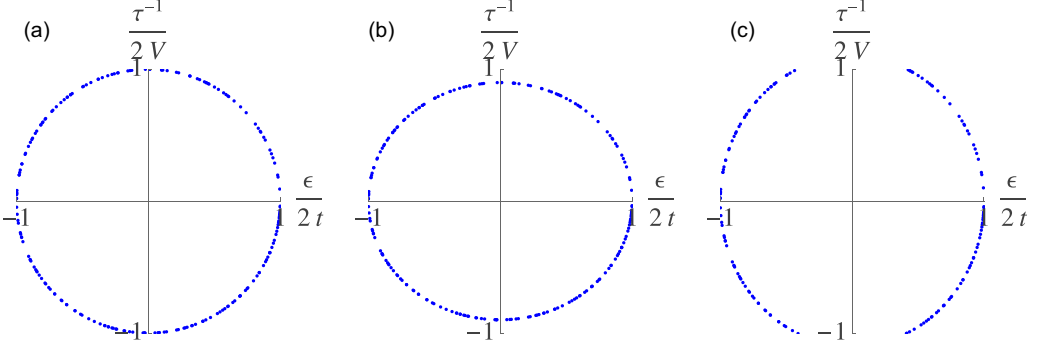


FIG. 8. Maps of the iteration of the recursion relations [Eq. (4.20)]. (a) The initial point is on the circle of radius 1 where the exceptional points are located. These points iterate chaotically along the circle. (b) For an initial point inside the circle with vorticity $\vartheta = 1$, region I in the phase diagram of Fig. 10. (c) For an initial point outside the circle, in region II of zero vorticity, but inside the square circumscribing the circle. In this case, the iterated points that correspond to real values of k do not form a closed curve but an arc. The same behavior is found for an initial point in region III. We used 300 iterations in each case.

relations:

$$\left(\frac{\epsilon}{2t}\right)' = 2\left(\frac{e}{2t}\right)^2 - 1, \quad \left(\frac{\tau^{-1}}{2V}\right)' = 2\left(\frac{\tau^{-1}}{2V}\right)\left(\frac{e}{2t}\right), \quad (4.20)$$

and an additional equation for the ratio $r = \tan k = (\tau^{-1}/2V)/(\epsilon/2t)$,

$$r' = \frac{2r}{1 - r^2}. \quad (4.21)$$

Equation (4.20) is iterated with the initial constraint $(\tau^{-1}/2V)_0^2 + (\epsilon/2t)_0^2 = 1$. Initial points in the circle, under iteration, remain in the circle of radius 1, as shown in Fig. 8(a). These points iterate chaotically along the circle [21]. In fact, the recursion relation for $(\epsilon/2t)$ in Eq. (4.20) is the well-known logistic map [21], which is chaotic. The distribution of iterated points along the circle is not uniform. This can be most easily seen by focusing on the RG equation for the ratio $r = (\tau^{-1}/2V)/(\epsilon/2t)$ in Eq. (4.21). This map has, in the real axis, a single repulsive fixed point at $r = 0$. We start with an initial point in the circle and calculate the image of this point for 10^5 iterations according to Eq. (4.21). Figure 9 shows the number of visits to a given value of the ratio $r = (\tau^{-1}/2V)/(\epsilon/2t)$ on the real axis, for an arbitrary initial point in the circle. Most of the points under renormalization concentrate at small values of the ratio, since the distribution of visits is peaked at $r = 0$. For very large number of iterations, this distribution converges to a Cauchy distribution centered at $r = 0$ with width 1, as evidenced by the fit in Fig. 9. Its normalized form is given by $p(r) = P(r)/(\pi P(0)) = (1/\pi)(1/(1 + r^2))$, where $P(r)$ is the bare, calculated distribution of visits shown in Fig. 9.

More generally, initial points with different values of parameters iterate to different attractors, as shown in Fig. 8. This allows to construct the phase diagram shown in Fig. 10. Points in region I, characterized by a finite vorticity, iterate along closed curves, as shown in Fig. 8(b). Initial points in regions II and III with zero vorticity renormalize along arcs in the case of real momenta, as shown in Fig. 8(c). Finally, initial points in region IV, with $|\epsilon/2t| > 1$, iterate to infinity in both variables. In the Hermitian case, with $(\tau^{-1}/2V) = 0$, initial points with $|\epsilon/2t| > 1$ iterate to $\pm\infty$.

The renormalization group equations [Eq. (4.20)] have two unstable fixed points, respectively, at $(\epsilon/2t) = -1/2$, $(\tau^{-1}/2V) = 0$, and $(\epsilon/2t)_c = 1$, $(\tau^{-1}/2V)_c = 0$. The latter, $(1, 0)$, is the interesting one. It controls the phase transition of the noninteracting Hermitian chain, from the topologically nontrivial to the trivial insulating phase, at $|\epsilon/2t| = 1$. It also describes the crossover from Hermitian

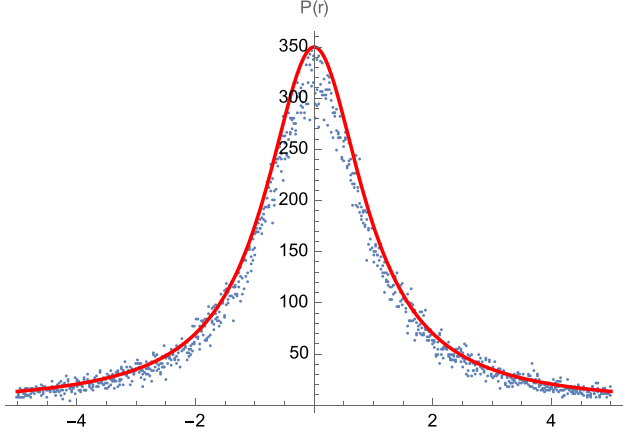


FIG. 9. (a) Number of visits $P(r)$ to a given value of the ratio $r = (\tau^{-1}/2V)/(\epsilon/2t)$ under iteration of the map in Eq. (4.21) with an initial point on the circle of radius 1. The number of visits, which is related to density of states of exceptional points, is represented by a Cauchy distribution of width 1. We used 10^5 iterations. The expression for the normalized distribution is given in the text.

to non-Hermitian behavior as interactions are turned on. The critical exponents associated with this fixed point are obtained from the Jacobian of the RG equations [Eq. (4.20)], calculated at this fixed point. We get

$$J = \begin{pmatrix} 4 & 0 \\ 0 & 2 \end{pmatrix}_{(1,0)}, \quad (4.22)$$

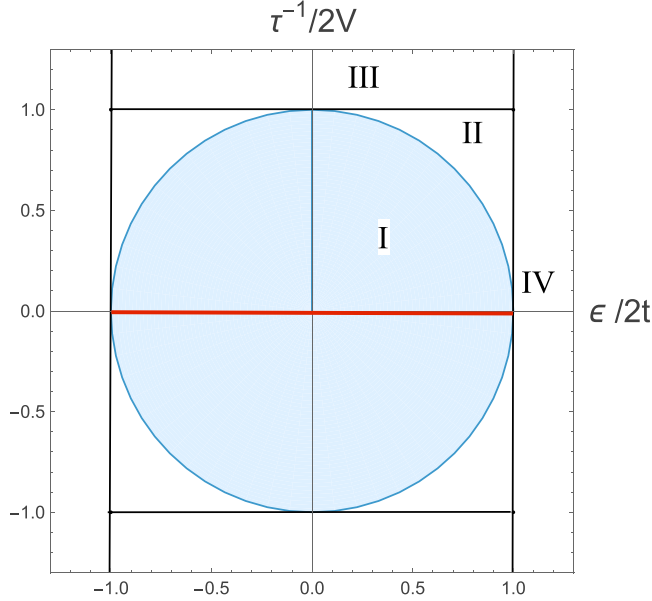


FIG. 10. Phase diagram of the non-Hermitian chain. Points in different regions labeled by I, II, II, and IV iterate to different attractors, as shown in the previous figure (Fig. 8). Points in region IV iterate to $(\epsilon/2t) = \infty$. The red line at $\tau^{-1} = 0$ represents the topologically nontrivial, noninteracting chain.

where we used $\partial(\epsilon/2t)'/\partial(\epsilon/2t) = 4$, $\partial(\tau^{-1}/2V)'/\partial(\tau^{-1}/2V) = 2$, and the cross derivatives are zero at this fixed point $(1, 0)$. Since the Jacobian is diagonal, the renormalization axes are orthogonal and correspond to the axis $(\epsilon/2t)$ and $(\tau^{-1}/2V)$. The correlation length exponent ν is obtained from the relation $b^{2/\nu} = 4$ [21]. Since the scale factor $b = 2$, we get $\nu = 1$ [21,38], in agreement with the analysis of the dispersion relation [43,44] discussed below. The power 2 of the scale factor b arises for the following reason. Near the fixed point at $k = 0$ or $(\epsilon/2t) = 1$, we get using Eq. (4.19), $(\epsilon/2t) = \cos k \approx 1 - k^2/2$ and $k = \sqrt{2(1 - (\epsilon/2t))}$. When we rescale the wave vector by the factor b , we obtain $bk = \sqrt{2b^2(1 - (\epsilon/2t))}$, implying that the renormalization factor for the variable $(\epsilon/2t)$ is b^2 . On the other hand, notice from Eq. (4.19) that the scale factor for $(\tau^{-1}/2V)$ is the same as for the wave vector k near the fixed point at $k = 0$. Since $(\tau^{-1}/2V)' = 2(\tau^{-1}/2V)$, close to $(1, 0)$, the second eigenvalue of the Jacobian equation (4.22) defines the dynamic exponent z through the relation $b^{\nu z} = 2$ that yields $\nu z = 1$. Since $\nu = 1$ at the Hermitian fixed point, we get the dynamic exponent $z = 1$. Then, non-Hermiticity, represented by the quantity $(\tau^{-1}/2V)$, is a relevant perturbation at the Hermitian quantum critical point (QCP) $(1, 0)$. Furthermore, comparing the ratio of renormalization of the variables $(\tau^{-1}/2V)$ and $(\epsilon/2t)$, $b^{\nu z}/b^{2\nu}$, we obtain a crossover exponent $b^{\nu z}/b^{2\nu} = b^{\vartheta}$, where $\vartheta = z/2 = 1/2$. This crossover exponent describes the shape of the critical line, $(\tau^{-1}/2V) = (1 - (\epsilon/2t)^2)^{1/2}$, separating regions with trivial and nontrivial values of the vorticity, close to the fixed point $(1, 0)$. The exponents we find from the RG approach are equal to those obtained from numerical approaches [38] and the analysis of the dispersion relations of the chain, as we discuss below.

The critical exponents for the Hermitian case, $(\tau^{-1}/2V) = 0$, can be directly obtained from the expansion of the dispersion relations [43,44] [Eq. (4.11)] close to the gap closing wave vector $k = 0$. These have the form

$$E^{\pm}(k, \epsilon/2t) = \pm 4t\sqrt{(\epsilon/2t - 1)^2 + (V/2t)^2 k^2}. \quad (4.23)$$

Since the gap for excitations $\Delta = E^+(k = 0, \epsilon/2t) - E^-(k = 0, \epsilon/2t) \propto (\epsilon/2t - 1)$, we get the gap exponent $\nu z = 1$. Also, $E^{\pm}(k, \epsilon/2t = 1)/4t = (V/2t)k$, implying the dynamic critical exponent $z = 1$. These relations yield $\nu = 1$, in agreement with the RG result. In the non-Hermitian case, we consider

$$E^{\pm}(k, \epsilon/2t) = \pm 4t\sqrt{(\epsilon/2t - 1 - i\tau^{-1}/2t)^2 + (V/2t)^2 k^2},$$

such that the real part of the gap at $k = 0$ vanishes with the critical exponent $\nu z = 1$. Its imaginary part $(\tau^{-1}/2t)$ is always finite. On the other hand, for $(\epsilon/2t) = 1$, we get

$$E^{\pm}(k, \epsilon/2t = 1) = \pm 4\sqrt{4V^2 k^2 - \tau^{-2}}, \quad (4.24)$$

which is linear in k for $\tau^{-1} = 0$, as expected from the dynamic exponent $z = 1$. For $\tau^{-1} \neq 0$, the dispersion is real for $k > k_c = (\tau^{-1}/2V)$. Solving the equation above for k , instead of E , we get

$$k = \frac{\tau^{-1}}{2V}\sqrt{1 + (\omega\tau)^2}, \quad (4.25)$$

where $\omega = E^+$. Then, in the limit $\omega\tau \gg 1$, we recover the dynamic exponent $z = 1$. Equation (4.24) is reminiscent of the theory of gapped momentum states [25] where propagating modes appear only above a critical momentum k_c . In this approach, the length scale related to the inverse critical momentum $\ell = 1/k_c$ is the *skin depth* characterizing the penetration of electromagnetic waves in the system [25].

For completeness, we show in Fig. 11, the real part of the dispersion relations of the quasiparticles in different regions of the phase diagram of Fig. 10. Compare the dispersion relation of the bare bands, with no hybridization [Fig. 11(a)], with those in regions II and III [Figs. 11(c) and 11(d)], where the bands are fully dehybridized. In region I, the real part of the dispersion relations shown in Fig. 11(b) is gapped. In this region of the phase diagram, as we show below using numerical

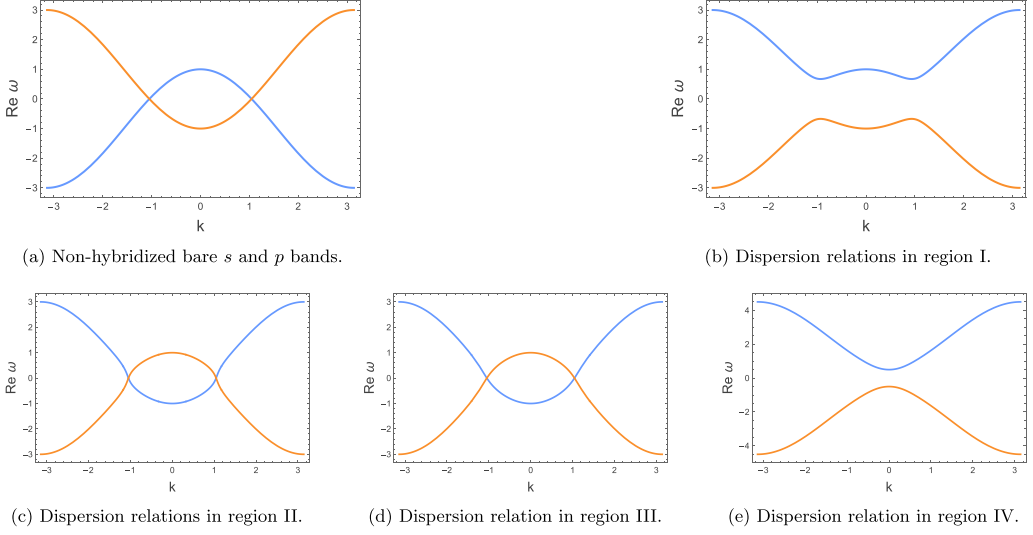


FIG. 11. Real part of the dispersion relations of the quasiparticles in different regions of the phase diagram in Fig. 10. Panel (a) shows the bare bands. (b) In region I, the system is a gapped topological insulator. (c) In region II, the system is *metallic* (see text). (d) In region III, the system is a fully dehybridized *metal* [compare with panel (a)]. In panel (c), the system is a gapped trivial insulator.

simulations in finite chains, the system is a topological insulator with surface states. In region IV, with dispersions shown in Fig. 11(e), the system is a gapped trivial insulator.

C. Finite chains

To fully characterize the effect of interactions on the topological properties of particle-hole-symmetric chains, we study them numerically under open boundary conditions. Due to non-Hermiticity, we distinguish in these calculations between right and left eigenvectors, specially when obtaining the wave functions of the relevant modes.

We start with the Hermitian chain, with $|\epsilon/2t| < 1$, i.e., inside its topological phase, where well-defined edge modes exist [38]. We turn on interactions and investigate how they modify these topological surface modes. We find that everywhere inside the circle of radius 1, corresponding to region I of the phase diagram in Fig. 10, where the vorticity is finite, there are well-defined edge modes in the non-Hermitian chain. This is shown in Fig. 12 for a representative point in region I of the phase diagram. Outside the circle, in regions II and III of the phase diagram with zero vorticity, the zero-energy modes have extended wave functions all along the chain, as shown also in Fig. 12. The same results are obtained when plotting only the left or the right eigenfunctions, or under periodic boundary conditions.

The wave functions of the edge states in region I decay exponentially inside the bulk, as can be seen in Fig. 13, i.e., $\psi(n) \propto e^{-n/\xi}$, where n is the side index and ξ a characteristic length. The penetration depth ξ of the surface modes increases as the phase boundary of region I with finite vorticity, represented by the circle of exceptional points, is approached from within. Furthermore, this characteristic length appears to diverge at this boundary, for a fixed $(\tau^{-1}/2V)$, varying $(\epsilon/2t)$. We find $\xi = |(\epsilon/2t)_c - (\epsilon/2t)|^{-\nu}$, with the exponent $\nu = 1$, as shown in Fig. 13. Then, the circle $(\tau^{-1}/2V)^2 + (\epsilon/2t)^2 = 1$ separating states with zero and finite vorticity is a critical boundary where a topological phase transition occurs accompanied by the divergence of a correlation length with a well-defined critical exponent ν . This critical boundary is also the location of the exceptional points of the system. It separates a topological insulating state with edge states inside the circle, from

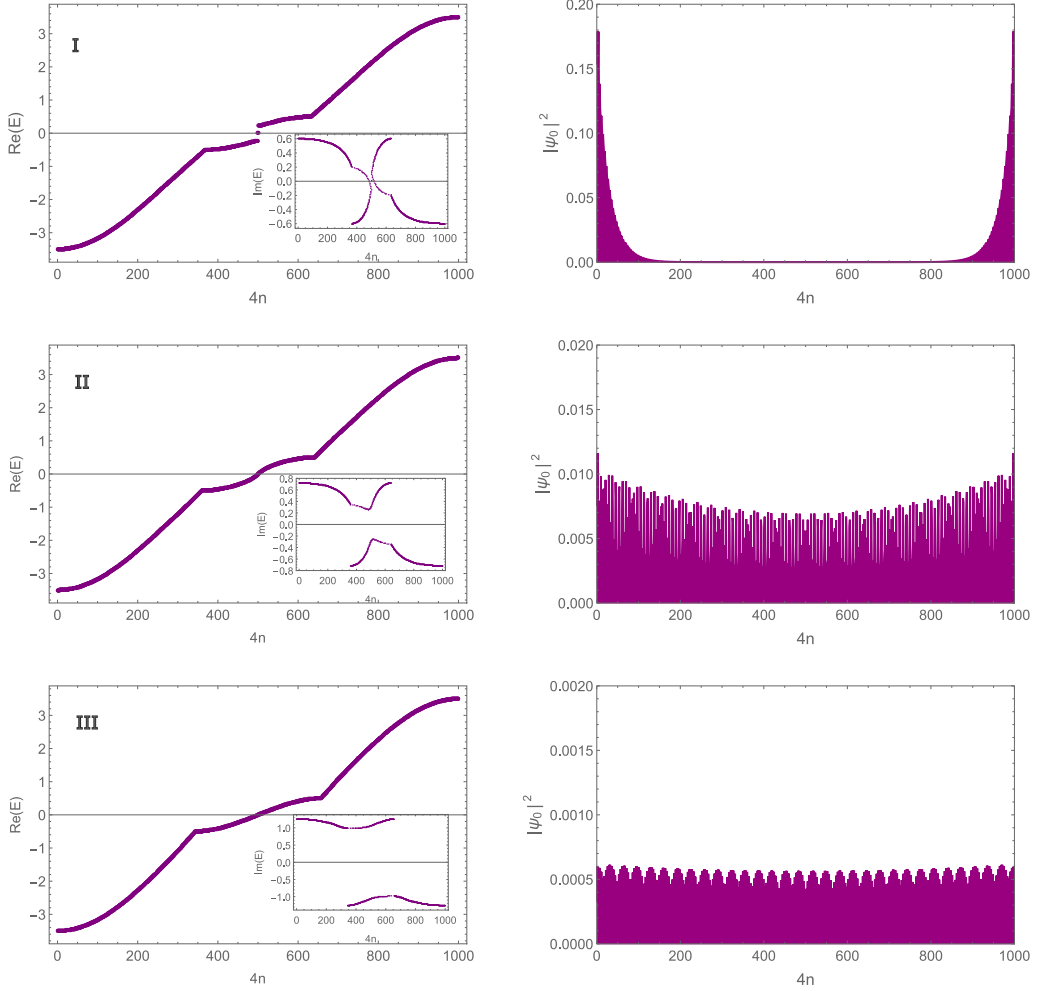


FIG. 12. Real part of eigenvalues and wave functions of the eigenmodes closest to zero energy of finite chains ($N = 250$), for fixed $\epsilon/2t = 0.75$ and $(\tau^{-1}/2V)$ corresponding to regions I, II, and III of the phase diagram shown in Fig. 10 [$(\tau^{-1}/2V) = 0.60, 0.72, 1.25$, respectively]. The insets show the imaginary part of the energy of the modes.

a trivial insulator for $|\epsilon/2t| > 1$, $(\tau^{-1}/2V) = 0$ region IV in Fig. 9, or from a fully dehybridized *metallic state*, corresponding to regions II and III of the phase diagram. By a metallic state, we mean a state of a fermionic system, with low energy excitations at the Fermi level that have extended wave functions and give rise to a linear specific heat at low temperatures, as shown in the Appendix B. These excitations are not well-defined quasiparticles [45], since they are overdamped with an imaginary part in their energy (see Fig. 12, larger than the real part.) In this sense, it would be appropriate to refer to this state as a *strange metal* with a blurred Fermi surface.

There is a remaining transition, from the metallic state in regions II and III to the trivial insulator in region IV, along boundaries represented by the vertical lines $|\epsilon/2t| = 1$. This transition is governed by strongly damped fixed points at $(\tau^{-1}/2V) = \pm\infty$, $(\epsilon/2t) = 1$. On the other hand, the boundary of the topological phase is a strange attractor and is not governed by a single fixed point. We emphasize that the characterization of the phase diagram presented above is consistent whether we use chains with periodic boundary conditions or finite chains with open boundary conditions, as in Figs. 11 and 12.

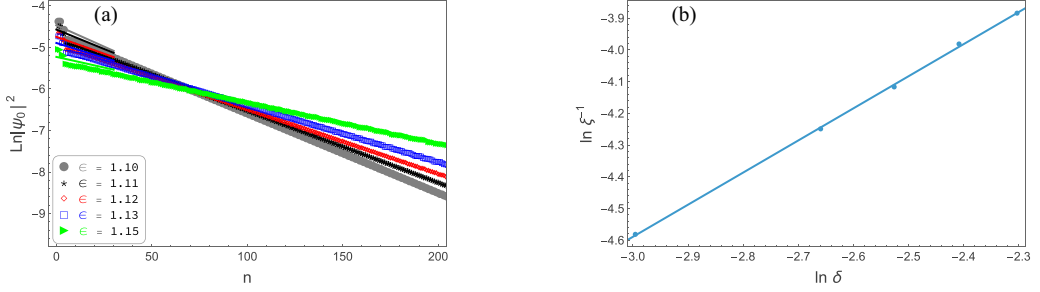


FIG. 13. (a) The wave function of the edge modes in region 1 decay exponentially in the bulk for large n , $\psi \propto e^{-n/\xi}$, where n is the site number and ξ the correlation length. This is shown by the straight lines on the semilog plots of $\psi(n)$, as a function of n . The inclinations of the lines give ξ for different distances to the critical boundary in Fig. 10. Here, we take $(\tau^{-1}/2V) = 0.8$ fixed and approach the boundary at $(\epsilon_c/2t) = 0.6$, varying ϵ . (b) Log-log plot of the inclinations of the lines in panel (a) as functions of $\delta = |\epsilon - \epsilon_c|$. The line through the points is $\text{Ln} \xi^{-1} = -1.56 + 1.01 \text{Ln} \delta$. Since $\xi \propto |\delta|^{-\nu}$, we get $\nu \approx 1$.

Our numerical calculations on finite chains with different boundary conditions do not provide any evidence for skin modes [42,46–51] in the non-Hermitian chain. This is supported by a straightforward analytic calculation using the same approach of Ref. [47]. Consistently, under periodic boundary conditions for the chain, we do not find a closed loop in the complex energy plane [50,51]. Figure 14 shows the iteration of an initial point in region III of the phase diagram of Fig. 10. The images of this point under the RG equations fall mostly in region II of the phase diagram that correspond to real values of momentum. Then, we do not expect any significant difference on the nature of the bulk states in regions II and III of the phase diagram.

The critical exponent ν of the correlation length, obtained when approaching the boundary at finite $(\tau^{-1}/2V)$, has the same numerical value $\nu = 1$ of the Hermitian case [38]. Since $(\tau^{-1}/2V)$ is a relevant perturbation at the Hermitian fixed point, we would expect, in the framework of the RG approach, to find another fixed point in the boundary, eventually in a different universality class that would govern the critical behavior along the unit circle. However, in the case studied here, the boundary is a chaotic attractor, which contains the Hermitian fixed point as the only fully unstable fixed point. Then, the equal values of ν for $(\tau^{-1}/2V) \geq 0$ may be more than a numerical coincidence, implying that the Hermitian fixed point controls the critical behavior all along the boundary.

V. DISCUSSION

In this work, we studied dehybridization in multiband materials from a perspective of the theory of non-Hermitian systems. This gives a insight into this phenomenon that has been invoked in several studies to explain the physical behavior of intermetallic, Kondo insulators [52], and heavy fermions. Concepts such as exceptional points and vorticity play an important role in this analysis, specially in understanding the modifications of the Fermi surface that accompany dehybridization.

Initially, we considered dehybridization as a finite-temperature effect that occurs as the lifetime of the quasiparticles decrease with increasing temperature due to interactions. This lifetime was obtained from numerical calculations using density-functional theory and DMFT methods and presents general features, such as Fermi liquid behavior at low temperatures and a saturation at high temperatures on the order of the bandwidth of the narrow correlated band. The crossover between these regimes sets the scale for dehybridization to occur. A relevant point concerns whether the dehybridization crossover temperature is a new energy scale in the case of strongly correlated multiband systems, or can be identified either with the Kondo or the coherence temperature on the phase diagram of heavy fermions. For the well-studied heavy-fermion compounds of the families

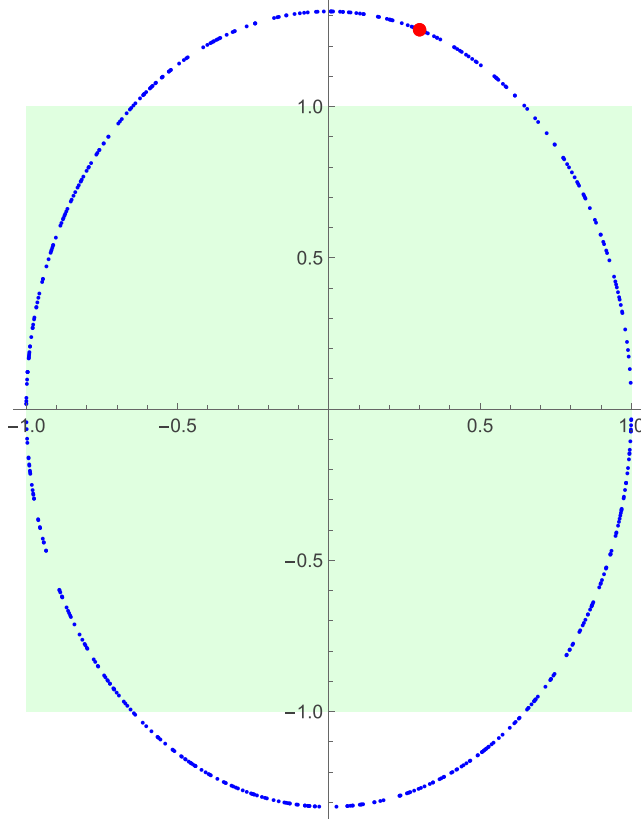


FIG. 14. Images of an initial point (red dot) in region III of the phase diagram under the RG equations, similar to Fig. 8(c). Most of the points fall in region II, the shaded square, of real momenta (see also Fig. 9).

$\text{Ce}_m\text{MIn}_{3m+2}$ ($\text{M} = \text{Co}, \text{Rh}, \text{Ir}; m = 1, 2$) [18,19], the characteristic temperature T^* and the dehybridization temperature T_D are much higher than the critical temperatures of the magnetic or superconducting phase transitions in these materials. In particular, T^* obtained in Sec. III for CeIrIn_5 is much higher than the superconducting critical temperature of this system [35,54]. Then, dehybridization is a relatively high-temperature phenomenon compared to any ordering temperature in these materials. It occurs in an incoherent regime where correlations between different types of quasiparticles are essentially local. In this case, it seems natural to identify the characteristic temperature with the Kondo temperature, i.e., $T^* = T_K$, since the latter marks the onset of the Fermi liquid regime. This is corroborated by the results of Ref. [29] showing that exceptional points associated with the dehybridization crossover appear at the Kondo temperature where magnetic moments are screened. Also the two-band correlation function, $\langle c^+ f \rangle$, varies in a temperature range settled by T^* . Notice that dehybridization implies a breakdown of the Kondo effect, as it decouples the different bands of the material. In the present approach, this occurs as a crossover with increasing temperature and is not associated with a different type of zero-temperature Kondo fixed point.

A different situation occurs if the ground state of the system is a coherent hybridized Fermi liquid close to, but on the noncritical side, the QCP of a magnetic instability [53,55] (see Fig. 15). The relevant scale for the onset of the Fermi liquid regime in this case is the coherence temperature [53], $T_{\text{coh}} = |g|^{\nu z}$, as shown in Fig. 15. Here, g is the distance to the QCP, and ν and z the correlation length and dynamical critical exponents, respectively [53]. In this region of the phase diagram, it is

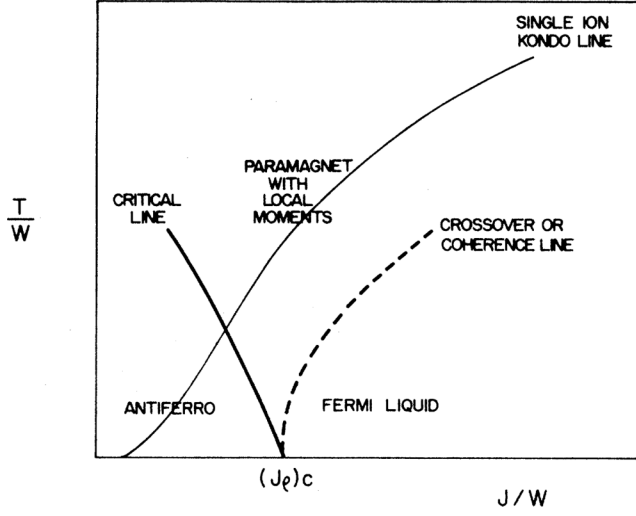


FIG. 15. Phase diagram of heavy-fermion systems as described by the Kondo lattice model [53]. J is the Kondo interaction between local and conduction electrons and $W = 1/\rho$ is the bandwidth of the conduction band. The quantum critical point, which separates the long range ordered magnetic state from the Fermi liquid ground state, is at $J/W = (J\rho)_c$. Reprinted from Ref. [53].

the coherence temperature that gives the scale for dehybridization, since it marks the crossover to the Fermi liquid regime.

An interesting case is that of a material *sitting* just at a magnetic QCP, $(J/W) = (J\rho)_c$ in Fig. 15. There is no characteristic scale in this case since the system is at criticality. The relaxation time according to simple scaling diverges with temperature as $\tau \propto 1/T$ or in more sophisticated approaches with different powers of T that depend on dimensionality d and the dynamic exponent z [56] of the QCP. Dehybridization will occur as temperature increases even at the QCP, as a consequence of critical slowing down.

We have also investigated in this paper dehybridization in a model system consisting of a chain of two orbitals per site with an intersite antisymmetric hybridization, in the presence of interactions. This is special model, since its Hermitian version is a well-characterized topological insulator with a topological quantum phase transition. We introduced a renormalization group to treat the non-Hermitian case from a perspective of critical phenomena [57]. In this model, dehybridization can be associated with a topological phase transition, where a topological index, the vorticity, changes from nontrivial to trivial values, at exceptional points. We obtained the relevant fixed points and calculated the critical exponents both analytically and numerically. Numerical and analytical results rule out the existence of skin modes in this system.

These various aspects of dehybridization show that it is an important concept to understand the physics of multiband strongly correlated electronic systems, with implications in experiments and theory.

VI. CONCLUSIONS

In this paper, we presented results on the phenomenon of dehybridization in multiband correlated electronic systems using the renormalization group. We have focused on the case of antisymmetric hybridization that occurs between electronic orbitals of different parities, like s - p , s - f , p - d , and d - f mixing orbitals. We considered, in particular, non-Hermitian particle-hole-symmetric chains that allow for a detailed study of the interplay between band effects, topology, and interactions in giving rise to dehybridization. We used a renormalization group method to relate the topological behavior

of these chains to a general approach based on the theory of critical phenomena. In this case, dehybridization can be associated with a topological phase transition driven either by interactions or band effects.

Dehybridization is a useful concept to describe the temperature dependence of transport and magnetic properties of multiband metals and Kondo insulators and f -band metals [4–17]. Also, it has provided a framework to understand the evolution with temperature of spectroscopic properties related to modifications of the Fermi surface, as observed in ARPES experiments [18]. The temperature dependence of the lifetime of the quasiparticles, which sets the scale for dehybridization, can be obtained by analytical [2] and numerical [31,35] methods. The new approach presented by the theory of non-Hermitian systems and critical phenomena gives a fresh and deep insight on the nature of this problem.

Generically, the dehybridization transition in many-body systems is a temperature-dependent crossover phenomenon related to gradual modifications of the Fermi surface and rearrangement of spectral weight. We discussed the energy scale relevant for this phenomenon relating it either to the Kondo or the coherence temperature depending on the region of the phase diagram the system is located. We pointed out that dehybridization occurs even for a system at a magnetic QCP as temperature increases. In this work, the breakdown of the Kondo effect is a consequence of dehybridization that decouples the narrow and conduction bands. It is a crossover effect, not related to a new kind of fixed point that occurs as a function of temperature.

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DATA AVAILABILITY

The data that support the findings of this article are openly available [26].

APPENDIX A

In this Appendix, we consider the zero-temperature correlation function $\langle s^+ f \rangle$ for parabolic bands in two dimensions, as in Secs. II and III. We assume the bands cross at the Fermi energy $E_F = 0$ and at the common Fermi wave vector $k_F = k_0$. The Green's function is given by

$$\mathbf{G}^{-1} = (w\mathbf{I} - \mathbf{H}), \quad (\text{A1})$$

with $\mathbf{H}$ as in Eq. (4.1). The dispersion relations are

$$\epsilon_k^s = k^2 - \mu, \quad \epsilon_k^f = \alpha(k^2 - \mu), \quad (\text{A2})$$

and vanish at the Fermi wave vector $k_0 = \sqrt{\mu}$. For the 2D system with parabolic bands, the off-diagonal Green's function is given by

$$G_{sf}(\omega, \mathbf{k}) = \frac{ikv(\cos \theta + \sin \theta)}{-\alpha k^4 + k^2 v^2 \sin(2\theta) + k^2(2\alpha\mu + v^2 + \alpha\omega + w) - (\mu + \omega)(\alpha\mu + \omega)}. \quad (\text{A3})$$

For $|\mathbf{k}| = k = k_0$, we get

$$G_{sf}(\omega, \theta) = -\frac{i\sqrt{\mu}v(\sin(\theta) + \cos(\theta))}{\mu v^2 \sin(2\theta) + \mu v^2 - \omega^2}. \quad (\text{A4})$$

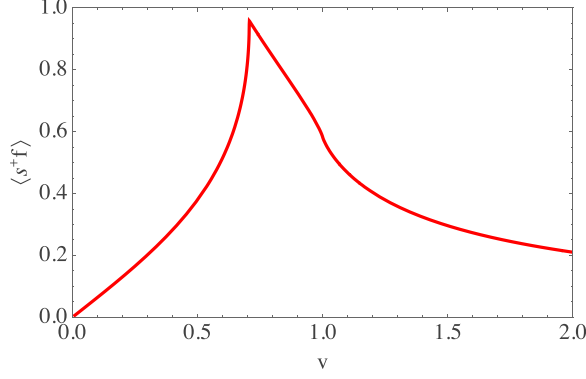


FIG. 16. Zero-temperature correlation function $\langle s^+ f \rangle$ as a function of hybridization $v/\sqrt{\mu}$ for $\mu = 1$.

The correlation function at zero temperature is given by

$$\langle s^+ f \rangle_\theta = -\frac{1}{\pi} \int_{-\mu}^0 G_{sf}(\omega, \theta) d\omega = \frac{1}{\pi} \coth^{-1} \left(\frac{v(\sin(\theta) + \cos(\theta))}{\sqrt{\mu}} \right). \quad (\text{A5})$$

Averaging over the Fermi surface, we obtain the result shown in Fig. 16. Notice that the spectral functions in Fig. 6 are obtained from $G(\mathbf{k}, \omega) = G_{ss}(\mathbf{k}, \omega) + G_{ff}(\mathbf{k}, \omega)$.

APPENDIX B

The specific heat of the non-Hermitian particle-hole-symmetric chain can be calculated using an expression appropriate for the case of complex energies [44]. A slightly modified form is given by [58]

$$C_V/T = \frac{1}{\pi N T^3} \sum_{j=1}^{2N} \int_{-\infty}^{\infty} d\omega \frac{w^2 e^{\omega/T}}{(1 + e^{\omega/T})^2} \frac{|\epsilon_I(j)|}{(\omega - \epsilon_R(j))^2 + \epsilon_I^2(j)}, \quad (\text{B1})$$

where N is the number of sites in the chain and $\epsilon_R(j)$, $\epsilon_I(j)$ the real and imaginary parts of the energy eigenvalues of the chain, respectively. Figure 17 shows the specific heat of the chain in region III of

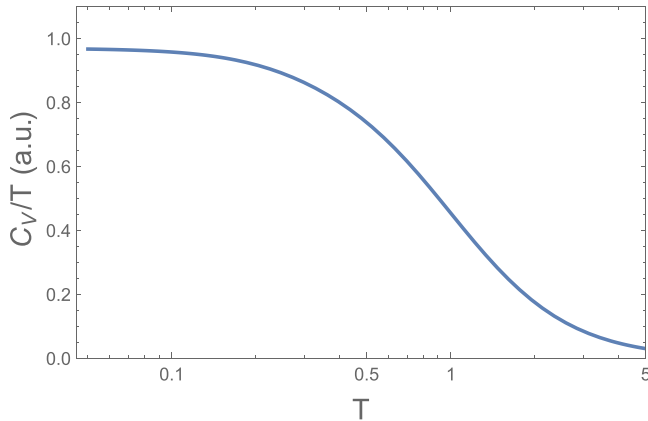


FIG. 17. Specific heat of a non-Hermitian particle-hole-symmetric chain with 200 sites in region III of the phase diagram in Fig. 10. The plot shows the low-temperature linear behavior of the specific heat. Energies are in units of hopping $t = 1$.

the phase diagram of Fig. 10. Specifically, we take $(\epsilon/2t) = 0.5$ and $(\tau^{-1}/2V) = 1.11$. Notice that at the lowest temperatures, the specific heat varies linearly with temperature. The relaxation time is considered independent of temperature in this calculation.

At the Hermitian fixed point, the specific heat can be obtained analytically. The imaginary part of the eigenvalues goes to zero and the Lorentzians become δ functions. The sum is mostly easily performed using periodic boundary conditions in k space. Equation (B1) yields

$$C_V/T = \frac{\pi N}{6v}, \quad (\text{B2})$$

where we used Eq. (4.23), and $v = 2V/t$ is the velocity of the Dirac fermions at $(\epsilon/2t) = 1$. Then, the specific heat is also linear in temperature in this case.

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