

Symmetric Tensor Scars with Tunable Entanglement from Volume to Area Law

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Teleportation of quantum information over long distances requires robust entanglement on the macroscopic scale. The construction of highly energetic eigenstates with tunable long-range entanglement can provide a new medium for information transmission. Using a symmetric superposition of the antipodal triplet states, we construct polynomially many exact zero-energy eigenstates for a class of nonintegrable spin-1/2 Hamiltonians with two-body interactions. These states exhibit nonthermal correlations, and hence, are genuine quantum many-body scars. By tuning the distribution of triplets, we induce extensive, logarithmic, or area-law entanglement and can observe a second-order entanglement phase transition. Quasiparticle excitations in this manifold converge to be exact quantum many-body scars in the thermodynamic limit. This framework has a natural extension to higher dimensions, where entangled states controlled by lattice geometry and internal symmetries can result in new classes of correlated out-of-equilibrium quantum matter. Our results provide a new avenue for entanglement control and quantum state constructions.

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Introduction—It was previously conjectured that all thermalizing nonintegrable quantum many-body systems satisfy a strong eigenstate thermalization hypothesis (ETH) [1–3], wherein initial conditions generically thermalize and are unable to retain long-range quantum correlations. In recent years, many exceptions have been found, in a wide range of models featuring ETH-violating eigenstates, known as quantum many-body scars (QMBSs) [4–15], following a ground-breaking experiment in a Rydberg atom quantum simulator [16]. This phenomenon can dramatically slow down thermalization for weakly entangled initial states and preserve quantum correlations, including topological order [10,11,17].

Traditionally, these examples have featured area-law or logarithmic entanglement scaling, which can be understood as an embedding of ground states of a particular Hamiltonian deep into the many-body spectrum of a nonintegrable model [18]. More recently, eigenstates with volume-law entanglement have attracted significant interest, as they usually satisfy ETH. Nonetheless, atypical volume-law states can encode quantum information in accessible operators, which can be robust despite being highly entangled. Understanding the exact structure of these states

can help identify observables that behave athermally hidden in the complex entanglement features and can potentially be utilized for information storage and communication. In this respect, Bell pairs have proven to be useful building blocks for such highly entangled states. Rainbow scars [19–22] are a particular example that are formed by a concentric arrangement of Bell pairs over arbitrary distances and exhibit a volume law for a large majority of bipartitions. Similar constructions play a role in preparing thermofield double states [23–26] describing the interior of black holes. In periodic spin chains, volume-law states can also be prepared using entangled antipodal pairs of spins [27–29] in both integrable [30–32] and nonintegrable models [33]. Such states have also been known as “crosscap” states in conformal field theory for a long time [34].

Many-body states formed by singlet coverings have a rich history in condensed matter physics. The celebrated Majumdar-Ghosh state [35,36] is the frustration-free ground state of a Heisenberg-type spin chain. This state consists of a nearest-neighbor (n.n.) singlet covering and exhibits valence bond order. In 2D, such order can be destroyed by superposing all n.n. singlet coverings, forming a quantum spin liquid [37–40], realized at the Rokhsar-Kivelson point in quantum dimer models [41–43]. The singlets can be long-ranged as well [44]. In fact, a spin liquid can be obtained from an ordered state by tuning the length distribution of the singlets [45]. The possibility of realizing states that are a superposition of Bell pair coverings at higher energies, particularly at infinite temperature, has attracted limited attention.

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Quantum many-body systems with spectral reflection symmetry can host a zero-energy manifold exponentially large in system size [4,5,46]. These states lie at infinite temperature and are expected to satisfy ETH. Nevertheless, a highly degenerate manifold offers the possibility of constructing QMBs by a suitable superposition of states in this manifold [7,47,48]. Staggered Hamiltonians have also recently emerged as possible hosts of quantum scars [9,49–51]. In this Letter, we construct a large variety of exact zero-energy eigenstates for a class of staggered Heisenberg models in one dimension by superposing long-range triplet coverings. Their entanglement can be tuned from volume law to area law by controlling the distribution of triplets. Our construction is generalizable to higher dimensions and is stable to several forms of symmetry-breaking perturbations.

Model and the symmetric tensor scars—Consider a class of bond-staggered Heisenberg Hamiltonians with long-range interactions on a periodic chain of size $2N$,

$$H = \sum_{i=1}^{2N} (-1)^i \mathbf{S}_i \cdot \mathbf{S}_{i+r}, \quad (1)$$

where $\mathbf{S} = (\sigma_x, \sigma_y, \sigma_z)/2$ is a vector of spin-1/2 operators, and $r \in [0, N-1]$ is the range of interactions. H is invariant under translation by two lattice sites, reflection over a bond, and has an internal $SU(2)$ symmetry. The nearest-neighbor version of the model has a long history; the ground state belongs to the Haldane phase of gapped symmetry protected topological order [52–54], which was also realized experimentally [55].

We define the following many-body states, which we designate as *root states*,

$$|\Psi(v)\rangle = v^{\otimes N}, \quad (2)$$

where v is a state of two spins at antipodal sites i and $i+N$. We find that Eq. (2) is a zero-energy eigenstate of Eq. (1), provided that N is odd and v is either the singlet state or any triplet state [56]. The singlet root state $|\Psi(S)\rangle$ belongs to $S_{\text{tot}}^z = 0$, the largest symmetry sector of Eq. (1) with maximal magnon number, where $N_{\text{mag}} = N - \langle |S_{\text{tot}}^z| \rangle$ (deficit in total spin polarization captures the average number of magnons in the state). Construction of exact zero-energy eigenstates of Eq. (1) with high magnon numbers has been a recent challenge [49,51], which we overcome in this Letter. The root states $|\Psi(T)\rangle$, constructed from some triplet representation T , are generically not eigenstates of S_{tot}^z and have no well-defined magnon numbers, but they will typically have an extensive average number of magnons. We note that certain root states can also be annihilated by infinitely many odd conserved charges of the XXX spin chain [30]; hence, they can be treated as integrable boundary states to construct new scar models [60,61].

The linear span of the root states (excluding the singlet root state) is isomorphic to the space of symmetric tensors of rank- N over the three-dimensional vector space of triplets, $\text{Sym}^N(V_{S=1})$, with dimension $\mathcal{N} = \binom{N+2}{N} = (N^2 + 3N + 2)/2$. Hence, this subspace scales quadratically with the system size and is much smaller than the exponentially large space of all zero-energy states.

To construct a family of complex zero-energy eigenstates, we take the innovative step of introducing superpositions of triplet coverings. We show that choosing different triplets for the antipodal pairs and organizing them in a symmetric linear superposition results in a family of exact zero-energy eigenstates with an intricate web of quantum correlations. These states form a complete basis for the *symmetric tensor scars* and can be labeled by three positive integers, n_1, n_2, n_3 (the number of triplets of each kind), which sum to N and are formally represented as

$$|\Psi_{n_1, n_2, n_3}^{\text{sym}}\rangle = \frac{1}{\mathcal{N}_c} \sum_{\pi \in S_N} \pi(|T_1\rangle^{\otimes n_1} \otimes |T_2\rangle^{\otimes n_2} \otimes |T_3\rangle^{\otimes n_3}), \quad (3)$$

where $\pi(\cdot)$ is a permutation of the tensor-power factors (an element of the symmetric group of N objects, S_N). The normalization $\mathcal{N}_c = \sqrt{N!n_1!n_2!n_3!}$ can be computed from the orbit-stabilizer theorem. We show one such state in Fig. 1(a). The anomalous nature of these states (scars) will become evident shortly.

Although there are infinitely many choices for the triplet basis, here we consider the following two, which serve as examples of the most and the least entangled states,

$$|T_X^B\rangle_i = \frac{1}{\sqrt{2}}(|\uparrow_i \uparrow_j\rangle + |\downarrow_i \downarrow_j\rangle); \quad |T_+^C\rangle_i = |\uparrow_i \uparrow_j\rangle, \quad (4a)$$

$$|T_Y^B\rangle_i = \frac{1}{\sqrt{2}}(|\uparrow_i \uparrow_j\rangle - |\downarrow_i \downarrow_j\rangle); \quad |T_-^C\rangle_i = |\downarrow_i \downarrow_j\rangle, \quad (4b)$$

$$|T_Z^B\rangle_i = \frac{1}{\sqrt{2}}(|\uparrow_i \downarrow_j\rangle + |\downarrow_i \uparrow_j\rangle); \quad |T_Z^C\rangle_i = |T_Z^B\rangle_i, \quad (4c)$$

where $j \equiv i+N$, while $|\uparrow\rangle$ and $|\downarrow\rangle$ are eigenvectors of S^z . We call T^B (T^C) the Bell pair (conventional) basis for the triplets. Subsequently, we represent a family of symmetric tensor states in these two bases by (n_X, n_Y, n_Z) and (n_+, n_-, n_Z) , respectively.

Entanglement entropy—We focus mainly on the Rényi-2 entanglement entropy $S^{(2)}$ for the symmetric tensor state, because unlike with the von Neumann entropy, S^{vN} , we have developed statistical mechanical techniques for its efficient computation and analytic treatment [56], enabling us to deduce the entanglement scaling. We are particularly interested in the scaling behavior with system size and whether it is thermal or athermal. S^{vN} is found to follow similar behavior as $S^{(2)}$ for smaller system sizes [56].

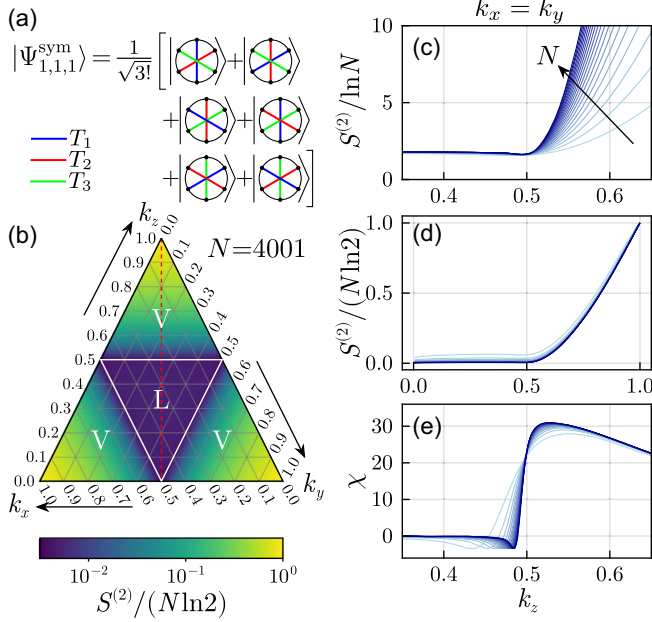


FIG. 1. (a) A symmetric tensor state for $N = 3$, a zero-energy eigenstate of Eq. (1). Each component is a tensor product of three triplet types (T_1 , T_2 , T_3) between antipodal sites. (b) Phase diagram for the Bell basis states. Color indicates the half-system entropy $S^{(2)}$ for $N = 4001$. White lines indicate the boundary between volume-scaling entanglement (V) and log-scaling entanglement phases (L). The red dashed line indicates $k_x = k_y$, which is shown in more detail in panels (c)–(e), wherein increasing saturation indicates increasing N from 201 to 4001 in steps of 200. Along this line, we show entanglement rescaled by (c) $\ln N$ and (d) N to demonstrate the two scalings. (e) Susceptibility $\chi = (1/N)d^2 S^{(2)}/dk_z^2$ exhibits a crossing evidencing a second-order phase transition.

We find that the half-chain $S^{(2)}$ assumes a range of scaling behaviors, summarized in Fig. 1 for the Bell basis and Fig. 2 for the conventional basis. In addition to the volume law, we find examples of log-law (both for the Bell and conventional bases) and area-law (for conventional basis only) states.

In the Bell basis, we find a precise rule to determine which states have volume-law or log-law entanglement. We represent a scaling family of states by a triplet (k_x, k_y, k_z) by setting $n_\alpha = k_\alpha N$. Whenever the largest k_α is greater than the sum of the other two, the states exhibit volume law, and otherwise, log law [see the phase diagram in Fig. 1(b)]. In Fig. 1(c)–(e), we show how the entropy changes along the red line marked on the diagram. We examine $S^{(2)}$ rescaled by (c) $\ln N$ and (d) N to show that the scaling is indeed logarithmic and volume in the two phases. In Fig. 1(e), we show the susceptibility, $\chi = (1/N)d^2 S^{(2)}/dk_z^2$ (identifying $S^{(2)}$ as related to the free energy of a statistical mechanical model [62–64]), which exhibits a crossing indicating a second-order entanglement phase transition (akin to famous examples in localized models [65,66] and monitored systems [67–70]).

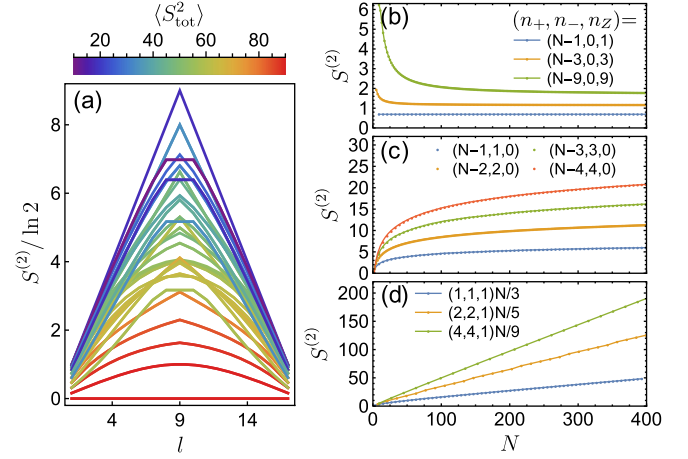


FIG. 2. Entanglement for the symmetric tensor states in the conventional basis. (a) $S^{(2)}$ vs subsystem size l for all states with $N = 9$. The color denotes $\langle S_{\text{tot}}^2 \rangle$ of the corresponding state. Example states showcasing the (b) area-law, (c) log-law, and (d) volume-law entanglement scaling of half-chain entropy. Lines in (c) are fits to $a \ln N + b$.

In the conventional basis, the entanglement ranges from zero to maximal [see Fig. 2(a)], and we find families of states covering the entire spectrum of entanglement scaling behaviors, including area, log, and volume law. An example of area law is shown in Fig. 2(b), where the states $(N - n_z, 0, n_z)$ with fixed n_z have entanglement bounded by a constant, $\ln[(\sqrt{\pi n_z!})/\Gamma(n_z + 1/2)]$. States of type $(N - n_-, n_-, 0)$ with fixed n_- exhibit logarithmic entanglement [see Fig. 2(c)], following the form $S^{(2)} = n_- \ln(N) - \ln n_-!$ (the $n_- = 1$ case was previously studied by us in [49]). Some volume-law states have a counter-intuitive behavior; for example, the state $(1, 1, 1)N/3$ shows a volume-law scaling with a small coefficient, but the entanglement grows faster when we increase the proportion of the two product triplets [see Fig. 2(d)]. Such anomalous growth of entanglement originates from the symmetric superposition. We find the entanglement scaling to become volume law when both the product triplet numbers $\{n_+, n_-\}$ scale extensively; e.g., the state $(N \pm 1, N \mp 1, 0)/2$ has $S^{(2)} \approx N \ln 2 - 1/(N + 1)$, and hence, is maximally entangled in the thermodynamic limit. This demonstrates how the symmetrization of triplet states can generate a hierarchy of entanglement scalings, which exemplifies the rich structure of the space of symmetric superpositions.

Local observables and correlation functions—We now look at the behavior of local observables and whether they imply a thermal behavior of symmetric tensor states. Expectation values of single-site observables are zero in the Bell basis, $\langle S_i^\alpha \rangle = 0$, since every site is maximally entangled with the rest of the system in all states. This matches the infinite-temperature average of these quantities, and one may be tempted to conclude these states as thermal.

However, there are subtleties depending on the range of observables due to the following reasons. First, the root states [Eq. (2)] are product states for specific nonlocal and noncontiguous bipartitions, and consequently, are strongly atypical. Secondly, antipodal correlators assume nonthermal values, such as $C[N] = 1/4$ (where $C[l] = \langle \mathbf{S}_i \cdot \mathbf{S}_{i+l} \rangle$) for all the symmetric tensor states. Conventionally, the thermal nature of a state is probed via local correlation functions while $C[N]$ is nonlocal, and thus, experimentally inaccessible in thermodynamically large systems, though within the setting of local operation and classical communication [71], their athermal character can be probed. Furthermore, we find that at any finite system size, $C[l < N]$ is a nonzero constant independent of l (except for the root states). Hence, all such states are long-range ordered [72–74] and atypical at any finite system size. As a consequence, $C[l]$ is a single-variable function of $\langle S_{\text{tot}}^2 \rangle$, increasing linearly with it, $C[l] = \langle S_{\text{tot}}^2 \rangle / [4N(N-1)] - 1/(2N-2)$.

The scaling of local correlation functions with system size is crucial for the ultimate fate of their thermal nature. In the Bell basis, we obtain $C[l] = (n_X n_Y + n_Y n_Z + n_Z n_X) / (2N^2 - 2N)$. This vanishes as $1/N$ for states with one of the n_i scaling as N while the other two are constant. However, as soon as two of the three n_i scale with N , $C[l]$ assumes a nonzero constant value even in the thermodynamic limit. For example, $C[l]$ for the states $(2N/3, N/3, 0)$ and $(N/3, N/3, N/3)$ approaches $1/9$ and $1/6$, respectively [56], in the thermodynamic limit; this proves (even in the conventional sense) that these states are nonthermal. Interestingly, though the latter state follows logarithmic entanglement scaling in agreement with the behavior of many atypical states, the former exhibits volume law, and hence, can be referred to as exceptional QMBS with volume-law entanglement, going beyond the previously studied cases of volume-law scars [27,28].

In the conventional basis, the correlations are given by $C[l] = [(n_+ - n_-)^2 + (n_+ + n_-)(2n_Z - 1)] / [2N(2N - 2)]$. Due to the presence of product states in this basis, $C[l]$ is generically nonthermal ($C[l] \neq 0$) even in the thermodynamic limit. There are two limiting cases when correlations exhibit thermal behavior: (1) n_+ , n_- are constant, and (2) n_+ , n_- are both extensive and differ by $O(1)$ with a constant n_Z . While both cases have volume-law entanglement, the former belongs to the same class as the entangled-antipodal-pair states, while the latter has a more complex structure. Note that $C[l]$ can be negative for some states (with $n_Z = 0$ and $|n_+ - n_-| < \sqrt{N}$) in this basis for any finite system size. To summarize, we have shown that the symmetric tensor states in general exhibit nonthermal behavior in their correlations.

Quasiparticle excitations and asymptotic QMBS—We consider quasiparticle excitations [6,7,15,73] on top of the singlet root state, generated by a local operator $Q_j = \sigma_{\alpha,j}$, where $\alpha \in \{X, Y, Z, +, -\}$, the action of which on one singlet is to replace it with the appropriate triplet. We find

$S_{\text{tot}}^2 Q_j |S\rangle = 2Q_j |S\rangle$, hence, this excitation carries spin-1 and should be referred to as a triplon. We can take linear combinations of these operators to find a single-particle wave function that represents an asymptotically stable excitation (or asymptotic QMBS [75]),

$$|Q\rangle = \frac{1}{\sqrt{N}} \sum_{j=1}^N Q_j |S\rangle. \quad (5)$$

This wave function is a square wave, as the superposition is on only half the system. We find $E_Q = \langle Q | H | Q \rangle = 0$ and the energy variance $\delta = \langle Q | H^2 | Q \rangle - E_Q^2 = 2/N$ (for $r = 1$). Therefore, although $|Q\rangle$ is not an eigenstate of H at any finite system size, the lifetime of the quasiparticle diverges in the thermodynamic limit. This can be understood intuitively as follows: the energy fluctuations are proportional to the gradient of the wave function, so the most stable excitations will vary spatially as little as possible. The antisymmetric nature of the singlet state, however, demands that the wave function be odd under translation by N as the even-transforming component annihilates $|S\rangle$. This mandates some minimum amount of variation into the wave function. In the thermodynamic limit, however, the energy fluctuations become negligible because almost all of the wave function is far away from the edges of the square wave. Intriguingly, a family of related quasiparticle eigenstates for this model can be constructed using a generalized Bethe ansatz [51].

Extension to higher dimensions—The state construction relies on the cancellation between the action of a local term in the Hamiltonian with its antipodal partner, while the structure of the rest of the lattice is mostly immaterial. Hence, higher-dimensional generalizations are fairly straightforward, but they also provide a new way to create antipodal pairs. For example, a rectangular lattice of $2N_x \times 2N_y$ sites (with a specific choice of the staggered interactions and odd N_x and even N_y) possesses two antipodal pairs: $(x, y) - (x + N_x, y + N_y)$ and $(x, y) - (x + N_x, y)$. Given this setup, all symmetric tensor states are zero-energy eigenstates of the Hamiltonian. Thus, higher-dimensional lattices allow the superposition of multiple “flavors” of symmetric states, potentially leading to a liquidlike behavior. Extensions to cubic and hypercubic lattices are similarly straightforward.

Stability against perturbations—We find that many symmetric tensor scars remain zero-energy eigenstates even in the presence of $SU(2)$ symmetry-breaking perturbations, such as uniform and staggered fields, or easy-axis anisotropy of the exchange interaction. States that are no longer eigenstates in the presence of such perturbations remain stable up to significantly long times for moderate values of the perturbations, while others exhibit coherent oscillations in the presence of an additional staggered field (see Ref. [56]).

Discussion—In this Letter, we have shown the existence of a class of exact quantum many-body scars at infinite temperature in a family of models with staggered Heisenberg interaction. These symmetric tensor scars are constructed by long-range triplet coverings of the system. Some of these (root) states exhibit a mixture of thermal (i.e., volume-law scaling of entanglement and thermal values of local observables) and nonthermal (product structure for certain bipartite entanglement cuts, strictly nonthermal antipodal correlation) properties. We demonstrate how to regularize such anomalous QMBs by inducing a nonthermal local expectation value via symmetric superposition of different triplet coverings. We note that the relationship between the frame of root states and the symmetric tensor scars mirrors the relationship between the mean-field time-dependent variational principle frame [76] used to study scar states in the PXP model and the permutation-invariant scar quasimodes [77] that result from the quantization of that semiclassical frame. The ergodic properties of the states are sensitive to the choice of the triplet basis. While the Bell pair basis admits log and volume-law scaling but no area-law states, the conventional basis supports the full spectrum of entanglement behaviors. Many of these states are also stable against symmetry-breaking perturbations [56]. Quasiparticles on top of the singlet root state are found to be asymptotic QMBs. We also discuss the recipe for generalization to higher dimensions.

Our results open up numerous diverse research directions. First and foremost, our recipe for creating highly entangled states unlocks the way of building exact many-body scars beyond the usual area-law paradigm. In fact, since in the conventional triplet basis, the states admit any behavior from area to volume law, one can construct states with exotic entanglement scaling [such as $\sim \ln^2 N$ or a fractal entanglement $\sim N^\alpha$, $\alpha \in (0, 1)$]. This allows analytical insights into unusual families of states, such as those exhibiting quantum critical properties or multifractality. Additionally, similar symmetric tensor constructions are viable, such as those involving a superposition of three or more spin state coverings, although whether such constructions would lead to new eigenstates is an open question. Beyond this, the possibility of easily creating scars in higher-dimensional systems raises the question of whether one can produce topological symmetric tensor scars, potentially providing analytical leverage on non-trivial anyonic excitations (extending our Letter on quasiparticles) and other topological phenomena. These novel state construction techniques can not only lead to theoretical insights into complex dynamical properties, but also provide a framework for stabilizing quantum order in thermal systems. The simulation of such symmetric tensor scars in near-term quantum computers is also an interesting future avenue to explore [78–80].

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Data availability—The data that support the findings of this article are openly available [81].

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