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Shaping causality: programmable nonlocal signal generation in long-range spin systems

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Understanding how information spreads in non-relativistic many-body systems is a central issue for quantum information processing [1, 2]. While short-range interactions confine information within a local light cone [3], long-range interactions typically lead to uncontrolled nonlocal spread across the entire system [4, 5]. Here, we demonstrate that this apparent dichotomy is not fundamental and that nonlocality in systems with long-range interactions can be deterministically controlled. By mapping spin dynamics to a hard-core boson chain, we identify a regime in which the causal space-time landscape can be precisely shaped. We show that placing spin excitations in a polarized background allows a local perturbation to trigger nonlocal signals exactly at the positions of these excitations. These pre-selected sites act as seeds for new, effective light cones, allowing information to bypass the bulk and re-emerge at distant, programmable locations. This mechanism avoids uncontrollable global nonlocality while circumventing the speed limits associated with local transport. By engineering these nonlocal communication channels, our findings offer a versatile framework for information distribution relevant to quantum memories, error correction, and programmable platforms such as trapped ions [2, 6, 7].

Keywords: emergent locality, long-range interacting spin systems, quantum control

Introduction— Although non-relativistic quantum systems are not fundamentally constrained to exhibit a finite speed for information propagation, information in many-body systems with short-range interactions nevertheless spreads in a locally bounded manner, as formalized by the Lieb–Robinson bounds [3]. These bounds dictate that the spread of information occurs at most ballistically within an effective light cone, with a velocity determined by the coupling strength, and with leakage of information outside the light cone being at least exponentially suppressed. Thus, locality limits information propagation and shapes our understanding of entanglement and correlation growth, thermalization, and causality [3, 8–11].

Many programmable quantum platforms [12] naturally realize long-range couplings. These include trapped ions [1, 2, 5, 6, 13, 14], Rydberg atom arrays [15–17], dipolar molecules [18, 19], and optical cavity QED systems [20–23]. Such interactions can enable faster-than-ballistic spread of information, suggesting a breakdown of locality [1, 4, 5, 24–26].

However, simply increasing the interaction range often leads to uncontrolled spread of information [27–31]. To control this behavior and engineer useful quantum channels, one requires a mechanism that suppresses generic nonlocal transport while enabling targeted signaling. We demonstrate that this can be achieved by exploiting the interplay between long- and short-range interactions.

Here, we show that for a class of long-range 1D spin chains, the causal space-time landscape can be precisely shaped. We demonstrate the existence of a parameter regime (large system size or strong long-range coupling) where the onset of nonlocality is fully programmable. To illustrate this mechanism, we consider a transmission protocol involving three parties: a sender (Alice), an intermediate observer (Eve), and a distant receiver (Bob) [Fig. 1(a)]. While short-range interactions require Alice’s signal to traverse every intermediate site to reach Bob, our long-range architecture breaks this constraint. By tuning the initial state, Alice can trigger a nonlocal signal that “jumps” over Eve entirely, appearing at Bob’s site while leaving Eve undisturbed. Once it emerges at the target site, the signal reverts to local propagation within a linear light cone. This creates a direct, nonlocal communication channel that bypasses the geometric bulk of the chain.

Crucially, this nonlocal signaling is deterministic and scalable: by placing additional background excitations

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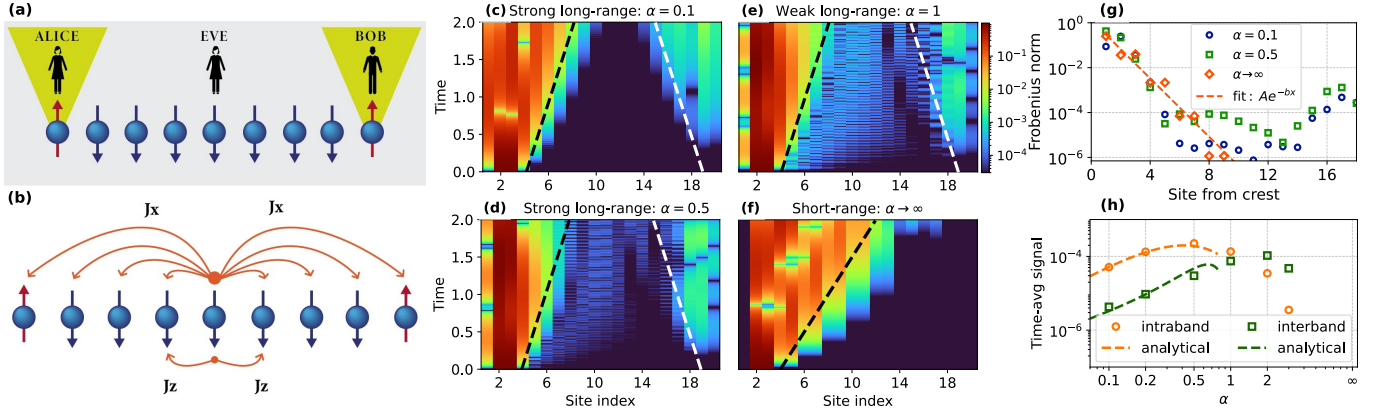


FIG. 1. **Nonlocal signal generation.** Illustration of: (a) nonlocal signaling, (b) the model. (c–h) Emergence of nonlocal light cones for $N = 20$, $J_x = 5$, and $J_z = 1$. (c–f) Density plots of the Frobenius norm $\|\Delta\rho_n(t)\|_F$ showing light-cone dynamics under the Hamiltonian (1) for initial states prepared in band 2, (2, 19; 3, 19) (notation defined in Eq. (4)). Dashed lines (black and white for contrast) indicate the analytically predicted light-cone velocities: $2J_z$ for interaction exponent $\alpha \leq 1$ and $4J_z$ for $\alpha \rightarrow \infty$ (see Supplementary Information). A nonlocal light cone is clearly visible in the long-range regime ($\alpha = 0.1, 0.5, 1$) and disappears in the nearest-neighbor coupling limit, $\alpha \rightarrow \infty$, (f). (g) Outside-the-cone profiles at $t = 0.7$, obtained by locating the crest of the outward-propagating wave and plotting the decay of $\|\Delta\rho_n(t)\|_F$ beyond the wave crest. All profiles decay exponentially (consistent with the Lieb–Robinson bound), with the long-range cases exhibiting an even steeper decay than the nearest-neighbor case whose profile is fitted by an exponential (orange dashed line with $A = 1.3$, $b = 1.5$). (h) Time-averaged ($t \in [0, 0.5]$) signals at sites 19 and 10 showing respectively the intraband and interband signals as a function of interaction exponent α . The dashed curves show the predicted analytical small- α scaling in Eqs. (9) and (10), up to an overall constant fitting factor.

at other positions (e.g., another observer, Charlie), Alice can broadcast the nonlocal signal to multiple arbitrarily chosen destinations simultaneously, [see, e.g., Fig. 3(a–c)]. Crucially, this mechanism extends to higher dimensions, see Supplementary Information.

The model— We demonstrate this physics using 1D long-range interacting spin-1/2 systems with open boundary conditions, Fig. 1(b), described by the Hamiltonian:

$$\hat{H} = \sum_{j=1}^{N-1} J_z \sigma_j^z \sigma_{j+1}^z + \sum_{i < j} \frac{J_x}{|j-i|^\alpha} \sigma_i^x \sigma_j^x, \quad (1)$$

where $J_z, J_x > 0$ (setting $J_z = 1$ in the following for convenience), and α controls the interaction range, interpolating between infinite-range ($\alpha = 0$) and nearest-neighbor ($\alpha \rightarrow \infty$) limits. Note that applying Kac rescaling [32–34] for thermodynamic consistency does not affect our main results (see Supplementary Information). To quantify the information spread, we use the Frobenius-norm distance between the reduced density matrices $\rho_n^{(1)}(t)$ and $\rho_n^{(2)}(t)$, respectively, of two evolved initial states at each site n :

$$\|\Delta\rho_n(t)\|_F = \sqrt{\text{Tr}(\Delta\rho_n(t)\Delta\rho_n^\dagger(t))}, \quad (2)$$

where $\Delta\rho_n(t) = \rho_n^{(1)}(t) - \rho_n^{(2)}(t)$. We normalize this value in the interval $[0, 1]$, see Methods. Unlike two-point correlators, OTOCs, or magnetization, which test specific observables in specific orientations, the Frobenius norm

in Eq. (2) captures the maximum possible change in any measurement outcome, making it sensitive to coherent differences that other probes can miss.

Information propagation can also be evaluated using the site-resolved x -magnetization $\frac{|\langle\psi_1|\sigma_n^x|\psi_1\rangle - \langle\psi_2|\sigma_n^x|\psi_2\rangle|}{2}$ or the two-point correlator $|\langle\sigma_2^x(t)\sigma_j^x(t)\rangle - \langle\sigma_2^x(t)\rangle\langle\sigma_j^x(t)\rangle|$, as has been done in ion trap experiments [5]. In the latter case, only a single initial state is required. Both figures of merit can be used to unveil the same physics described by the Frobenius norm, see Supplementary Information. Specifically, to compute the normalized Frobenius norm, we compare the time evolution of two nearby initial states, $|\psi_1\rangle$ and $|\psi_2\rangle$. As an illustrative example, consider a chain of $N = 20$ spins and

$$\begin{aligned} |\psi_1\rangle &= |\downarrow_1 \uparrow_2 \downarrow_3 \downarrow_4 \cdots \uparrow_{19} \downarrow_{20}\rangle, \\ |\psi_2\rangle &= |\downarrow_1 \downarrow_2 \uparrow_3 \downarrow_4 \cdots \uparrow_{19} \downarrow_{20}\rangle, \end{aligned} \quad (3)$$

which differ by a single-site shift of one spin-up excitation (site $2 \leftrightarrow 3$, with the second excitation fixed at site 19). All spins are specified in the σ^x basis, i.e., along the direction defined by the long-range interaction. We denote such pairs of initial states compactly as

$$(i_1, i_2; j_1, j_2), \quad (4)$$

where (i_1, i_2) and (j_1, j_2) specify the excitation positions in $|\psi_1\rangle$ and $|\psi_2\rangle$, respectively; here (2, 19; 3, 19). Physically, this setup corresponds to a differential signaling protocol: Alice encodes information by choosing

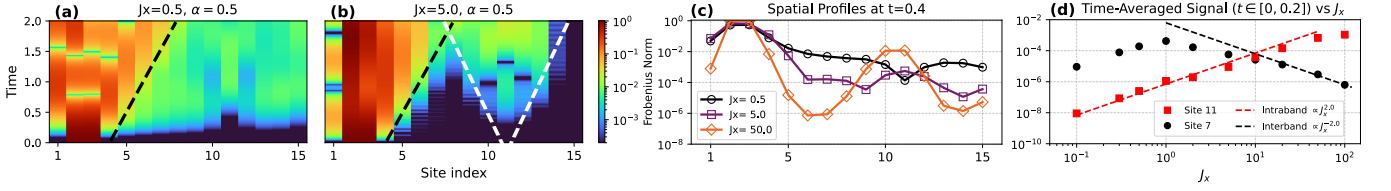


FIG. 2. **Emergence of nonlocal light-cones.** (a,b) Density plots of the Frobenius norm for $N = 15$, $\alpha = 0.5$ showing light-cone dynamics for initial states prepared in band 2, (2, 11; 3, 11) (notation defined in Eq. (4)). Dashed lines (black and white for contrast) indicate the analytically predicted light-cone velocities (see Supplementary Information). (c) Spatial profiles of the light cone at a fixed time $t = 0.4$. Increasing J_x suppresses interband contributions, leading to the clear emergence of a well-defined nonlocal light cone. (d) Scaling behavior of the time-averaged signal. At site 11 (red squares), the time-averaged signal scales as J_x^2 , in conformance with the analytical intraband scaling formula, Eq. (10). Conversely, at site 7 (black circles), the signal exhibits an asymptotic J_x^{-2} scaling in the large-bandgap limit, see Eq. (9) and Supplementary Information for more details.

between two different local excitations at the source (site 2 vs. 3). Bob retrieves this information by monitoring how the system's evolution (state $|\psi_2\rangle$) deviates from a known calibration baseline (state $|\psi_1\rangle$).

An example of nonlocal light cones emerging from the interplay of long- and short-range interactions for strong long-range coupling J_x is shown in Fig. 1(c-e). We compute the Frobenius norm by evolving the states in Eq. (3) with the Hamiltonian in Eq. (1) for different values of the interaction exponent α . For long-range interactions (c-e), a second light cone erupts at the target site 19, despite the two initial states being identical in that region. As a comparison, for the same initial conditions with nearest-neighbor interactions ($\alpha = \infty$), a single light cone emerges from the source perturbation, Fig. 1(f). The information leaking outside the nonlocal, receiver-centered light cone (site 19) decays faster than exponentially, matches the behavior observed for the light cone originating at the source, see Fig. 1(g) and Supplementary Information for more details. The intensity of the signal both within the nonlocal light cone and in the intermediate region between the source and receiver depends on several parameters. In particular, their dependence on the interaction exponent α is shown in Fig. 1(h). As α increases, both the signal inside the nonlocal light cone and between the two light cones increase at first, while for $\alpha \gg 1$ both signals decrease and only the light cone at the source remains.

Band structure and mapping to a hard-core boson model— The emergence of these programmable nonlocal light cones can be rigorously understood by analyzing the band structure of the excitation spectrum and by mapping the spin Hamiltonian to a hard-core boson chain, as detailed below.

Let us consider the band structure of the spectrum of long-range interacting systems. For $\alpha = 0$, the long-range term depends only on the total x -magnetization $M_x = \frac{1}{2} \sum_{j=1}^N \sigma_j^x$, so the spectrum organizes into well-separated degenerate “bands” (or excitation manifolds) labeled by the number $b = 0, 1, 2, \dots$ of up spins along x . The band energies E_b are quadratic in b with large

interband gaps ΔE_b that strongly suppress transitions between them. Explicitly, one has:

$$E_b = \frac{J_x}{2}(N - 2b)^2 - \frac{J_x N}{2}, \quad (5)$$

$$\Delta E_b = E_b - E_{b-1} = -2J_x(N - 2b + 1). \quad (6)$$

The J_z nearest-neighbor term couples only bands $b \leftrightarrow b \pm 2$.

For $0 < \alpha < 1$, the many-body bands remain, but intraband degeneracies lift. For band 1, the band width δE_1 is set by the energy difference between excitations at the chain center and edge, which for small α behaves as

$$\frac{\delta E_1}{|\Delta E_1|} \sim \frac{2J_x(\ln 2)\alpha N^{1-\alpha}}{2J_x N^{1-\alpha}} = (\ln 2)\alpha. \quad (7)$$

As $\alpha \rightarrow 1^-$, the ratio becomes $\mathcal{O}(1)$; for $\alpha > 1$ the band width exceeds the gap, causing spectral overlap and interband mixing (see Supplementary Information). Since higher bands broaden by an additional factor $\sim b$, the overlap criterion for band 1 already determines the onset of overlap for the full spectrum.

This band structure suggests a mapping in terms of hard-core bosons hopping on a lattice of size N , where a boson corresponds to spin up in the σ^x basis. Explicitly, Eq. (1) takes the equivalent form of a bosonic Hamiltonian (see Methods for derivations):

$$\begin{aligned} \hat{H}_{\text{eff}} = & J_z \sum_i \left(\hat{a}_i^\dagger \hat{a}_{i-1} + \hat{a}_{i-1}^\dagger \hat{a}_i \right) + J_z \sum_i \left(\hat{a}_i^\dagger \hat{a}_{i-1}^\dagger + \hat{a}_i \hat{a}_{i-1} \right) \\ & + \sum_i V(i) \hat{n}_i + \sum_{i < j} W_{i,j} \hat{n}_i \hat{n}_j + U \sum_i \hat{n}_i (\hat{n}_i - 1) + E_{b=0}, \end{aligned} \quad (8)$$

where

$$\begin{aligned} V(i) &= -2J_x \sum_{j \neq i} \frac{1}{|i - j|^\alpha}, & W_{i,j} &= \frac{4J_x}{|i - j|^\alpha}, \\ U \rightarrow \infty \text{ (hard-core)}, & & E_{b=0} &= J_x \sum_{i < j} \frac{1}{|i - j|^\alpha}. \end{aligned}$$

Here $V(i)$ acts as a local potential and $W_{i,j}$ represents the instantaneous density-density interaction between distant bosons. In this representation, the band number becomes the total boson occupation number $\sum_i \hat{n}_i$, and the second term in $\hat{H}_{\text{eff}}$ contains the coupling between bands.

Manifestly, the only source of nonlocality in $\hat{H}_{\text{eff}}$ is the two-body interaction $W_{i,j} \propto J_x/|i-j|^\alpha$. Indeed, if the W term is removed from the Hamiltonian, the dynamics exhibit only local information spread (see Supplementary Information).

Interband and intraband nonlocality and the emergence of nonlocal light cones—The long-range interaction term (W) drives two distinct nonlocal mechanisms: a programmable *intraband* process operating strictly within a fixed excitation manifold, and an unprogrammable *interband* process driven by virtual transitions between different manifolds.

The strengths of both mechanisms can be estimated by noting that the normalized Frobenius norm is exactly equal to the absolute difference in local boson densities: $|\langle \hat{n}_n^{(1)}(t) \rangle - \langle \hat{n}_n^{(2)}(t) \rangle|$ for two different initial states with a fixed boson number, see Supplementary Information. The *interband* nonlocality is generated by the second term in the r.h.s. of Eq. (8), which creates or annihilates a pair of bosons, giving rise to a transition from band b to band $b \pm 2$. The probability to create or annihilate a boson pair will depend, through the W term, on the distribution of bosons in the initial band- b state. In the presence of large band gaps, $\Delta E \sim J_x N^{1-\alpha} \gg J_z$, the interband signal can be modeled via first-order time-dependent perturbation theory, yielding an interband signal strength that scales as (see Supplementary Information):

$$\|\Delta\rho_n(t)\|_F^{\text{interband}} \sim \frac{\alpha J_z^2}{J_x^2 N^{3(1-\alpha)} r^{\alpha+1}}, \quad (9)$$

where r is the distance from the source. Eq. (9), which is valid for $\alpha < 1$, establishes that the interband background vanishes entirely in the thermodynamic limit ($N \rightarrow \infty$) or for strong long-range coupling ($J_x \rightarrow \infty$). The case $\alpha = 0$ is special and is discussed in detail in the Supplementary Information.

In the limit of small interband nonlocality, information spread can be described by the band-projected effective Hamiltonian H_{eff}^b , obtained from Eq. (8) by omitting the second term. In the projected Hamiltonian, which serves the total boson number, the interplay of the kinetic hopping term proportional to J_z and the density-density interaction $W_{i,j}$ determines the emergence of nonlocal light cones at the boson positions. Physically, when a local perturbation shifts the source boson, it instantly modulates the interaction potential felt by all remote background bosons. This sudden quench of the local energy landscape triggers a response: the remote bosons begin to propagate via the nearest-neighbor hopping term J_z , gener-

erating independent, localized light cones that erupt directly from their respective positions. This *intraband* signal can be estimated by computing the time-evolved density operator $\hat{n}_q(t) = U^\dagger(t) \hat{n}_q U(t)$, $U(t) = e^{-iH_{\text{eff}}^b t}$. Expanding the evolution operator to fourth order in time, one obtains:

$$\|\Delta\rho_n(t)\|_F^{\text{intraband}} = \frac{t^4 J_z^2 (4J_x)^2}{3} \alpha^2 (\alpha + 1) \frac{1}{r_0^{2\alpha+3}} + \mathcal{O}(t^5), \quad (10)$$

independent of both the system size N and the specific band index b .

Fig. 2 demonstrates how increasing the long-range interaction strength J_x gives rise to a non-local light cone. At small J_x , information spreads unconstrained across the entire lattice [Fig. 2(a)], but as J_x increases, a distinct non-local light cone originates at the target excitation [Fig. 2(b)]. This trend is confirmed by the fixed-time slice in Fig. 2(c): increasing J_x suppresses the interband signal and amplifies the intraband signal, consistent with the scaling of Eqs. (9) and (10) shown in Fig. 2(d).

To characterize the signal at nonlocal light cones, we study $\hat{H}_{\text{eff}}^b$. In Fig. 3(a-c), the Frobenius norm is evaluated for two initial states in $b = 3$ that differ by a displacement of the first excitation from site 2 to site 3, with the positions of the other two excitations held fixed. Fig. 3(a-c) clearly shows how the causal spacetime structure can be precisely shaped: nonlocal light cones erupt precisely at the positions of the fixed excitations and propagate independently, each exhibiting local behavior as seen in their respective profiles [Fig. 3(d-f)]. In order to study how the intraband signal (inside nonlocal light cones) decays with distance, we consider a pair of initial states in band $b = 2$ and vary the target excitation position. The intraband signal decays with distance with an exponent predicted by Eq. (10), $p(\alpha) = 2\alpha + 3$, see Fig. 3(g,h). Also, the scaling with other parameters predicted by Eqs. (9, 10) is confirmed numerically: the α -dependence for $\alpha < 1$ is verified in Fig. 1(h) and the scaling with J_x is verified in Fig. 2(d). For the scaling with N and more details, see Supplementary Information.

Crucially, the emergence of nonlocal light cones requires a spatially dependent potential; for infinite-range interactions ($\alpha = 0$), W reduces in each band to an overall spectral shift, and no nonlocal light cones emerge, see Supplementary Information. Also note that the nearest-neighbor interaction J_z is vital to this architecture; in the absence of this term, or upon substituting it with a local external field, nonlocal (and local) light cones disappear entirely, see Supplementary Information.

The competition between the intraband and interband signals determines the controllability of the causal landscape. By evaluating the intraband signal at the characteristic transport time $t \sim 1/J_z$, we identify an upper bound for the critical distance r_c below which the causal landscape is controllable. Demanding that the intraband signal overcomes the static interband background,

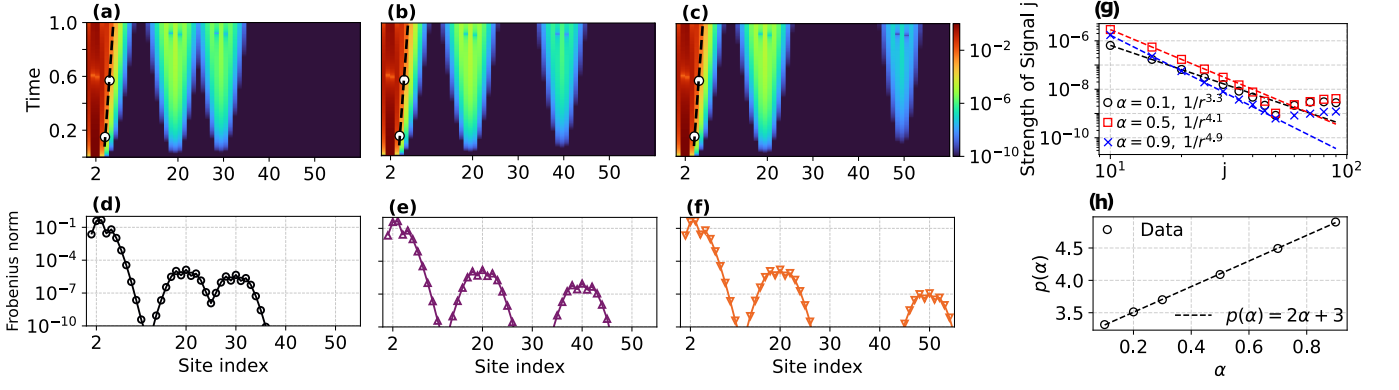


FIG. 3. **Shaping the space-time causal landscape.** Density plots of the Frobenius norm showing light-cone dynamics under the projected Hamiltonian $\hat{H}_{\text{eff}}^{b=3}$, obtained by varying the target position j for initial states prepared in band 3, $(2, 20, j; 3, 20, j)$ (notation defined in Eq. (4)), for $j = 30$ (a), $j = 40$ (b), and $j = 50$ (c). Other parameters are: $J_x = 0.5$, $\alpha = 0.5$, and $N = 100$ (only 60 sites are shown). (d,e,f) The respective light-cone profiles at $t = 1$. The amplitude of the nonlocal light cone at j decreases as j increases, while the nonlocal cone originating at site 20 remains unaffected by the variation of j . In panels (g,h), $\hat{H}_{\text{eff}}^{b=2}$ is used to study the decay of the intraband signal with j using the initial states $(2, j; 3, j)$ with $N = 100$, $J_z = 1$, and $J_x = 0.1$. (g) Power-law scaling of the intraband signal in the nonlocal light cone. Summing the signal over bulk sites (sites 10–90) at $t = 0.5$ and varying the excitation position j yields a power-law scaling $\sim 1/r^{p(\alpha)}$, where $r = j - 3$ is the distance between message and target excitations. Note that when the distance exceeds $N/2$, the signal increases due to boundary conditions (see Supplementary Information). (h) Extracted exponent p as a function of α , showing the relation $p(\alpha) = 3 + 2\alpha$.

$F_{\text{intraband}}(1/J_z) > F_{\text{interband}}$, leads to the condition, for $\alpha < 1$:

$$r_0^{\alpha+2} < \left(\frac{J_x}{J_z}\right)^4 N^{3(1-\alpha)}. \quad (11)$$

This boundary defines the critical radius r_c . In the thermodynamic limit ($N \rightarrow \infty$), $r_c \rightarrow \infty$; the unprogrammable background noise is perfectly suppressed by the large energy gap $\Delta E \sim J_x N^{1-\alpha}$, making the entire chain controllable. In the Supplementary Information, a more detailed derivation is shown together with the critical radius in the presence of Kac rescaling.

Emergence of locality— Numerous studies have shown that the presence of long-range interactions does not guarantee enhanced information propagation; frequently, dynamics remain confined within a linear light cone for timescales that grow with system size [12, 27–31, 35–40]. Some aspects of this phenomenon, also termed *cooperative shielding*, have been confirmed experimentally in trapped ion systems [36]. Nevertheless a unified microscopic understanding of when and why locality emerges from inherently nonlocal Hamiltonians has remained elusive.

Our analysis clarifies that the emergence of local versus nonlocal behavior in long-range interacting systems is dictated directly by the initial conditions, a result that generalizes previous findings [41, 42].

As derived above, the limit $N^{1-\alpha} J_x \gg J_z$ suppresses interband leakage. Dynamics are then dictated by the Hamiltonian projected onto manifolds of conserved excitation number, leaving the intraband density-density

interaction $W_{i,j}$ as the unique source of nonlocality. Crucially, the role of $W_{i,j}$ depends on how these excitations are spatially distributed.

If the initial state consists of excitations spatially confined to a specific region, the nonlocal W term vanishes outside this region. Consequently, information spread remains strictly local in the bulk, and nonlocal light cones emerge *only* at the specific positions of the pre-existing excitations. Conversely, if the initial condition contains excitations distributed across the entire chain (e.g., a global quench from a z -polarized state), the W term connects all sites simultaneously, leading to highly nonlocal propagation, see Supplementary Information for a detailed comparison of states initialized in different bases. This framework elegantly reconciles seemingly contradictory experimental results.

Conclusions and discussion— We have demonstrated that nonlocality in long-range interacting systems can be precisely engineered, allowing the causal space-time landscape of many-body quantum spin systems to be programmably shaped. By exploiting the interplay between long-range and short-range interactions, we established a protocol to induce nonlocal light cones at deterministic locations, effectively allowing information to bypass the geometric bulk of the system.

Our findings are immediately actionable in state-of-the-art quantum simulators. Trapped ion platforms are particularly well-suited to verify these predictions: while they natively realize long-range interactions, the specific Hamiltonian required here can be engineered using established digital or hybrid quantum simulation protocols [2, 6, 7]. Furthermore, recent advances in multimode cavity

QED have enabled tunable-range and sign-changing interactions [22, 23, 43], offering an alternative pathway to realize the connectivity required for our protocol. Ultimately, engineered nonlocal light cones elevate connectivity to a programmable resource. By shaping where and how information propagates, our protocol enables targeted long-range communication without sacrificing local isolation. This opens new directions for quantum communication, state transfer, simulation, error correction, and distributed quantum computing, while providing a general framework for designing many-body systems with tailored information-propagation landscapes.

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METHODS

Conventions— Throughout this work, we use a normalized Frobenius norm, defined by dividing $\|\Delta\rho_n(t)\|_F$ by its global maximum over all sites and times, thus mapping all signals to the interval $[0, 1]$. For brevity, we refer to this normalized quantity simply as the Frobenius norm.

Numerical methods— Real-time dynamics are computed by direct integration of the time-dependent Schrödinger equation using a fourth-order Runge-Kutta (RK4) scheme. Time evolution is performed either in the full many-body Hilbert space or within a band-projected manifold, obtained by restricting the Hamiltonian to a suitably chosen manifold with a fixed excitation number. The time step Δt is chosen sufficiently small to ensure numerical convergence of all reported observables. Convergence is assessed by systematically reducing Δt until relative changes in local observables, Frobenius norms, and extracted light-cone velocities between successive runs fall below numerical precision. Typical simulations employ $\Delta t \in [10^{-7}, 10^{-5}]$, depending on system size and interaction strength. While the RK4 scheme is not exactly unitary and accumulates errors at long times, all results reported in the main text correspond to evolution times $t \lesssim 10$, over which norm conservation is explicitly monitored at each time step, with the maximum deviation satisfying $|\langle\psi(t)|\psi(t)\rangle - 1| < 10^{-14}$.

Exact diagonalization is employed to compute spectral quantities and to benchmark RK4 simulations. Initial states are chosen as product states with a localized perturbation in the σ_x basis, prepared within a fixed excitation sector unless stated otherwise. All numerical

results are cross-validated using independent diagnostics, yielding consistent behavior across different system sizes (further details can be found in the Supplementary Information).

Formal derivation: Holstein–Primakoff transformation and linear spin wave theory— We derive the effective Hamiltonian for the spin model of Eq. (1) using the Holstein–Primakoff (HP) transformation. Crucially, because spin-1/2 systems strictly forbid double occupancy, the higher-order corrections in the Holstein–Primakoff transformation identically vanish. Consequently, this mapping to a hard-core boson model is not restricted to the low-excitation regime, but is fully exact across all excitation manifolds.

The standard HP transformation assumes a fully +z polarized ferromagnetic vacuum, and the HP mapping reads:

$$\hat{S}_z = S - \hat{n}, \quad (12)$$

$$\hat{S}_+ = \sqrt{2S} \sqrt{1 - \frac{\hat{n}}{2S}} \hat{a}, \quad (13)$$

$$\hat{S}_- = \sqrt{2S} \hat{a}^\dagger \sqrt{1 - \frac{\hat{n}}{2S}}, \quad (14)$$

where S denotes the total spin quantum number of the local magnetic moments, $\hat{n} = \hat{a}^\dagger \hat{a}$ is the local boson number operator, and the ladder operators are defined as $\hat{S}_\pm = \hat{S}_x \pm i\hat{S}_y = \frac{1}{2}(\sigma_x \pm i\sigma_y)$. In the low-excitation (dilute) regime, where $\hat{n} \ll 2S$, the square root can be expanded as a power series:

$$\sqrt{1 - \frac{\hat{n}}{2S}} = 1 - \frac{\hat{n}}{4S} - \frac{\hat{n}^2}{32S^2} - \dots \quad (15)$$

In the zeroth-order (linearized) approximation of the HP transformation, which keeps only the leading behaviour in $\langle \hat{a}^\dagger \hat{a} \rangle$, the spin operators become $\hat{S}_+^{(0)} = \sqrt{2S} \hat{a}$, $\hat{S}_-^{(0)} = \sqrt{2S} \hat{a}^\dagger$, $\hat{S}_z^{(0)} = S - \hat{a}^\dagger \hat{a}$, and so $\sigma_x \approx \hat{a} + \hat{a}^\dagger$, $\sigma_y \approx i(\hat{a}^\dagger - \hat{a})$, $\sigma_z = 1 - 2\hat{a}^\dagger \hat{a}$.

However, our model is defined with a ground state polarized along the x -axis, so we apply a $\pi/2$ rotation about the y -axis, $\hat{R}_y(\frac{\pi}{2}) = e^{-i\frac{\pi}{4}\sigma_y}$. Substituting the linearized HP expressions into the rotated frame gives $\sigma_z \approx -(\hat{a} + \hat{a}^\dagger)$, $\sigma_y \approx i(\hat{a}^\dagger - \hat{a})$, $\sigma_x = 1 - 2\hat{n}$. Note that the terms involving σ_x (i.e., the J_x terms) are exact. The hopping terms, on the other hand, are the zeroth-order contributions in the expansion shown in Eq. (15). Applying the rotated bosonic operators to the Hamiltonian of Eq. (1), we obtain:

$$\begin{aligned} \hat{H}^{(0)} = & J_z \sum_j \left(\hat{a}_j^\dagger \hat{a}_{j+1} + \hat{a}_{j+1}^\dagger \hat{a}_j + \hat{a}_j^\dagger \hat{a}_{j+1}^\dagger + \hat{a}_j \hat{a}_{j+1} \right) \\ & - 2 \sum_{i < j} \frac{J_x}{|i - j|^\alpha} (\hat{n}_i + \hat{n}_j) + 4 \sum_{i < j} \frac{J_x}{|i - j|^\alpha} \hat{n}_i \hat{n}_j \\ & + \sum_{i < j} \frac{J_x}{|i - j|^\alpha}. \end{aligned} \quad (16)$$

Beyond linear order, the HP expansion introduces density-dependent hopping terms $\sqrt{1 - \frac{\hat{n}}{2S}} \approx 1 - \frac{\hat{n}}{4S}$, leading to interaction-assisted hopping, which we denote $\delta\hat{H}^{(1)}$. So

$$\hat{H} = \hat{H}^{(0)} + \delta\hat{H}^{(1)}, \quad (17)$$

where $\hat{H}^{(0)}$ is given by Eq. (16), and the additional terms at the next order are:

$$\begin{aligned} \delta\hat{H}^{(1)} = & J_z \sum_{j=1}^{L-1} \left[-\frac{1}{4} (\hat{a}_j + \hat{a}_j^\dagger) (\hat{n}_{j+1} \hat{a}_{j+1} + \hat{a}_{j+1}^\dagger \hat{n}_{j+1}) \right. \\ & - \frac{1}{4} (\hat{n}_j \hat{a}_j + \hat{a}_j^\dagger \hat{n}_j) (\hat{a}_{j+1} + \hat{a}_{j+1}^\dagger) \\ & \left. + \frac{1}{16} (\hat{n}_j \hat{a}_j + \hat{a}_j^\dagger \hat{n}_j) (\hat{n}_{j+1} \hat{a}_{j+1} + \hat{a}_{j+1}^\dagger \hat{n}_{j+1}) \right]. \end{aligned}$$

Since in spin- $\frac{1}{2}$ ($S = 1/2$) systems double occupancy is forbidden, every term in $\delta\hat{H}^{(1)}$ necessarily attempts either over-annihilation or creation on an already occupied site, and therefore all such contributions vanish identically in every band. Thus we recover the effective Hamiltonian, Eq. (8).

Subspace projection derivation of the effective Hamiltonian— We begin with the single-excitation band ($b = 1$). The nearest-neighbor Ising term induces hopping, giving

$$\hat{H}_{\text{eff}}^{b=1} = J_z \sum_j (|j\rangle\langle j+1| + \text{h.c.}) - 2J_x \sum_j \epsilon_j |j\rangle\langle j|. \quad (18)$$

Thus the $b = 1$ band maps to a local tight-binding model with position-dependent onsite potential.

Higher bands— A basis state in band b has b flipped spins. Let us define the set of positions D with spins flipped, e.g., $D = \{3, 4, 7\}$ for $b = 3$ if spins in position 3, 4, 7 are flipped. The contribution to the energy from the J_x interaction can be decomposed as

$$\begin{aligned}
E &= J_x \sum_{i < j} \frac{\sigma_i^x \sigma_j^x}{(j-i)^\alpha} = J_x \left[\sum_{\substack{i < j \\ i, j \notin D}} \frac{1}{(j-i)^\alpha} - \sum_{\substack{i < j \\ i \in D, j \notin D}} \frac{1}{(j-i)^\alpha} - \sum_{\substack{i < j \\ i \notin D, j \in D}} \frac{1}{(j-i)^\alpha} + \sum_{\substack{i < j \\ i, j \in D}} \frac{1}{(j-i)^\alpha} \right] \\
&= J_x \left[\sum_{i < j} \frac{1}{(j-i)^\alpha} - 2 \sum_{j \in D} \sum_{i \neq j} \frac{1}{|j-i|^\alpha} + 4 \sum_{\substack{i, j \in D \\ i < j}} \frac{1}{|j-i|^\alpha} \right] = E_{b=0} + \sum_{j \in D} V(j) + \sum_{\substack{i, j \in D \\ i < j}} W_{i,j}, \tag{19}
\end{aligned}$$

where

$$E_{b=0} = J_x \sum_{i < j} \frac{1}{(j-i)^\alpha}, \quad V(j) = -2J_x \sum_{i \neq j} \frac{1}{|j-i|^\alpha},$$

and

$$W_{i,j} = \frac{4J_x}{|j-i|^\alpha}.$$

Thus, the projected Hamiltonian for band b contains a local onsite potential V and a nonlocal two-body term W .

It is interesting to note that in the case of Kac rescaling (explicitly addressed in the Supplementary Information), for which $J_x \sim J_{\text{long}}/N^{1-\alpha}$, the one-body potential scales as $O(\alpha J_{\text{long}})$ for small α , while the two-body term is at most $O(b^2 J_{\text{long}}/N^{1-\alpha})$. Hence for $\alpha < 1$, fixed b , and $N \rightarrow \infty$, W becomes negligible and the effective Hamiltonian approaches locality.

Second-quantized form— Introducing creation and annihilation operators $\hat{a}_i^\dagger, \hat{a}_i$ for up spins in the σ_x basis, one finds

$$\hat{V} = \sum_i V(i) \hat{a}_i^\dagger \hat{a}_i + \sum_{i < j} W_{i,j} \hat{a}_i^\dagger \hat{a}_i \hat{a}_j^\dagger \hat{a}_j, \tag{20}$$

with a hard-core constraint forbidding double occupancy. The Ising term provides the kinetic energy

$$\hat{T} = J_z \sum_i \left(\hat{a}_i^\dagger \hat{a}_{i-1} + \hat{a}_{i-1}^\dagger \hat{a}_i \right). \tag{21}$$

Collecting terms gives the effective Hamiltonian (8), a hard-core Bose-Hubbard model with one-body and two-body potentials. All contributions are local except for W .

Data availability: The data that support the findings of this study are available from the corresponding author upon reasonable request.

Code availability: The code that supports the findings of this study is available from the corresponding author upon reasonable request.

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