Supercritical entanglement in local systems: Counterexample to the area law for quantum matter

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Supercritical entanglement in local systems: Counterexample to the area law for quantum matter

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Quantum entanglement is the most surprising feature of quantum mechanics. Entanglement is simultaneously responsible for the difficulty of simulating quantum matter on a classical computer and the exponential speedups afforded by quantum computers. Ground states of quantum many-body systems typically satisfy an “area law”: The amount of entanglement between a subsystem and the rest of the system is proportional to the area of the boundary. A system that obeys an area law has less entanglement and can be simulated more efficiently than a generic quantum state whose entanglement could be proportional to the total system's size. Moreover, an area law provides useful information about the low-energy physics of the system. It is widely believed that for physically reasonable quantum systems, the area law cannot be violated by more than a logarithmic factor in the system's size. We introduce a class of exactly solvable one-dimensional physical models which we can prove have exponentially more entanglement than suggested by the area law, and violate the area law by a square-root factor. This work suggests that simple quantum matter is richer and can provide much more quantum resources (i.e., entanglement) than expected. In addition to using recent advances in quantum information and condensed matter theory, we have drawn upon various branches of mathematics such as combinatorics of random walks, Brownian excursions, and fractional matching theory. We hope that the techniques developed herein may be useful for other problems in physics as well.

Significance

We introduce a class of exactly solvable models with surprising properties. We show that even simple quantum matter is much more entangled than previously believed possible. One then expects more complex systems to be substantially more entangled. For over two decades it was believed that the area law is violated by at most a logarithm in the system's size for quantum matter (i.e., interactions satisfying physical reasonableness criteria clearly stated in the article). In this work we introduce a class of physically reasonable models that we can prove violate the area law by a square root, i.e., exponentially more than the logarithm.

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are solvable. This makes the model truly local (Eq. 6). We remark that the particle spins can be as low as $s = 2$ for $\sqrt{n}$ violation. We now describe this class of models and detail the proofs and further discussions in the SI Appendix. Let us consider an integer spin-$s$ chain of length $2n$. It is convenient to label the $d = 2s + 1$ spin states by $\{/.,/.,\ldots,/,0,\ldots,\cdot,\cdot\cdot\cdot\}$ as shown in Fig. 1. Equivalently, and for better readability, we instead use the labels $\{u^+,u^-,\ldots,u^s,0,d^s,d^s,\ldots,d^s\}$ where $u$ means a step up and $d$ a step down. We distinguish each type of step by associating a color from the $s$ colors shown as superscripts on $u$ and $d$.

A Motzkin walk on $2n$ steps is any walk from $(x,y) = (0,0)$ to $(x,y) = (2n,0)$ with steps $(1,0)$, $(1,1)$, and $(-1, -1)$ that never passes below the $x$ axis, i.e., $y \geq 0$. An example of such a walk is shown in Fig. 2. The height at the midpoint is $0 \leq m \leq n$, which results from $m$ steps up with the balancing steps down on the second half of the chain. In our model the unique ground state is the $s$-colored Motzkin state which is defined to be the uniform superposition of all $s$ colorings of Motzkin walks on $2n$ steps. The nonzero heights in the middle are the source of the mutual information between the two halves and the large entanglement entropy of the half-chain (Fig. 3).

The Schmidt rank is $(s^{|+|} - 1) / s - 1 \approx (s^{|+|}) / s - 1$, and using a 2D saddle-point method, the half-chain entanglement entropy asymptotically is (see SI Appendix for derivation)

$$S = 2 \log_2(s) \sqrt{\frac{2m}{\pi} + 1} + \log_2(2s\pi n) + \left(\gamma - \frac{1}{2}\right) \log_2 e \text{ bits,}$$

where $\gamma = \sqrt{s} / (2\sqrt{s} + 1)$ is constant and $\gamma$ is the Euler constant.

The ground state is a pure state (which we call the Motzkin state), whose von Neumann entropy is zero. However, the entanglement entropy quantifies the amount of disorder produced by the local terms, and the spin dimension is quite high (21).

As noted previously, a generic state violates the area law maximally (9). It was largely believed that the ground state of a generic state violates the area law by at least a $\log(n)$ factor, where $n$ is the number of particles (see ref. 28 for a review). Physically reasonable models need to have Hamiltonians that are (i) local, (ii) translationally invariant, and (iii) have a unique ground state. These requirements, among other things, eliminate highly fine-tuned models. This implies that $\log(n)$ is the maximum expected entanglement entropy in realistic physical spin chains.

In an earlier work, Bravyi et al. (29) proposed a spin-1 model with the ground-state half-chain entanglement entropy $S = (1/2) \log n + c$, which is logarithmic factor violation of the area law as expected during a phase transition. This model is not truly local as it depends crucially on boundary conditions. The scaling of the entanglement is exactly what one expects for critical systems.

We have found an infinite class of exactly solvable integer spin-$s$ chain models with $s \geq 2$ that are physically reasonable and exact calculation of the entanglement entropy shows that they violate the area law by the leading order by $\sqrt{n}$ (Eq. 1). The proposed Hamiltonian is local and translationally invariant in the bulk but the entanglement of the ground state depends on boundary projectors. We prove that it has a unique ground state and give a technique for proving the gap that uses universal convergence of random walks to a Brownian motion. We prove that the energy gap scales as $n^{-c}$, where using the theory of Brownian excursions we show that the constant $c \geq 2$. This bound rules out the possibility of these models being describable by a CFT. The Schmidt rank of the ground state grows exponentially with $n$.

We then introduce an external field. In the presence of the external field the boundary projectors are no longer needed. The model has a frustrated ground state, and its gap and entanglement
Therefore, the local Hamiltonian, with projectors as interactions, that has the Motzkin state as its unique zero-energy ground state is

$$H = \Pi_{\text{boundary}} + \sum_{j=1}^{2n-1} \Pi_{ij+1} + \sum_{j=1}^{2n-1} \Pi_{ij+1}^{\text{exact}},$$

where $\Pi_{ij+1}$ implements the local changes discussed above and is defined by

$$\Pi_{ij+1} = \frac{1}{2} \left[ |D^k\rangle_{ij+1} \langle D^k| + |U^k\rangle_{ij+1} \langle U^k| + |\varphi^k\rangle_{ij+1} \langle \varphi^k| \right].$$

Fig. 3. Motzkin walk with $s=2$ colors of length $2n=10$. The height $m$ quantifies the degree of correlation between the two halves.

The difference between the ground-state energy and the energy of the first excited state is called the gap. One says a system is gapped when the difference between the two smallest energies is at least a fixed constant in the thermodynamical limit ($n \to \infty$). Otherwise the system is gapless.

Whether a system is gapped has important implications for its physics. When it is gapless, the scaling by which the gap vanishes as a function of the system’s size has important consequences for its physics. For example, gapped systems have exponentially decaying correlation functions (22), and quantum critical systems are necessarily gapless (31). Moreover, systems that obey a CFT are gapless but the gap must vanish as $1/n$ (32). Therefore, to quantify the physics, it is desirable to find new techniques for analyzing the gap that can be applied in other scenarios.

The model proposed here is gapless and the gap scales as $n^{-c}$ where $c \geq 2$ is a constant. We prove this by finding two functions, both of which are inverse powers of $n$ such that the gap is always smaller than one of them (called an upper bound) and greater than the other (called a lower bound). We use techniques from mathematics such as Brownian excursions and universal convergence of random walks to a Brownian motion, as well as other ideas from computer science such as linear programming and fractional matching theory. We describe the ideas and leave the details of the proofs for SI Appendix.

To prove an upper bound on the gap, one needs a state $\langle \phi \rangle$ that has a small constant overlap with the ground state and such that $\langle \phi | H | \phi \rangle \geq O(n^{-2})$. Take

$$\langle \phi \rangle = \frac{1}{\sqrt{M_{2n}}} \sum_{m_p} e^{2\pi i m_p} |m_p\rangle,$$

where the sum is over all Motzkin walks, $M_{2n}$ is the total number of Motzkin walks on $2n$ steps, $A_{n}$ is the area under the Motzkin walk $m_p$, and $\theta$ is a constant to be determined by the condition of a small constant overlap with the ground state. The overlap with the ground state is defined by $\langle M_{2n} | \phi \rangle = (1/M_{2n}) \sum_{m_p} e^{2\pi i m_p}$. As $n \to \infty$, the random walk converges to a Wiener process (33) and a random Motzkin walk converges to a Brownian excursion (34). We scale the walks such that they take place on the unit interval. The scaled area is denoted by $A$ and $\theta \to \theta$. In this limit, the overlap becomes (see Fig. 4 for the density and Fig. 5 for its Fourier transform; $F_A(\theta)$ is the Fourier transform of the probability density function, which is called the characteristic function.)

$$\lim_{n \to \infty} \langle M_{2n} | \phi \rangle = F_A(\theta) \equiv \int_0^\infty f_A(x) e^{2\pi i x \theta} dx,$$

where $f_A(x)$ is the probability density function for the area of the Brownian excursion (35) shown in Fig. 4. In Eq. 5, taking $\theta \ll O(1)$, gives $\lim_{n \to \infty} \langle M_{2n} | \phi \rangle \approx 1$ because it becomes the integral of a probability distribution. However, taking $\theta \gg O(1)$ gives a highly oscillatory integrand that nearly vanishes. To have a small constant overlap with the ground state, we take $\theta$ to be the standard of deviation of $f_A(x)$. Direct calculation then gives $\langle \phi | H | \phi \rangle = O(n^{-2})$. This upper bound decisively excludes the possibility of the model being describable by a conformal field theory (18).

Using various ideas in perturbation theory, computer science, and mixing times of Markov chains we obtain a lower bound on the gap that scales as $n^{-c}$, where $c \gg 1$. Because it might be of independent interest in other contexts, we present a combinatorial and self-contained exposition of the proof in the SI Appendix, different in some aspects from that given in ref. 29.

The model above has a unique ground state because the boundary terms select out the Motzkin state among all other walks with different fixed initial and final heights. Without the boundary projectors, all walks that start at height $m_1$ and end at height $m_2$ with $-2n \leq m_1, m_2 \leq 2n$ are ground states. For example, when $s = 1$, the ground-state degeneracy grows quadratically with the system’s size $2n$ and exponentially when $s > 1$.

For the $s = 1$ case, if we impose periodic boundary conditions, then the superposition of all walks with an excess of $k$ up (down) steps is a ground state. This gives $4n+1$ degeneracy of the ground state, which includes unentangled product states.

![Fig. 4. Plot of the probability density of the area under a Brownian excursion $f_A(x)$ on [0,1).](image-url)
When \( s > 1 \), each one of the walks with \( k \) excess up (down) steps can be colored exponentially many ways; however, generically they will not be product states. Consider an infinite chain \( (-\infty, \infty) \) and take \( s > 1 \). There is a ground state of this system that corresponds to the balanced state, where on average for each color, the state contains as many \( u \) as \( d \). Suppose we restrict our attention to any block of \( n \) consecutive spins. This block contains the sites \( j, j + 1, \ldots, j + n - 1 \), which is a section of a random walk. Let us assume that it has initial height \( m_j \) and final height \( m_{j+n-1} \). Further, let us assume that the minimum height of this section is \( m_k \) with \( j \leq k \leq j + n - 1 \). From the theory of random walks, the expected values of \( m_j - m_k \) and of \( m_{j+n-1} - m_k \) are \( \Theta(\sqrt{n}) \). The color and number of any unmatched step-ups in this block of \( n \) spins can be deduced from the remainder of the infinite walk. Thus, a consecutive block of \( n \) spins has an expected entanglement entropy of \( \Theta(\sqrt{n}) \) with the rest of the chain. A similar argument shows that any block of \( n \) spins has an expected half-block entanglement entropy of \( \Theta(\sqrt{n}) \).

If we take \( s = 1 \), where the ground state can be a product state, the \( \sqrt{n} \) unmatched step-up just mentioned can be matched anywhere on the remaining left and right part of the chain. Two consecutive blocks of \( n \) spins can be unentangled because the number of unbalanced steps that are matched in the next block is uncorrelated with the number of unbalanced steps in the first block. However, when \( s > 1 \) the ordering has to match. Even though the number of unbalanced steps in two consecutive blocks is uncorrelated, the order of the types of unbalanced steps in them agrees.

The Hamiltonian without the boundary terms is truly translationally invariant, yet has a degenerate ground state. We now propose a model with a unique ground state that has the other desirable properties of the model with boundaries, such as the gap and entanglement entropy scalings as before. To do so, we put the system in an external field, where the model is described by the Hamiltonian

\[
\hat{H} \equiv H + \epsilon F \\
F \equiv \sum_{n=1}^{2} \sum_{k=1}^{n} \langle \sigma^+ \rangle_{\sigma^x} + \langle \sigma^z \rangle_{\sigma^z},
\]

where \( H \) is as before but without the boundary projectors and \( \epsilon = \epsilon_0/n \) with \( \epsilon_0 \) being a small positive constant. It is clear that \( F \) treats \( u \) and \( d \) symmetrically; therefore, the change in the energy as a result of applying an external field depends only on the total number of unbalanced steps denoted by \( m \). We denote the change in the energy of \( m \) unbalanced steps by \( \Delta E_m \). When \( s = 1 \), the degeneracy after applying the external field will be one for the Motzkin state, two-fold when there is a single imbalance, three-fold for two imbalances, etc. Because the energies are equal for all \( m \) imbalance states, it is enough to calculate the energy for an excited state with \( m \) imbalances resulting only from excess step-ups. We denote these states by \( |g_m\rangle \), where \( 0 \leq m \leq 2n \).

The first-order energy corrections, obtained from first-order degenerate perturbation theory, are analytically calculated to be

\[
\epsilon \langle g_m | F | g_m \rangle \approx 4 \sigma \epsilon_0 + \frac{m e}{8 \sqrt{s}} \left( \frac{m}{n} \right). \tag{7}
\]

The physical conclusion is that the Hamiltonian without the boundary projectors, in the presence of an external field \( F \), has the Motzkin state as its unique ground state with energy \( 4 \sigma \epsilon_0 \).

Moreover, what used to be the rest of the degenerate zero-energy states acquire energies above \( 4 \sigma \epsilon_0 \) that for first elementary excitations scale as \( 1/n^2 \). Moreover, the numerical calculations indicate that the spin–spin correlation functions are flat (36). We leave further investigations for future work.

The energy corrections just derived do not mean that the states with \( m \) imbalances will make up for all of the low-energy excitations. For example, when \( s > 1 \), in the presence of an external field, the energy of states with a single crossed term will be lower than those with large \( m \) imbalances and no crossings.

Because \( |\epsilon F| \ll |H| \), the ground state will deform away from the Motzkin state to prefer the terms with more zeros in the superposition. But, as long as \( \epsilon \) is small, the universality of Brownian motion guarantees the scaling of the entanglement entropy. It is, however, not yet clear to us whether \( \epsilon \) can be tuned to a quantum critical point where the ground state has a sharp transition from highly entangled to nearly a product state. It is possible that the transition is smooth and that the entanglement continuously diminishes as \( \epsilon \) becomes larger. For example, in the limit where \( |\epsilon| \gg |H|/|F| \), the effective unperturbed Hamiltonian is approximately \( F \), whose ground state is simply the product state \( |\psi\rangle \otimes \mathcal{B} \).

Our model shows that simple physical systems can be much more entangled than expected. From a fundamental physics perspective, it is surprising that a 1D translationally invariant quantum spin chain with a unique ground state has about \( \sqrt{n} \) entanglement entropy. Moreover, this adds to the collection of exactly solvable models from which further physics can be extracted. Such a spin chain can in principle be experimentally realized, and the large amount of entanglement may be used as a resource for quantum technologies and computation.

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