

Integrable spin chains with random interactions

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We study a Yang-Baxter integrable quantum spin-1/2 chain with random interactions. The Hamiltonian is local and involves two- and three-spin interactions with random parameters. We show that the energy eigenstates of the model are never localized and in fact exhibit perfect energy and spin transport at both zero and infinite temperatures. By considering the vicinity of a free fermion point in the model we demonstrate that this behavior persists under deformations that break Yang-Baxter integrability but preserve the free fermion nature of the Hamiltonian. In this case the ballistic behavior can be understood as arising from the correlated nature of the disorder in the model. We conjecture that the model belongs to a broad class of models avoiding localization in 1D.

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I. INTRODUCTION

After the fundamental paper by Anderson [1] it was believed for a long time that in one-dimensional random potentials all eigenstates are localized in the thermodynamic limit for arbitrarily weak disorder [2–4]. If Azbel resonances [5,6], which form a set of measure zero, are neglected, the above statement is rigorously speaking valid only for white-noise spatially uncorrelated randomness [7]. Later it was realized that the spatial correlations of the disorder potential can profoundly influence Anderson localization [8–10]. In this case localization can be partially suppressed, at least for weak disorder [11]. In this context a delocalization-localization transition in 1D for *long-range* correlated disorder potentials has been intensively discussed in the literature [12–17]. On the other hand, it was found that models with specific short-range correlated disorder, so-called dimer models, exhibit conducting states [18–21]. In recent years considerable efforts have been made to understand the combined effects of disorder and interactions, which leads to the phenomenon of *many-body localization* (MBL) [22–29]; see Refs. [31–35] for recent reviews. The MBL transition generally occurs at finite energy densities and is characterized by ergodicity breaking, the existence of an extensive number of quasilocal integrals of motion in the localized phase [28–30], and Poissonian level statistics. This is reminiscent of Yang-Baxter integrable many-body systems [36,37], which also feature Poissonian level statistics and extensive numbers of conservation laws. In Yang-Baxter integrable systems the conserved charges are extensive but have (quasi)local densities. An interesting question is then whether there are any connections between Yang-Baxter integrability and MBL. An example of a Yang-Baxter integrable model that is localized is provided by disordered Richardson models [38]. However, this class of models is infinite-ranged whereas studies of MBL have focused on models like the spin-1/2 Heisenberg chain with a random white-noise correlated magnetic field. Other forms of disorder such as random exchange interactions [39,40] have been explored as well [41–43], and there appears to be a widespread

belief that MBL behavior is a rather generic phenomenon in the strong-disorder regime. Non-MBL behavior has been found in a disordered Hubbard chain [44], but this could be related to the presence of non-Abelian symmetries [43,45]. Another little explored issue is what effects correlations in the disorder have on MBL [46–48].

In this work we study a Yang-Baxter integrable model of a Heisenberg-like spin chain with tuneable randomness and Abelian symmetry. We employ a number of standard tools used to probe for (many-body) localized behavior: inverse participation ratios, local quantum quench dynamics, and transport properties in energy eigenstates. All methods point to the same conclusion: the model does not exhibit any traces of localization irrespective of the magnitudes of the interactions and disorder. On the contrary, we find that the model is an ideal conductor for both spin and energy. Moreover, we show in a noninteracting limit that by deforming the model by tuning the correlations between the random interaction parameters (the resulting model is no longer Yang-Baxter integrable) it is possible to induce localization. This suggests that the model we study here is a particular example of a broader class of strongly disordered models in one dimensions that do not exhibit MBL.

II. THE MODEL

The Hamiltonian of our integrable chain contains nearest-neighbor, next-nearest-neighbor, and three-spin interactions with random couplings, cf. Fig. 1, and can be expressed in the form

$$\begin{aligned}
 H = \sum_{j=1}^{L/2} & J_{2j}^{(1)} ([\vec{\sigma}_{2j-1} \cdot \vec{\sigma}_{2j}]_{\Delta_{2j}} + [\vec{\sigma}_{2j} \cdot \vec{\sigma}_{2j+1}]_{\Delta_{2j}}) \\
 & + K_{2j} ([\vec{\sigma}_{2j} \cdot (\vec{\sigma}_{2j-1} \times \vec{\sigma}_{2j+1})]_{\Delta_{2j}^{-1}} + \Delta_{2j}^{-1}) \\
 & + J_{2j}^{(2)} (\vec{\sigma}_{2j-1} \cdot \vec{\sigma}_{2j+1} - 1), \quad (1)
 \end{aligned}$$

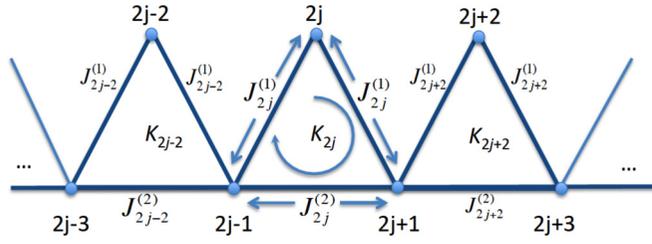


FIG. 1. Schematic representation of the disordered interacting spin chain studied here. Three types of position-dependent interactions with random parameters are present: a nearest-neighbor exchange $J_{2j}^{(1)}$, a next-nearest-neighbor coupling $J_{2j}^{(2)}$, and a three-spin interaction K_{2j} . Explicit expressions for the various terms in the Hamiltonian are given in the text. The ratios $J_{2j}^{(1)}/J_{2j}^{(2)}$ and $J_{2j}^{(1)}/K_{2j}$ are correlated.

where $[\vec{\sigma}_j \cdot \vec{\sigma}_k]_{\Delta} = \sigma_j^x \sigma_k^x + \sigma_j^y \sigma_k^y + \Delta(\sigma_j^z \sigma_k^z - 1)$. The exchange couplings are parametrized as

$$J_{2j}^{(1)} = \frac{\sin^2 \eta \cosh \xi_{2j}}{\sin^2 \eta + \sinh^2 \xi_{2j}}, \quad J_{2j}^{(2)} = \frac{\cos \eta \sinh^2 \xi_{2j}}{\sin^2 \eta + \sinh^2 \xi_{2j}},$$

$$K_{2j} = \frac{\sin \eta \cos \eta \sinh \xi_{2j}}{\sin^2 \eta + \sinh^2 \xi_{2j}}, \quad \Delta_{2j} = \frac{\cos \eta}{\cosh \xi_{2j}}, \quad (2)$$

where ξ_{2j} and $\Delta = \cos(\eta)$ are free parameters of the model. By construction we recover the spin-1/2 Heisenberg XXZ Hamiltonian if we set all inhomogeneities to zero $\xi_2 = \xi_4 = \dots = \xi_L = 0$. In the following we mainly consider the case where ξ_{2k} are independent random variables drawn from a flat distribution

$$P_W(\xi) = \frac{1}{2W} \theta(W - |\xi|). \quad (3)$$

The derivation of the Hamiltonian (1) is summarized in Appendix A. The model (1) is a variant of a class of disordered impurity models previously studied by Klümper and Zvyagin [49], who in particular determined thermodynamic properties [49–53]. Yang-Baxter integrability imposes severe restrictions on the form of the Hamiltonian. This results in all three kinds of interactions involving the same random parameters and can be viewed as short-range correlated disorder in a model of interacting spins.

As a sufficiently strong next-nearest-neighbor exchange can induce dimerization our model can in some sense be considered as an interacting analog of the “dimer models” mentioned above.

Higher conservation laws

As shown in Appendix A the Hamiltonian (1) is related to the transfer matrix $\tau(\mu)$ of an inhomogeneous six-vertex model. This connection is useful for constructing higher conservation laws, which are a characteristic feature of Yang-Baxter integrable models. In the case at hand they can be obtained by taking logarithmic derivatives of the transfer matrix at $\mu = 0$:

$$Q^{(n)} = i^n \left. \frac{d^{n-1}}{d\mu^{n-1}} \right|_{\mu=0} \ln[\tau(\mu)], \quad n = 2, 3, \dots \quad (4)$$

The Hamiltonian is by construction proportional to $Q^{(2)}$:

$$H = -2i \sin \eta Q^{(2)}. \quad (5)$$

Importantly the higher conservation laws are also (ultra)local in the following sense: they can be expressed in the form

$$Q^{(n)} = \sum_{j=1}^L Q_j^{(n)}, \quad (6)$$

where $Q_j^{(n)}$ act nontrivially only on a finite number of neighboring sites. We note that the structure of these conservation laws is very different from that on the “ l bits” in many-body localized systems.

In the following we will make use of the first higher conservation law $Q^{(3)}$. To that end we require an explicit expression in terms of the L operator (A2) and its derivatives. For the operator $Q^{(2)}$ this is readily done:

$$Q^{(2)} = - \sum_{j=1}^{L/2} Q_{2j-1,2j}^{(2,1)} + Q_{2j-1,2j,2j+1}^{(2,2)}, \quad (7)$$

where

$$[Q_{1,2,3}^{(2,2)}]_{\alpha_1 \alpha_2 \alpha_3}^{\beta_1 \beta_2 \beta_3} = [L(-x_2)]_{\alpha_2 d}^{\alpha_1 c} [L'(0)]_{\alpha_3 e}^{\beta_3} [L(x_2)]_{d \beta_2}^{\beta_1},$$

$$[Q_{1,2}^{(2,1)}]_{\alpha_1 \alpha_2}^{\beta_1 \beta_2} = [L'(-x_2)]_{\alpha_2 d}^{\alpha_1 c} [L(x_2)]_{d \beta_2}^{\beta_1}. \quad (8)$$

The conservation law $Q^{(3)}$ can be expressed as a sum of terms that involve spin interactions on two, three, four, and five neighboring sites, respectively:

$$Q^{(3)} = -i \sum_{j=1}^{L/2} [Q_{2j-1,2j}^{(3,1)} + Q_{2j-1,2j,2j+1}^{(3,2)} + Q_{2j-1,2j,2j+1,2j+2}^{(3,3)} + Q_{2j-1,2j,2j+1,2j+2,2j+3}^{(3,4)}], \quad (9)$$

where

$$[Q_{1,2}^{(3,1)}]_{\alpha_1 \alpha_2}^{\beta_1 \beta_2} = [L''(-x_2)]_{\alpha_2 d}^{\alpha_1 c} [L(x_2)]_{d \beta_2}^{\beta_1} - [Q_{1,2}^{(2,1)} Q_{1,2}^{(2,1)}]_{\alpha_1 \alpha_2}^{\beta_1 \beta_2},$$

$$[Q_{1,2,3}^{(3,2)}]_{\alpha_1 \alpha_2 \alpha_3}^{\beta_1 \beta_2 \beta_3} = 2[L'(-x_2)]_{\alpha_2 d}^{\alpha_1 c} [L'(0)]_{\alpha_3 e}^{\beta_3} [L(x_2)]_{d \beta_2}^{\beta_1} - [Q_{1,2}^{(2,1)} Q_{1,2,3}^{(2,2)} + Q_{1,2,3}^{(2,2)} Q_{1,2}^{(2,1)}]_{\alpha_1 \alpha_2 \alpha_3}^{\beta_1 \beta_2 \beta_3},$$

$$Q_{1,2,3,4}^{(3,3)} = Q_{3,4}^{(2,1)} Q_{1,2,3}^{(2,2)} - Q_{1,2,3}^{(2,2)} Q_{3,4}^{(2,1)},$$

$$Q_{1,2,3,4,5}^{(3,4)} = Q_{3,4,5}^{(2,2)} Q_{1,2,3}^{(2,2)} - Q_{1,2,3}^{(2,2)} Q_{3,4,5}^{(2,2)}. \quad (10)$$

The operator $Q^{(3)}$ can be expressed in terms of Pauli matrices using (A2), but this is not particularly useful for our purposes.

III. NONINTERACTING LIMIT

The particular case $\eta = \pi/2$ maps to noninteracting spinless fermions by means of a Jordan-Wigner transformation [54]. The resulting Hamiltonian (1) is block-diagonal $H = P_+ H_+ + P_- H_-$, where $P_{\pm} = \frac{1}{2}[1 \pm (-1)^F]$ are projection operators onto the subspaces with even and odd numbers of fermions, respectively (F is the fermion number operator). We find

$$H_+ = \sum_{j < k=1}^L c_j^\dagger A_{jk} c_k + \text{H.c.}, \quad (11)$$

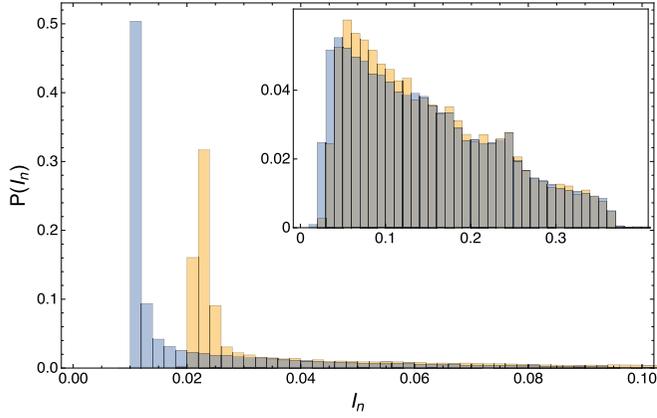


FIG. 2. Histograms of the inverse participation ratios for single-particle energy eigenstates for system sizes $L = 64$ (yellow) and $L = 128$ (blue) averaged over 1000 disorder realizations with probability distribution for inhomogeneities given by the box distribution with $W = 3$. Inset: Same for eigenstates of (11), (13) with $s = 0$, $s' = 0.2$.

where $A_{1,L-1} = 2i \tanh(\xi_L)$, $A_{1,L} = \frac{2}{\cosh(\xi_L)}$, and

$$A_{2j\pm 1,2j} = -\frac{2}{\cosh(\xi_{2j})}, \quad A_{2j-1,2j+1} = 2i \tanh(\xi_{2j}). \quad (12)$$

Fermions tunnel between neighboring sites with amplitudes that are random apart from the constraints $A_{2j-1,2j} = A_{2j,2j+1}$. In addition there is a next-nearest-neighbor hopping on the sublattice of all odd sites. Importantly the corresponding tunneling amplitudes $A_{2j-1,2j+1}$ are not independent random variables, but are related to the amplitudes $A_{2j-1,2j}$. The fermion hopping (12) can therefore be thought of as realizing a particular kind of *correlated disorder*. As we will see, this has important consequences for the physical properties of energy eigenstates. Single-particle energy eigenstates are constructed as $|\Psi_n\rangle = \sum_{j=1}^L \phi_{n,j} c_j^\dagger |0\rangle$, where ϕ_n are the (orthonormal) eigenvectors of the matrix A and $|0\rangle$ is the state without fermions. In order to investigate whether the model (11) is localized we have determined the inverse participation ratio (IPR) of single-particle energy eigenstates $I_n = \sum_{j=1}^L |\phi_{n,j}|^4$. We have considered several probability distributions of the random parameters ξ_{2j} , all of which lead to the same conclusion. We therefore focus on (3). In Fig. 2 we show normalized histograms of I_n averaged over 1000 disorder realizations for $W = 3$ and two different system sizes. We see that the inverse participation ratios are strongly peaked at a value that we find to scale inversely with system size as $1/L$. This indicates that the eigenstates are not localized. At this point the question arises of whether the model (11), (12) is delocalized as a result of fine-tuning, or whether it is representative of a broader class of theories. To investigate this issue we have considered free fermion models of the type (11) with tunneling amplitudes

$$A_{2j\pm 1,2j} = -2|x_{2j}|, \\ A_{2j-1,2j+1} = 2is \operatorname{sgn}(x_{2j}) \sqrt{1 - x_{2j}^2} + s'y_{2j}, \quad (13)$$

where we take x_{2j} and y_{2j} to be independent random variables with probability distribution $P_1(x)$ (3). The tuning parameters $0 \leq s, s' \leq 1$ allow us to interpolate between the ‘‘Yang-

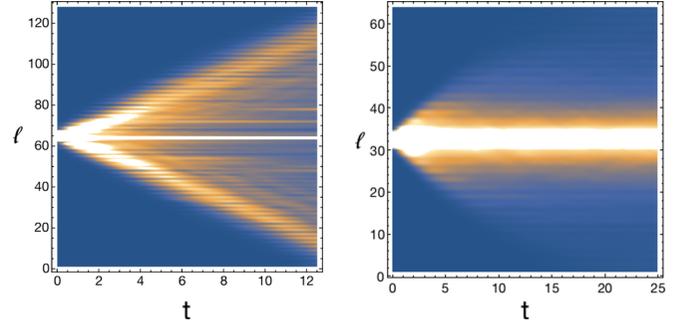


FIG. 3. Left plot: $\langle S_\ell^z(t) \rangle$ averaged over 30 disorder realizations from the box probability distribution for a system of size $L = 128$ and initial thermal state with $\beta = 1$. There is a clear light-cone effect. Right plot: The same for the modified free fermion model (13) with $s = s' = 0$, $L = 64$. Picture is consistent with localization.

Baxter’’ case in which the next-nearest-neighbor tunneling amplitudes $A_{2j-1,2j+1}$ are fixed in terms of the $A_{2j,2j+1}$ and the limit in which they become independent random variables. We have analyzed IPRs for a range of values s and s' . The results suggest that for $s \approx 1$ and small values of s' , i.e., Hamiltonians close to the Yang-Baxter point, eigenstates are delocalized. On the other hand for small values of s and s' , i.e., weak uncorrelated next-nearest-neighbor tunneling, the data are more consistent with localization as shown in the inset of Fig. 2. This suggests that the ‘‘Yang-Baxter’’ model (11), (12) does not correspond to an isolated point in parameter space but is representative of a delocalized region that arises as a result of the correlation between the nearest-neighbor and next-nearest-neighbor tunneling.

Local quantum quench

A second way of investigating localization properties of energy eigenstates is by considering the spreading of correlations after a local quantum quench. We prepare the system in the initial finite energy density state and then overturn two neighboring spins. This choice of initial state allows us to work in the even fermion parity sector of the Hilbert space, $(-1)^{\hat{F}} = 1$. In order to investigate the spreading of correlations we determine the expectation value of the z component of spin at site ℓ . Using Wick’s theorem we obtain compact expressions for $S_\ell^z(t)$ that can be evaluated numerically for systems of hundreds of spins. In Fig. 3 we show results for a representative example, where a system of size $L = 128$ is initially prepared in an energy eigenstate corresponding to inverse temperature $\beta = 1$. We see that the perturbation, which is initially localized at sites $L/2$ and $L/2 + 1$, propagates ballistically through the system, as can be seen from the presence of a ‘‘light cone’’ outside of which our observable remains negligibly small. The velocity characterizing this ballistic propagation depends on the disorder distribution and can be determined exactly in the thermodynamic limit. The spreading of a local perturbation in energy eigenstates of the modified free fermion model (13) can be analyzed in an analogous way. As shown in Fig. 3, for small values of s and s' the perturbation remains localized at sites $L/2$ and $L/2 + 1$ in an extended time window even though a weak light-cone effect occurs at early times. This again indicates that

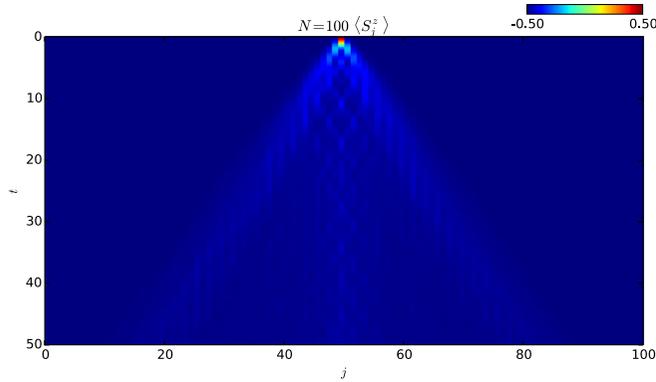


FIG. 4. Spreading of a single spin flip on top of the saturated ferromagnetic state for $L = 100$ and $\eta = 0$ in (1), averaged over 50 disorder realizations with distribution $P_1(\xi)$.

the modified free fermion model is localized at small values of s, s' .

IV. STRONGLY INTERACTING REGIME

Examination of the IPR of (1) away from the free fermion point for small system sizes $L = 10, 12$ is compatible with delocalized behavior of energy eigenstates. We also have studied the spreading of local perturbations in energy eigenstates. (i) We have considered a single spin flip at an odd site on top of the saturated ferromagnetic state. Representative results for the subsequent dynamics on an $L = 100$ site system are shown in Fig. 4. There is a clear light-cone effect that signals ballistic spreading of the perturbation. (ii) We have flipped two neighboring spins in the ground state; cf. Ref. [55] for a discussion of the analogous protocol in the clean system. In this case however our numerics are limited to small systems of up to $L = 16$. We find that there again is a clear light-cone effect; see Fig. 5.

V. BOUNDS ON SPIN AND ENERGY TRANSPORT

We will now demonstrate that the eigenstates of (1) exhibit ballistic energy and spin transport for any anisotropy η and disorder strength W . We employ a combination of two meth-

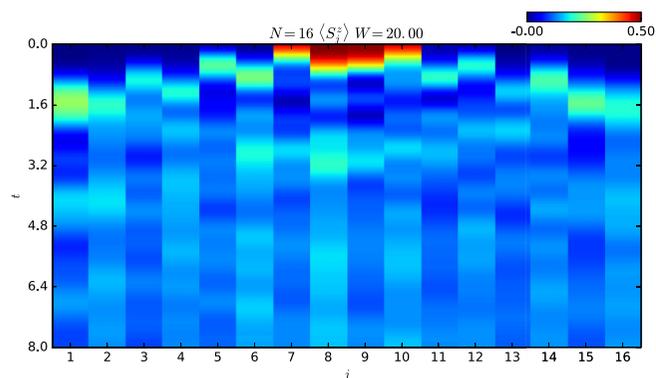


FIG. 5. Spreading of spin flips at two neighboring sites on top of the ground state for $N = 16$ and $\cos(\eta) = 2$ in (1), averaged over 20 disorder realizations with distribution $P_{20}(\xi)$.

ods: the first is based on Mazur's inequality [56] and was previously employed to establish the existence of a finite-temperature Drude weight in the clean case [57], while the second is based on the recently developed hydrodynamic approach to transport in integrable models [58–61]. The starting point of the first approach is the existence of a set of conserved quantities $[H, Q_n] = 0$ that are orthogonal in the sense that $\langle Q_n, Q_m \rangle_\beta = \delta_{n,m} \langle Q_n^2 \rangle_\beta$. Here $\langle \cdot \rangle_\beta$ denotes a thermal expectation value. As the z component of total spin $\sigma^z = \sum_{j=1}^L \sigma_j^z$ is a conserved quantity in our model we employ a magnetic field term to fix the magnetization in our thermal ensemble. Given an operator $A = A^\dagger$ with $\langle A \rangle_\beta = 0$ the following inequality due to Mazur [56] then holds:

$$\lim_{T_0 \rightarrow \infty} \frac{1}{T_0} \int_0^{T_0} dt \langle A(t)A \rangle_\beta \geq \sum_n \frac{\langle A Q_n \rangle_\beta^2}{\langle (Q_n)^2 \rangle_\beta}. \quad (14)$$

A positive bound for the right-hand side of (14) implies that the autocorrelation function of the operator A does not decay to zero at late times. This implies that the Fourier transform has a nonvanishing (generalized) Drude weight:

$$\frac{1}{L} \int_0^\infty dt \cos(\omega t) \langle A(t)A \rangle_\beta = 2\pi D_A \delta(\omega) + \dots \quad (15)$$

When A is the spin current or the energy current operator the nondecay of the autocorrelation functions shows that the system is an ideal conductor of spin/energy. The Hamiltonian (1) has an extensive number of integrals of motion $Q^{(n)}$ (4). The conservation laws relevant to us here have local densities and we focus on the most local of these, $Q^{(3)}$, which involves interactions between spins on at most five neighboring sites; cf. Eq. (9). We furthermore constrain our discussion to infinite temperatures $\beta = 0$. For local operators the corresponding thermal average equals the expectation value in typical energy eigenstates at the associated energy density, which allows us to draw conclusions about the local properties of the eigenstates of (1). In order to use the Mazur inequality (14) we carry out a subtraction $Q_3 = Q^{(3)} - \langle Q^{(3)} \rangle_{\beta=0}$, which ensures that the expectation value of Q_3^2 is extensive, i.e., $\lim_{L \rightarrow \infty} L^{-1} \langle Q_3^2 \rangle_{\beta=0} = a_1 > 0$. The expression for a_1 is very cumbersome so that we do not report it here.

The spin and energy current operators $J^{S,E}$ associated with the Hamiltonian (1) $H = \sum_j H_{2j-1,2j,2j+1}$ are obtained from the continuity equations

$$\begin{aligned} i \sum_{j=-\infty}^{\ell} [\sigma_j^z, H] &= J_\ell^S, \\ i \sum_{j=-\infty}^{\ell} [H_{2j-1,2j,2j+1}, H] &= J_{2\ell}^E. \end{aligned} \quad (16)$$

Evaluating the commutators and then summing over all sites gives

$$\begin{aligned} J^E &= 4i \sin^2(\eta) \sum_j Q_{2j-1,2j,2j+1,2j+2}^{(3,3)} \\ &\quad + Q_{2j-1,2j,2j+1,2j+2,2j+3}^{(3,4)}, \end{aligned} \quad (17)$$

where $Q^{(3,3)}$ and $Q^{(3,4)}$ are given in (10). The spin current operator can be written in the form

$$\begin{aligned} J^S = & 2 \sum_j J_{2j}^{(1)} (T_{2j-1,2j}^{xy} - T_{2j-1,2j}^{yx} + T_{2j,2j+1}^{xy} - T_{2j,2j+1}^{yx}) \\ & + 2J_{2j}^{(2)} (T_{2j-1,2j+1}^{xy} - T_{2j-1,2j+1}^{yx}) \\ & - \frac{2K_{2j}}{\Delta_{2j}} (T_{2j-1,2j,2j+1}^{xzx} + T_{2j-1,2j,2j+1}^{zyz}) \\ & + K_{2j} (T_{2j-1,2j,2j+1}^{zxx} + T_{2j-1,2j,2j+1}^{zyy} + T_{2j-1,2j,2j+1}^{yyz}) \\ & + T_{2j-1,2j,2j+1}^{xxz}, \end{aligned} \quad (18)$$

where we have defined

$$T_{j_1, \dots, j_n}^{\alpha_1 \dots \alpha_n} = \prod_{k=1}^n \sigma_{j_k}^{\alpha_k}. \quad (19)$$

We find that in contrast to the homogeneous case, the energy current is not conserved; i.e., $[H, J^E] \neq 0$.

At infinite temperature and finite magnetization m a tedious but straightforward calculation gives the following result for the overlap of the spin current with the third conserved charge:

$$\langle J^S Q_3 \rangle_{\beta=0} = \frac{\Delta}{1 - \Delta^2} \sum_n \frac{4m(1 - 4m^2) f(\xi_{2n})}{[\cosh(2\xi_{2n}) - (\cos(2\eta))]^3}, \quad (20)$$

where

$$\begin{aligned} f(z) = & \cos(2\eta) \cosh(6z) \\ & - 2[\cos(4\eta) - \cos(2\eta) + 3] \cosh(4z) \\ & + [6 \cos(4\eta) - \cos(6\eta) + 10] \cosh(2z) \\ & - 18 \cos(2\eta) + 2 \cos(4\eta) + 6. \end{aligned} \quad (21)$$

For a very large system we may replace the sum by an integral so that

$$\begin{aligned} \langle J^S Q_3 \rangle_{\beta=0} = & a_S L + o(L), \\ a_S = & \frac{\Delta}{1 - \Delta^2} \int d\xi \frac{4m(1 - 4m^2) f(\xi) P(\xi)}{[\cosh(2\xi) - \cos(2\eta)]^3}. \end{aligned} \quad (22)$$

Here $P(\xi)$ is the probability distribution on the random variables ξ_{2n} . Importantly we have $a_S \neq 0$ unless we fine-tune the probability distribution. This in turn provides a positive bound for the Mazur inequality:

$$\lim_{L \rightarrow \infty} \lim_{T_0 \rightarrow \infty} \frac{1}{T_0 L} \int_0^{T_0} dt \langle J^S(t) J^S \rangle_{\beta=0} \geq \frac{a_S^2}{a_1}. \quad (23)$$

In the case of the energy current for simplicity we consider the zero-magnetization sector $m = 0$. Applying Mazur's inequality we find

$$\begin{aligned} & \lim_{L \rightarrow \infty} \lim_{T_0 \rightarrow \infty} \frac{1}{T_0 L} \int_0^{T_0} dt \langle J^E(t) J^E \rangle_{\beta=0} \\ & \geq \lim_{L \rightarrow \infty} \frac{1}{L} \frac{\langle J^E Q_3 \rangle_{\beta=0}^2}{\langle Q_3 \rangle_{\beta=0}} = \frac{64[2 + 2 \cos(2\eta)]^2}{16 \sin^4(\eta) a_1}. \end{aligned} \quad (24)$$

Interestingly the bound (24) is independent of the inhomogeneities. The generalization to $m \neq 0$ is very tedious but straightforward and provides a nonzero bound as well.

The above calculation proves that at energy densities corresponding to infinite temperature the model (1) exhibits (i) a nonzero Drude weight at any finite magnetization; (ii) ballistic energy transport.

VI. SPIN AND ENERGY TRANSPORT FROM GENERALIZED HYDRODYNAMICS

Generalized Drude weights (15) can be analyzed in full by means of the approach introduced in Ref. [58]. The starting point is the existence of a basis of local charges \hat{Q}_i and associated currents J_i . Using these charges a *generalized Gibbs ensemble* is defined by the density matrix $\rho_{\text{GGE}} \sim \exp(-\sum_n \mu_n \hat{Q}_n)$, where μ_i are ‘‘chemical potentials.’’ The generalized Drude weights D_A are then obtained from appropriate expectation values in this ensemble and are determined by using the thermodynamic Bethe ansatz (TBA) method [62]. According to Ref. [58], in integrable models D_A can be expressed as

$$D_A = \sum_n \int d\lambda \frac{\eta_n(\lambda)}{\rho_n^{\text{tot}}(\lambda)} \left(\frac{\epsilon'_n(\lambda) q_A^{\text{eff}}(\lambda)}{2\pi [1 + \eta_n(\lambda)]} \right)^2, \quad (25)$$

where $\eta_n(\lambda) = \bar{\rho}_n(\lambda)/\rho_n(\lambda)$ is the ratio of hole and particle densities, $\rho_n^{\text{tot}}(\lambda, \{\xi_{2j}\}) = \rho_n + \bar{\rho}_n$, $\epsilon_n(\lambda)$ are the energies of n -string excitations over the state of thermal equilibrium [63], and $q_A^{\text{eff}} = \partial_{\mu_A} \ln \eta_n$ are effective transport charges. The implementation of this approach in our ‘‘inhomogeneous’’ case reveals (see Appendix B for more details) that the disorder merely renormalizes the Drude weight through the disorder dependence of the velocity of the elementary excitations over the equilibrium state under consideration, which enters (15) via the factor $1/\rho_n^{\text{tot}}(\lambda)$. It follows then that the disorder average can be exchanged with the integration and summation in (25). The disorder-averaged Drude weight is then given by

$$\bar{D}_A = \sum_n \int d\lambda \overline{[\rho_n^{\text{tot}}(\lambda)]^{-1}} \eta_n(\lambda) \left[\frac{\epsilon'_n(\lambda) q_A^{\text{eff}}(\lambda)}{2\pi [1 + \eta_n(\lambda)]} \right]^2, \quad (26)$$

where $\overline{[\rho_n^{\text{tot}}(\lambda)]^{-1}} = \int P(\{\xi\}) \frac{1}{\rho_n^{\text{tot}}(\lambda, \{\xi\})}$ denotes the disorder average with probability distribution function $P(\xi)$. As the total density $\rho_n^{\text{tot}}(\lambda)$ is a positive quantity this average is nonzero for generic $P(\xi)$. Therefore, the Drude weight is only renormalized due to the disorder dependence of string particle and hole densities. We note that in contrast to the Mazur bound calculation the TBA approach takes into account the full set of conserved quantities. These observations can be universally extended to any integrable model with disorder of the type described here.

VII. CONCLUSIONS

In this paper we studied a Yang-Baxter integrable interacting spin system with controllable short-range correlated disorder. Using a combination of diagnostics we have demonstrated the absence of many-body localization. We find that the model is in fact an ideal conductor for both energy and magnetization. For particular parameter values the model can be mapped to noninteracting fermions and we have established the absence of Anderson localization in this case. In contrast, a sufficiently strong deformation of the free-fermion Hamiltonian away from

the Yang-Baxter point shows signatures of localization. We expect that in the interacting case small perturbations away from the Yang-Baxter point will lead to diffusive behavior, while sufficiently strong deformations will be required to induce an MBL transition.

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APPENDIX A: INHOMOGENEOUS XXZ CHAIN

The quantum inverse scattering method (QISM) [36] provides a simple way of introducing “impurities” into Yang-Baxter integrable models. This has been used in the literature to construct a variety of models with impurities embedded in both noninteracting and correlated hosts [64–73], as well as models with “disorder” [49–53]. Here we focus on the simplest case, which is related to the spin-1/2 Heisenberg XXZ chain. The basic ingredients in the QISM are the R matrix $R(\mu) \in \text{End}(V_A \otimes V_A)$ and the L operator $L(\mu) \in \text{End}(V_A \otimes V_Q)$, where V_A and V_Q are finite-dimensional “auxiliary” and “quantum” vector spaces. In the cases we are interested in the Yang-Baxter relations read

$$R(\lambda - \mu)[L(\lambda) \otimes L(\mu)] = [L(\mu) \otimes L(\lambda)]R(\lambda - \mu). \quad (\text{A1})$$

In the case of the spin-1/2 XXZ chain we have [36]

$$\begin{aligned} [(L(\lambda))_{\alpha\beta}^{ab}] &= \frac{1 + \tau_{ab}^z \sigma_{\alpha\beta}^z}{2} + b(\lambda) \frac{1 - \tau_{ab}^z \sigma_{\alpha\beta}^z}{2} \\ &\quad + c(\lambda)(\tau_{ab}^- \sigma_{\alpha\beta}^+ + \tau_{ab}^+ \sigma_{\alpha\beta}^-), \\ b(\lambda) &= \frac{\sinh(\lambda)}{\sinh(\lambda + i\eta)}, \quad c(\lambda) = \frac{i \sin(\eta)}{\sinh(\lambda + i\eta)}, \end{aligned} \quad (\text{A2})$$

where η is a free parameter and τ^α , σ^α are Pauli matrices acting on the auxiliary and quantum spaces respectively. The QISM provides a commuting family of transfer matrices $[\tau(\mu), \tau(\lambda)] = 0$ of the form

$$\tau(\mu)_{\alpha_1, \dots, \alpha_L}^{\beta_1, \dots, \beta_L} = \prod_{j=1}^L [L(\mu - \xi_j)]_{\alpha_j \beta_j}^{c_j c_{j+1}}, \quad (\text{A3})$$

where the free parameters ξ_j are known as “inhomogeneities” and where we have defined $c_{L+1} = c_1$. In order to obtain a local Hamiltonian we now set

$$\xi_{2j+1} = 0, \quad (\text{A4})$$

and then take the logarithmic derivative of the transfer matrix at $\mu = 0$

$$H = 2i \sin \eta \left. \frac{d}{d\mu} \right|_{\mu=0} \ln[\tau(\mu)]. \quad (\text{A5})$$

The explicit expression for the resulting Hamiltonian is given by (1).

Spectral properties

The Hamiltonian (1) is readily diagonalized by algebraic Bethe ansatz [36]. The energy eigenvalues are given by

$$E = - \sum_{j=1}^N \frac{4 \sin^2(\eta)}{\cosh(2\lambda_j) - \cos(\eta)}, \quad (\text{A6})$$

where the rapidities $\lambda_1, \dots, \lambda_N$ are solutions of the Bethe ansatz equations

$$\begin{aligned} &\left(\frac{\sinh(\lambda_j + i\eta/2)}{\sinh(\lambda_j - i\eta/2)} \right)^{\frac{L}{2}} \prod_{k=1}^{L/2} \frac{\sinh(\lambda_j - \xi_{2k} + i\eta/2)}{\sinh(\lambda_j - \xi_{2k} - i\eta/2)} \\ &= \prod_{k \neq j} \frac{\sinh(\lambda_j - \lambda_k + i\eta)}{\sinh(\lambda_j - \lambda_k - i\eta)}, \quad j = 1, \dots, N. \end{aligned} \quad (\text{A7})$$

Equations (A6) and (A7) establish a peculiar property of the model (1): the spectrum is invariant under arbitrary permutations of the inhomogeneities $\{\xi_2, \xi_4, \dots, \xi_L\}$, i.e.,

$$\text{spec } H[\{\xi_2, \dots, \xi_L\}] = \text{spec } H[\{\xi_{P(2)}, \dots, \xi_{P(L)}\}] \quad (\text{A8})$$

for any permutation P of the integers $2, 4, \dots, L$. This property is not apparent from the explicit expression (1) and Hamiltonians corresponding to different permutations of the inhomogeneities generally do not commute.

a. Free fermion point

The Hamiltonian (1) has a free fermion point at $\eta = \frac{\pi}{2}$. The corresponding Hamiltonian is

$$\begin{aligned} H &= \sum_{j=1}^{L/2} \frac{1}{\cosh(\xi_{2j})} \sum_{\alpha=x,y} [\sigma_{2j-1}^\alpha \sigma_{2j}^\alpha + \sigma_{2j}^\alpha \sigma_{2j+1}^\alpha] \\ &\quad - \sum_{j=1}^{L/2} \tanh(\xi_{2j}) [\sigma_{2j-1}^y \sigma_{2j}^z \sigma_{2j+1}^x - \sigma_{2j-1}^x \sigma_{2j}^z \sigma_{2j+1}^y]. \end{aligned} \quad (\text{A9})$$

By applying the Jordan-Wigner transformation one can bring Eq. (A9) into the form of Eq. (11).

b. Isotropic (XXX) limit

The SU(2) invariant versions of the Hamiltonian and the Bethe ansatz equations are obtained by redefining

$$\lambda_j = \frac{\eta}{2} \Lambda_j, \quad \xi_{2j} = \frac{\eta}{2} \gamma_j, \quad (\text{A10})$$

and then taking the limit $\eta \rightarrow 0$. This gives a Hamiltonian of the form

$$\begin{aligned} H &= \sum_j \frac{4}{\gamma_j^2 + 4} [\vec{\sigma}_{2j-1} \cdot \vec{\sigma}_{2j} + \vec{\sigma}_{2j} \cdot \vec{\sigma}_{2j+1} - 2] \\ &\quad - \sum_j \frac{2\gamma_j}{\gamma_j^2 + 4} \vec{\sigma}_{2j} \cdot (\vec{\sigma}_{2j-1} \times \vec{\sigma}_{2j+1}) \\ &\quad + \sum_j \frac{\gamma_j^2}{\gamma_j^2 + 4} (\vec{\sigma}_{2j-1} \cdot \vec{\sigma}_{2j+1} - 1). \end{aligned} \quad (\text{A11})$$

The Bethe ansatz equations become

$$\left(\frac{\Lambda_j - i}{\Lambda_j + i}\right)^{\frac{M}{2}} \prod_{k=1}^{M/2} \frac{\Lambda_j - \gamma_k - i}{\Lambda_j - \gamma_k + i} = \prod_{k \neq j} \frac{\Lambda_j - \Lambda_k - 2i}{\Lambda_j - \Lambda_k + 2i}. \quad (\text{A12})$$

The energy corresponding to a solution of (A12) is

$$E = - \sum_j \frac{8}{\Lambda_j^2 + 1}. \quad (\text{A13})$$

APPENDIX B: DRUDE WEIGHTS FROM THE TBA CALCULATIONS

Let us consider first the case of $|\Delta| > 1$. While any eigenstate in a finite system of size L is assigned a unique set of rapidities $\{\lambda_j\}_{k=1}^N$ taken from solutions of Bethe equations (A7), in the thermodynamic limit (defined as $L \rightarrow \infty$, $N \rightarrow \infty$ with N/L finite), the solutions to Bethe equations organize into regular patterns which indicate the presence of well-defined particle excitations. These correspond to magnons and their bound states, so-called Bethe strings [62]. A general string solution reads $\{\lambda_\alpha^{k,m}\} = \{\lambda_\alpha^k + (k+1-2m)\frac{i\eta}{2}\}$, where $m = 1, 2, \dots, k$; α numerates different k strings; and m runs over internal rapidities. Scattering of different magnonic particles is characterized by the amplitudes

$$\begin{aligned} S_j &= \frac{\sin(\lambda - j\frac{i\eta}{2})}{\sin(\lambda + j\frac{i\eta}{2})}, \quad S_{jk} = \prod_{m=-\frac{k-1}{2}}^{\frac{k-1}{2}} \prod_{n=-\frac{j-1}{2}}^{\frac{j-1}{2}} S_{2m+2n+2} \\ &= S_{|j-k|} S_{j+k} \prod_{m=1}^{\min(j,k)-1} S_{|j-k|+2m}^2, \end{aligned} \quad (\text{B1})$$

with the convention that $S_0 \equiv 1$. In the thermodynamic limit particle rapidities become densely distributed along the real axis in the rapidity plane. This permits us to introduce distributions $\rho_k(\lambda)$ of k -string particles, along with the dual hole distributions $\bar{\rho}_k(\lambda)$ (holes are solutions to Bethe ansatz equations which differ from Bethe roots λ_k). The discrete Bethe equations (A7) get replaced by the integral Bethe-Yang equations. Assuming validity of the string solution in the presence of M inhomogeneities ($M/N \leq 1/2$), we can write these integral equations for the densities of string particles and holes in the thermodynamic limit of the inhomogeneous case. The Bethe-Yang equations for particles $\rho_n(\lambda)$ and holes $\bar{\rho}_n(\lambda)$ are given by

$$\begin{aligned} \frac{1}{N} \left(\sum_{j=1}^M a_n(\lambda + \xi_j) + (N-M)a_n(\lambda) \right) \\ = \bar{\rho}_n(\lambda) + A_{nm} \star \rho_m(\lambda). \end{aligned} \quad (\text{B2})$$

Here, the explicit form of the functions $a_n(\lambda)$ and $A_{nm}(\lambda) = \delta_{nm}\delta(\lambda) + a_{nm}$, which depend on the anisotropy parameter Δ , can be obtained from the following relations:

$$\begin{aligned} a_n(\lambda) &= \frac{1}{2\pi i} \partial_\lambda \ln S_n(\lambda), \\ a_{nm}(\lambda) &= \frac{1}{2\pi i} \partial_\lambda \ln S_{nm}(\lambda), \end{aligned} \quad (\text{B3})$$

where indexes n, m label corresponding stringy content. The \star operation refers to the convolution with the kernel A_{nm} ,

$$A_{mn} \star \rho_m(x) \equiv \sum_m \int_{-Q}^Q dy A_{mn}(x-y) \rho_m(y), \quad (\text{B4})$$

where the integration and summation limits depend on the value of anisotropy parameter. Explicitly, for $\Delta > 1$ we have

$$a_n(\lambda) = \frac{1}{2\pi} \frac{\eta \sinh(n\eta)}{\cosh(n\eta) - \cos(\eta\lambda)}. \quad (\text{B5})$$

For the isotropic (XXX) situation, when $\eta \rightarrow 0$, the driving function and the kernel are given by

$$a_n(\lambda) = \frac{1}{\pi} \frac{n}{(n^2) + \lambda^2}, \quad (\text{B6})$$

$$\begin{aligned} A_{nm}(\lambda) &= \delta(\lambda)\delta_{nm} + (1 - \delta_{nm})a_{|n-m|}(\lambda) \\ &\quad + 2a_{|n-m|+2}(\lambda) + \dots + 2a_{n+m-2}(\lambda) + a_{n+m}(\lambda), \end{aligned} \quad (\text{B7})$$

while in this case $Q = \infty$ and sum runs to infinity as well.

Classification of the particle content in the gapless regime $|\Delta| < 1$ is more involved; details can be found in [62,74]. Here, in addition to the magnon-type label k , an extra parity label $v \in \pm$ is required. Importantly, integers k now no longer coincide with the length of a string, i.e., a number of magnons forming a bound state. Instead, the k th particle consists of n_k Bethe roots and carries parity v_k (see [74] for further details). Setting $\Delta = \cos(\gamma)$, where $\gamma/\pi = m/l$ (with m, l co-prime integers) is a root of unity, the number of distinct particles in the spectrum is finite. Changing the parametrization $\lambda \rightarrow i\lambda$, $\eta \rightarrow i\gamma$ and incorporating the additional parity label, the elementary scattering amplitudes and kernels read

$$S_k(\lambda) \rightarrow S_{(n_j, v_j)} = \frac{\sinh\left[\lambda - n_j\frac{i\gamma}{2} + (1 - v_j)\frac{i\pi}{4}\right]}{\sinh\left[u + n_j\frac{i\gamma}{2} + (1 - v_j)\frac{i\pi}{4}\right]}, \quad (\text{B8})$$

and the whole set of scattering kernels is obtained, as in the case of $\Delta > 1$, from Eqs. (B1) and (B3). The Bethe-Yang equations become modified,

$$a_j(x) = \text{sgn}(q_j)(\rho_j + \bar{\rho}_j) + a_{jk} \star \rho_k, \quad (\text{B9})$$

where the summation in the convolution expression runs from 1 to m_l defined as $m_0 = 0$, $m_i = \sum_{k=1}^i v_k$, and numbers $v_1, \dots, v_{l-1} \geq 1$, $v_l \geq 2$ participate in the continuum fraction expression for γ/π , e.g., $\gamma/\pi = 1/\{v_1 + 1/v_2 + 1/(v_3 + \dots)\}$. Numbers q_j are defined recursively as [62], $q_0 = \pi/\gamma$, and

$$\begin{aligned} q_j &= \frac{1}{2}[(1 - \delta_{m_i, j})q_{j-1} + q_{j+1}], \quad m_i \leq j \leq m_{i+1} - 2, \\ q_j &= (1 - 2\delta_{m_{i-1}, j})q_{j-1} + q_{j+1}, \quad j = m_i - 1, i < l. \end{aligned}$$

Explicitly, the kernels $a_j(\lambda)$ are given by

$$a_j(\lambda) = \frac{1}{2\pi} \frac{\gamma \sin(\gamma q_j)}{\cosh(\gamma\lambda) + \cos(\gamma q_j)}. \quad (\text{B10})$$

The most important thing to notice here is that the left-hand side (driving terms) of the Bethe-Yang equations depends on the inhomogeneities while the right-hand side (convolution kernel) *does not* depend on inhomogeneities. This can also

be checked by explicit re-derivation of steps leading to these equations (B2).

The second set of equations is derived using the variation of the free energy (per particle) $f = e - Ts$ with respect to ρ_n and $\bar{\rho}_n$. Here

$$e = 2\pi \sum_n a_n(\lambda) \rho_n(\lambda) \quad (\text{B11})$$

is the energy density and the entropy density is

$$s = \sum_{n=1}^{\infty} \int_{-\infty}^{\infty} d\lambda [(\rho_n + \bar{\rho}_n) \ln(\rho_n + \bar{\rho}_n) - \rho_n \ln \rho_n - \bar{\rho}_n \ln \bar{\rho}_n]. \quad (\text{B12})$$

Variation of (B2) leads to the relationship between $\delta\rho_n$ and $\delta\bar{\rho}_n$,

$$\delta\bar{\rho}_n = -A_{nm} \star \delta\rho_n, \quad (\text{B13})$$

which finally leads to the second TBA equation

$$\ln(1 + \eta_n) = \frac{2\pi J}{T} a_n + A_{nm} \star \ln(1 + \eta_m^{-1}), \quad (\text{B14})$$

where $\eta_n = \bar{\rho}_n/\rho_n$. Importantly, since the right-hand side of (B2) does not depend on $\delta\rho_n$ or on $\delta\bar{\rho}_n$, Eq. (B14) and hence η_n are *independent of the inhomogeneities*. It is customary to recast (B14) in terms of the dressed energies defined by $\varepsilon_n = T \ln(\eta_n)$,

$$\frac{\varepsilon_j}{T} = \frac{\varepsilon_j^{(0)}}{T} + a_{nm} \star \ln(1 + e^{-\varepsilon_m/T}), \quad (\text{B15})$$

where the bare energies are $\varepsilon_j^{(0)} = 2\pi J a_n$. In [58] a hydrodynamic approach to the Drude weight(s) has been formulated based on the TBA approach. The starting point is the existence of a basis of local charges \hat{Q}_i and associated currents J_i . Using these charges a *generalized Gibbs ensemble* is defined by the density matrix

$$\rho_{\text{GGE}} \sim \exp\left(-\sum_n \mu_n \hat{Q}_n\right), \quad (\text{B16})$$

where μ_i are ‘‘chemical potentials.’’ The generalized Drude weights D_A are then obtained from appropriate expectation

values in this ensemble and are determined by using the TBA method [62]. According to Ref. [58] in integrable models D_A can be expressed as

$$D_A = \sum_n \int d\lambda \frac{\eta_n(\lambda)}{\rho_n^{\text{tot}}(\lambda)} \left(\frac{\varepsilon'_n(\lambda) q_A^{\text{eff}}(\lambda)}{2\pi[1 + \eta_n(\lambda)]} \right)^2, \quad (\text{B17})$$

where

$$q_A^{\text{eff}} = \partial_{\mu_A} \ln \eta_n \quad (\text{B18})$$

are effective transport charges. The functions ε'_n are derivatives of the energies of elementary excitations over the state of thermal equilibrium and were calculated in Ref [63]. They are obtained from the dressed energies by solving a set of linear integral equations:

$$\begin{aligned} \varepsilon'_j \star (1 - K)_{jk}(\lambda) &= \frac{d\varepsilon_k^{(0)}(\lambda)}{d\lambda}, \\ K_{jk}(x, y) &= -\text{sgn}(q_j) a_{jk}(x - y) (1 + e^{\varepsilon_j/T})^{-1}. \end{aligned} \quad (\text{B19})$$

The only quantities in (B17) that depend on the inhomogeneities ξ_{2j} are the total densities $\rho_n^{\text{tot}}(\lambda)$. This can be seen from (B2) once the disorder-independent equations (B14) for η_n have been solved. It follows that the disorder averaging of the generalized Drude weights can be interchanged with the integration and summation in (B17). Introducing

$$\overline{\left(\frac{1}{\rho_n^{\text{tot}}(\lambda)}\right)} = \int P(\{\xi\}) \frac{1}{\rho_n^{\text{tot}}(\lambda, \{\xi\})}, \quad (\text{B20})$$

where $P(\xi)$ is a disorder probability distribution, we then can express disorder-averaged Drude weights in the form

$$\bar{D}_A = \sum_n \int d\lambda \overline{\left(\frac{1}{\rho_n^{\text{tot}}(\lambda)}\right)} \eta_n(\lambda) \left(\frac{\varepsilon'_n(\lambda) q_A^{\text{eff}}(\lambda)}{2\pi[1 + \eta_n(\lambda)]} \right)^2.$$

As the total density $\rho_n^{\text{tot}}(\lambda)$ is a positive quantity D_A is only renormalized due to the dependence of the string particle and string hole densities on disorder, and will not vanish unless the disorder probability distribution is fine-tuned.

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