

Distributing entanglement at the quantum speed limit in Rydberg chains

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We numerically study the transport of Rydberg excitations in chains of neutral atoms. We realize an effective flip-flop interaction using off-resonant driving fields. By tuning the relative distances between atoms and applying atom-selective detuning fields, we fulfill the perfect transport condition. This condition enables the transfer of a single Rydberg excitation from one end of the chain to the other, allowing the distribution of entanglement across the chain at the quantum speed limit. Through numerical simulations, we identify the set of control parameters that maximizes the transport probability under experimentally relevant conditions. We study the various competing trade-offs involved in the hierarchy of approximations used to map the native Rydberg spin model onto the effective model driving spin transport. Our results suggest that entanglement can be distributed over chains of more than 50 atoms spanning hundreds of microns at room temperature. This study informs the selection of parameters for the experimental realization of perfect transport in Rydberg chains, providing an approach to distribute entanglement among distant atoms in quantum processors.

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I. INTRODUCTION

The efficient distribution of quantum information between distant quantum registers is a fundamental challenge in the development of scalable quantum information processors [1]. A typical approach relies on quantum channels capable of transmitting quantum information with high fidelity [2]. In quantum processors with static architectures, where registers are fixed in space and their interactions decay according to a power law, linear chains of spins have been proposed as a quantum channel for enabling short-distance quantum communication [3,4]. By precisely tuning the interaction strengths to achieve the perfect transport condition, the quantum state of a spin can be perfectly transferred to another spin located at a mirror-symmetric location in the chain at the quantum speed limit (QSL) [5–7]. This coherent state transfer enables faster information transport than sequential gate-based protocols and serves as a primitive for entanglement distribution between quantum registers, which in turn enables quantum teleportation protocols [8].

While the theoretical framework for perfect state transfer is well established, its experimental realization has been limited to one-dimensional transport over short distances [9,10]. Experimental realization is difficult due to limitations imposed by decoherence, disorder, and the difficulty of precisely engineering interactions in solid-state quantum systems

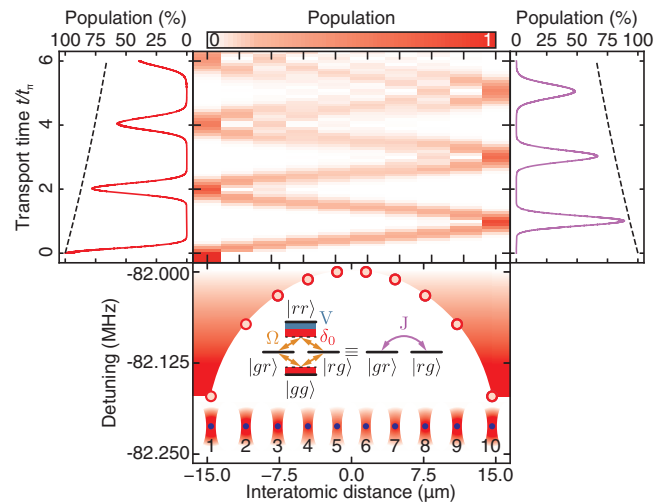


FIG. 1. Perfect transport in a Rydberg chain. (Top) A single-spin excitation prepared at one end of the chain is coherently transported to the other end at time t_π . The probability of measuring the single-spin excitation on the first (left) and last (right) spins undergoes collapses and revivals, which are damped by radiative decay (dashed line) and the loss of coherence caused by approximation errors and long-range interactions. (Bottom) The interatomic distances and local detunings are tuned to realize the perfect transport condition. (Inset) The effective spin exchange between two atoms is realized by detuning the driving field away from the single-spin resonance. The spin exchange is mediated through two-photon Raman transitions between the zero-excitation manifold and the two-excitation manifold, whose energy is shifted due to the Rydberg blockade. The large detuning suppresses the injection of spin excitations into high-excitation manifolds.

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[11,12]. Recent advances in the control of large, interacting quantum many-body systems offer an imminent opportunity for practical realization in larger, two-dimensional systems [13–17]. A particularly promising platform is offered by configurations of neutral atoms excited to a Rydberg state, also known as Rydberg atom arrays [18–21]. By realizing quantum spin models with variable parameters on arbitrary geometries with over a thousand atoms, these quantum simulators act as analog twins of solid-state spin systems, enabling the validation and cross-benchmarking of various control protocols across different parameter regimes.

Rydberg atom arrays support several routes to transport quantum states and distribute entanglement, including atom displacement [22], sequences of Rydberg-mediated swap gates [23,24], facilitated hopping via modulation of interatomic couplings [25], and perfect state transfer (PST) studied in this work. The key advantage of PST is that it only requires time-independent control on static trap arrays, which is particularly relevant for systems based on microlens arrays without active diffractive elements [26,27]. For transporting a Bell state over $60\text{ }\mu\text{m}$ with at least 50% fidelity, we estimate characteristic speeds of $0.3\text{ }\mu\text{m }\mu\text{s}^{-1}$ for swap-gate chains, $0.75\text{ }\mu\text{m }\mu\text{s}^{-1}$ for shuttling, $3\text{ }\mu\text{m }\mu\text{s}^{-1}$ for PST, and $51\text{ }\mu\text{m }\mu\text{s}^{-1}$ for facilitated hopping. Details of the methods used for these estimates are provided in the Supplemental Material [28]. Under the same target and geometry, PST is about four times faster than shuttling (and an order of magnitude faster than swap-gate chains) while requiring only static control on a static trap array. This balance of speed and simplicity highlights PST as a practical candidate for entanglement distribution in Rydberg arrays.

In this paper, we investigate the conditions required to achieve perfect transport in chains of neutral atoms with Rydberg interactions. Our goal is to identify the set of control parameters that maximizes the probability of perfect transport under realistic experimental conditions, as well as to understand how this probability scales with physical parameters such as the principal quantum number, temperature, and chain length. Through extensive numerical simulations, we identify and quantify the key trade-offs involved in optimizing transport probability, including the competition between transport speed and spin-excitation injection, as well as the interplay between longer lifetimes, stronger interactions, and reduced driving rates at higher principal quantum numbers. Our approach relies on mapping the transverse-field Ising spin Hamiltonian realized by Rydberg chains onto an effective Heisenberg XX Hamiltonian [29,30], and engineering the perfect transport condition by tuning the interatomic spacing and applying atom-selective detuning fields.

Beyond finding the set of parameters that maximizes the transport probability, we aim to set a baseline for the experimental study of spin transport on Rydberg platforms and to quantify their robustness against experimental imperfections. For example, the transport probability may be used as a metric to certify the realization of Heisenberg models in Rydberg systems. We also seek to identify the key challenges involved in relying on spin transport to distribute entanglement between distant registers in quantum processors.

The outline of the paper is as follows. We first review the conditions required to realize perfect transport in Rydberg

chains and discuss the hierarchy of approximations involved in mapping the theoretical model onto the experimentally relevant Rydberg model (see Sec. II). We then present our numerical simulation results identifying the set of control parameters that maximizes the transport probability under realistic experimental conditions (see Sec. III). These results highlight some of the trade-offs involved when optimizing over the detuning and principal quantum number. We finally explore the maximum chain length over which the system can serve as a coherent quantum channel for entanglement distribution (see Sec. IV). These findings lay the groundwork for implementing quantum state transfer protocols in programmable Rydberg platforms.

II. PERFECT TRANSPORT IN RYDBERG CHAINS

We consider a quantum system formed by a finite chain of L rubidium-87 atoms with variable, possibly nonuniform, interatomic spacings. Each atom encodes a qubit between its ground state, $|g\rangle = |5S_{1/2}\rangle$, and excited Rydberg state, $|r\rangle = |nS_{1/2}\rangle$. The principal quantum number is taken to be $50 \leq n \leq 80$, following typical tweezer experiments [31,32]. Each qubit is coherently driven via a two-photon Raman transition mediated through an intermediate electronic state, $|e\rangle$. The detuning of the optical field driving the $|g\rangle$ to $|e\rangle$ transition, δ_{ge} , is chosen to be sufficiently large to allow for the adiabatic elimination of $|e\rangle$ in the dynamics. For simplicity, we refer to the dressed states of the effective two-level qubit as $|g\rangle$ and $|r\rangle$. The effective Rabi frequency between $|g\rangle$ and $|r\rangle$ is $\Omega = (\Omega_{ge}\Omega_{er})/(2\delta_{ge})$, where Ω_{ge} and Ω_{er} are the Rabi frequencies for the transitions from $|g\rangle$ to $|e\rangle$ and from $|e\rangle$ to $|r\rangle$, respectively. The residual admixture of $|e\rangle$ in $|r\rangle$ is $|\epsilon_r|^2 \approx \Omega_{er}^2/4\delta_{ge}^2$, so that the effective decay rate of $|r\rangle$ is $\Gamma'_r(n, \theta) \approx \Gamma_r(n, \theta) + |\epsilon_r|^2\Gamma_e$, where Γ_e is the decay rate of the intermediate excited state and $\Gamma_r(n, \theta) = \Gamma_r^0(n) + \Gamma_r^{\text{BBR}}(n, \theta)$ is the state-dependent decay rate of the bare Rydberg state, which depends on the ambient temperature, θ , due to coupling to the blackbody radiation (BBR) [33]. For n ranging from 50 to 80, $\Gamma_r(n, 300\text{ K})$ spans 3.7–16.1 kHz and $\Gamma_r(n, 4\text{ K})$ spans 1.2–7.8 kHz, with larger values occurring at smaller n . Similarly, the decay rate of the dressed ground state $|g\rangle$ is $\Gamma_g \approx |\epsilon_g|^2\Gamma_e$, with $|\epsilon_g|^2 \approx \Omega_{ge}^2/4\delta_{ge}^2$; however, we neglect this effective decay rate in our calculations.

The Hamiltonian describing the dynamics of the system is analogous to the transverse-field Ising model [19],

$$\hat{\mathcal{H}}_{\text{Ryd}} = \sum_{l=1}^L \left(\frac{\Omega_l}{2} \hat{\sigma}_x^l - (\delta_0 + \delta_l) \hat{n}_l \right) + \sum_{l < l'} V_{ll'} \hat{n}_l \hat{n}_{l'}, \quad (1)$$

where $\hat{\sigma}_\alpha^l$ is the α -component Pauli operator for the l th spin, $\hat{n}_l = |r_l\rangle\langle r_l| = (\hat{1} - \hat{\sigma}_z^l)/2$ is the projection operator onto $|r_l\rangle$, and $V_{ll'} = V(n, \Delta x_{ll'}) = C_6(n)/\Delta x_{ll'}^6 > 0$ is the repulsive van der Waals interaction strength between two Rydberg atoms located at x_l and $x_{l'}$ with interatomic spacing $\Delta x_{ll'} = x_l - x_{l'}$. We define the global detuning as $\delta_0 = \delta_{gr}$ and include an additional site-selective detuning, δ_l , arising from second-order light shifts induced by far-off-resonant, atom-selective lasers [34]. For the sake of completeness, we choose the Rabi frequency Ω_l to be tunable on each site, although our protocol does not require it.

The native spin Hamiltonian does not drive spin transport when the state of the chain is initialized in an eigenstate of the $\bigotimes_j \hat{\sigma}_z^j$ operator [35,36]. However, an effective flip-flop interaction can be realized through second-order perturbation theory [29,30]. In the limit of large detunings, $\Omega_l \ll |\delta_0|$, $|\delta_0 - V_{ll'}|$ and $|\delta_l| \ll |\delta_0|$, Eq. (1) maps onto the effective XX Hamiltonian [29],

$$\hat{\mathcal{H}}_{XX} = \sum_{l=1}^L \mu_l \hat{n}_l + \sum_{l < l'} J_{ll'} (\hat{\sigma}_+^l \hat{\sigma}_-^{l'} + \hat{\sigma}_-^l \hat{\sigma}_+^{l'}), \quad (2)$$

where $\hat{\sigma}_\pm^l = (\hat{\sigma}_x^l \pm i\hat{\sigma}_y^l)/2$ are the raising and lowering operators for the l th spin, μ_l is the effective on-site potential, and $J_{ll'}$ is the effective flip-flop rate. This Hamiltonian is the Heisenberg XX model in a transverse field, which generates spin transport while conserving the number of spin excitations ($U(1)$ symmetry). An analytical expression for the parameters of the effective model, $\{\mu_l, J_{ll'}\}$, expressed in terms of the set of control parameters, $\{x_l, \Omega_l, \Delta_l = \delta_0 + \delta_l\}$, can be obtained from second-order perturbation theory [29],

$$\mu_l = -\Delta_l - \frac{\Omega_l^2}{2\Delta_l} + \sum_{l' \neq l} \frac{\Omega_l^2 V_{ll'}}{4\Delta_{l'}(\Delta_{l'} - V_{ll'})}, \quad (3)$$

$$J_{ll'} = \frac{\Omega_l \Omega_{l'} V_{ll'}}{8\Delta_l(\Delta_l - V_{ll'})} + \frac{\Omega_l \Omega_{l'} V_{ll'}}{8\Delta_{l'}(\Delta_{l'} - V_{ll'})}. \quad (4)$$

The effective Hamiltonian from Eq. (2) can be further simplified by restricting the sum over the dominant nearest-neighbor (NN) interactions,

$$\hat{\mathcal{H}}_{XX}^{\text{NN}} = \sum_{l=1}^L \mu_l \hat{n}_l + \sum_{l=1}^{L-1} J_{l,l+1} (\hat{\sigma}_+^l \hat{\sigma}_-^{l+1} + \hat{\sigma}_-^l \hat{\sigma}_+^{l+1}). \quad (5)$$

This Hamiltonian is isomorphic to the hard-core limit of the Bose-Hubbard Hamiltonian via the Holstein-Primakoff transformation [37], where the spin excitations play the role of hard-core bosons with at most one particle per site. By tuning the effective on-site potential, $\{\mu_l\}$, it is thus possible to use the Rydberg chain to study the evolution of bosons in a tilted optical lattice, leading to observables such as Bloch oscillations [14,38]. This Hamiltonian also maps onto a system of free fermions using the Jordan-Wigner transformation [39], enabling an analytical treatment of transport dynamics through the dispersion relation for finite-length chains with open boundary conditions [40].

A. Perfect transport in Heisenberg chains

The Heisenberg XX Hamiltonian generates ballistic transport of single-spin excitations on chains. For equidistant chains with uniform couplings, a single-spin excitation prepared at one end of the chain is not perfectly transported to the other end, unless the effective control parameters are tuned to realize the perfect transport condition [5]. Mathematically, the probability of the system evolving from $|l\rangle$ to its mirror-symmetric counterpart $|L+1-l\rangle$ after a time t is quantified by the transport probability,

$$p(t) = | \langle L+1-l | \hat{U}(t) | l \rangle |^2, \quad (6)$$

where $\hat{U}(t) = \exp(-i\hat{\mathcal{H}}t/\hbar)$ is the evolution operator generated by $\hat{\mathcal{H}}$ and $|l\rangle = |g_1 g_2 \cdots r_l \cdots g_L\rangle$ is the many-body

spin state hosting a single-spin excitation on the l th site for $l \geq 1$. We denote the many-body ground state containing no excitations as $|0\rangle = |g_1 g_2 \cdots g_L\rangle$.

Perfect state transfer is realized at a time t_π if there exists a set of control parameters such that $p_\pi = p(t_\pi) = 1$. Multiple choices of parameters may lead to perfect transport [5,6,41], including for systems with power-law decaying interactions [4]. These solutions distinguish themselves by their perfect transport time, which is lower bounded by the QSL [7]. For the Heisenberg XX model with NN interactions, the linear spectrum solution [5,6] is obtained for

$$\mu_l = \mu, \quad (7)$$

$$J_{l,l+1} = 2J_{\text{max}} \sqrt{l(L-l)}/\bar{L}, \quad (8)$$

where

$$J_{\text{max}} = \frac{\Omega^2}{4\delta_0} \left(\frac{\delta_0 \Delta x_{\text{min}}^6}{C_6} - 1 \right)^{-1} \quad (9)$$

is the maximum coupling strength between the central atoms separated by Δx_{min} and $\bar{L}$ is the characteristic length of the chain, which is defined as $\bar{L} = L$ for even chains and $\bar{L} = \sqrt{L^2 - 1}$ for odd chains. This solution can be understood as the necessary condition to achieve coherent population inversion in a collective spin system [5,42]. It saturates the QSL [7] with perfect transport time,

$$t_\pi(L) = \frac{\pi \bar{L}}{4J_{\text{max}}}. \quad (10)$$

B. Solving the inversion problem for Rydberg chains

In Rydberg chains, finding the set of control parameters that realizes the perfect transport condition for the linear spectrum solution involves solving an inversion problem. This problem requires finding the set of control parameters $\{\Omega_l, \Delta_l, V_{ll'}\}$ that satisfy Eqs. (3) and (4) under the constraints imposed by Eqs. (7) and (8). However, this system of equations is underdetermined, featuring $2L - 1$ equations and $3L - 1$ variables, thereby admitting multiple solutions.

A unique solution to the perfect transport problem is obtained for uniformly spaced chains by solving for the restricted set of parameters $\{\Delta_l, \Omega_l\}$ [29]. Here, we choose to solve for $\{\Delta_l, V_{ll'}\}$ instead, as it is typically easier to adjust the relative positions of neutral atoms using optical traps than to apply atom-selective driving fields. We set the interatomic distance between the central spins to the minimum achievable value, Δx_{min} , thereby fixing the interaction strength at the center of the chain to $V_{\text{max}} = C_6(n)/\Delta x_{\text{min}}^6$. We also set the on-site detuning of the central spins to zero, $\delta_l = 0$, so that $\Delta_l = \delta_0$. The problem is further simplified by exploiting the mirror symmetry of the perfect transport condition, which imposes $\Delta_l = \Delta_{L+1-l}$ and $x_l = -x_{L+1-l}$. Details of the numerical methods used to solve this inversion problem are provided in the Supplemental Material [28].

A representative numerical solution for a chain of $L = 10$ atoms excited to the $n = 70$ Rydberg state is shown in Fig. 1. Both the atom-selective detunings δ_l and the interatomic spacing $\Delta x_{ll'}$ increase away from the center of the chain, ensuring that the on-site potential μ_l remains constant, while satisfying the linear spectrum condition for the $J_{ll'}$ coefficients. The

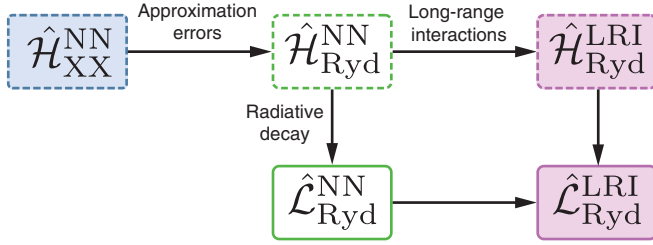


FIG. 2. Hierarchy of approximations. The perfect transport condition is realized for the Heisenberg XX Hamiltonian with nearest-neighbor (NN) interactions in the absence of radiative decay (blue). Its mapping onto the experimentally relevant Rydberg model incurs errors from the perturbative approximation (green), long-range interactions (LRIs; purple), and radiative decay (solid line).

coherent collapses and revivals of the single-spin excitation at the ends of the chain are damped by radiative decay and the loss of coherence resulting from approximation errors and long-range interactions. We now seek to quantify the competing influence of these effects under realistic experimental conditions.

III. OPTIMIZED CONTROL PARAMETERS AND EXPERIMENTAL TRADE-OFFS

The analytical solution for perfect transport is obtained when the chain evolves under $\hat{\mathcal{H}}_{XX}^{NN}$ in the absence of decay, control errors, and other imperfections. In practice, however, the system evolves under the full dissipative Rydberg Liouvillian $\hat{\mathcal{L}}_{Ryd}^{LRI}$, which includes both radiative decay and long-range interactions (LRIs). Realizing the perfect transport condition in a chain of Rydberg atoms requires mapping $\hat{\mathcal{L}}_{Ryd}^{LRI}$ onto $\hat{\mathcal{H}}_{XX}^{NN}$. This mapping involves a hierarchy of approximations (see Fig. 2), whose validity correlates with the probability of perfect transport, p_π . Maximizing p_π involves a trade-off between competing physical effects, which we quantify using numerical simulations.

A. Simulation parameters

We compute p_π for different sets of physical and control parameters under realistic experimental conditions. We set the minimum interatomic distance to $\Delta x_{\min} = 3 \mu\text{m}$ [19]. We further set the maximum optical power available to drive the $|e\rangle$ to $|r\rangle$ transition to $P_{er} = 4 \text{ W}$, which corresponds to a peak intensity of $I = 127 \text{ kW/cm}^2$ when delivered using a Bessel beam with 10% of its power contained within a core radius of $w_0 = 10 \mu\text{m}$. These constraints effectively set an upper bound for $J_{\max}(n)$, and thus a lower bound for $t_\pi \sim J_{\max}^{-1}(n)$.

Fixing n , we compute the interaction coefficients, $C_6(n)$, and dipole matrix elements, $d_{er}(n)$, using the ARC package [43]. We use these values to calculate the effective Rabi frequency, $\Omega(n)$, assuming $\Omega_{ge} = \Omega_{er}$ for a given maximum achievable laser power P_{er} . Then fixing L and δ_0 , we solve the inversion problem described in Sec. II B to find the set of local detuning values, δ_l , and atom positions, x_l , that realize the perfect transport condition under $\hat{\mathcal{H}}_{XX}^{NN}$. The global detuning is taken to be $-120 \text{ MHz} \leq \delta_0 \leq -40 \text{ MHz}$, which

is achievable with current hardware [44]. Using a reduced subspace method for all chains with $L > 4$ spins, we compute the time evolution of the initial state $|1\rangle$ generated by $\hat{\mathcal{H}}_{Ryd}^{LRI}$, $\hat{\mathcal{H}}_{Ryd}^{NN}$, and $\hat{\mathcal{H}}_{XX}^{NN}$ using QuTiP [45]. From the time-evolved state, $|\psi(t)\rangle = \hat{U}(\delta_0, t)|1\rangle$, we compute the transport probability, $p_\pi(\delta_0) = |\langle L|\psi(t_\pi)\rangle|^2$, as well as the population leaked into the many-body ground state containing no Rydberg excitation, $p_\pi^{(0)} = |\langle 0|\psi(t_\pi)\rangle|^2$, and the higher-excitation subspaces containing more than one excitation, $p_\pi^{(2+)} = 1 - (p_\pi^{(0)} + p_\pi^{(1)})$, where $p_\pi^{(1)} = \sum_{l=1}^L |\langle l|\psi(t_\pi)\rangle|^2$. These calculations are repeated for different values of δ_0 to find the optimal detuning, $\delta_0^* = \text{argmax}_{\delta_0} p_\pi(\delta_0)$, that maximizes the transport probability, $p_\pi^* = p_\pi(\delta_0^*)$. We multiply these values by $\exp(-\Gamma_r(n, \theta)t_\pi)$ to account for the radiative decay from the Rydberg state. This full procedure is then applied for various n and L to quantify the dependence of $p_\pi^*(n, L)$ on the principal quantum number, n , and chain length, L .

B. Optimizing over the global detuning

The first trade-off involves a competition between the speedup of transport dynamics and the breakdown of second-order perturbation theory when sweeping the global detuning, δ_0 . As $|\delta_0|$ decreases toward resonance, the effective interaction coefficients increase, leading to a speedup of transport, $t_\pi \sim J_{\max}^{-1}$ [see Fig. 3(a)]. This faster dynamics leads to a reduction in the probability of decay to the zero-excitation ground state, at the cost of an increase in the number of spin excitations being injected into the chain [see the crossover between $p_\pi^{(0)}$ and $p_\pi^{(2+)}$ in Fig. 3(b)]. As δ_0 is detuned away from resonance, the transport dynamics slows down, increasing the probability of decay to the zero-excitation ground state. We therefore find that for a given set of parameters, there exists an optimal detuning value, δ_0^* , for which the transport probability is maximized [Fig. 3(c)].

The effect of these competing factors can be quantified by comparing the transport probability obtained from different models over varying values of the global detuning δ_0 [Fig. 3(c)]. As the detuning is decreased toward resonance, the gradual breakdown of second-order perturbation theory is observed by the drop in p_π when comparing the evolution generated by $\hat{\mathcal{H}}_{XX}^{NN}$ (dashed blue) and $\hat{\mathcal{H}}_{Ryd}^{NN}$ (dash-dot green). As the detuning is increased away from resonance, the slowdown of transport dynamics causes a loss of probability between the exact model generated by $\hat{\mathcal{H}}_{XX}^{NN}$ (dashed blue) and its decaying version generated by $\hat{\mathcal{L}}_{XX}^{NN}$ (dotted blue). The competition between radiative decay and the injection of spin excitations results in a convex curve for $\hat{\mathcal{L}}_{Ryd}^{NN}$ (long-dash green) and $\hat{\mathcal{L}}_{Ryd}^{LRI}$ (solid purple), each of which exhibits a single maximum value at δ_0^* .

Given that the effective Rabi frequency Ω_l is bounded by the maximum available laser power (see Sec. III A), this variation of δ_0 is equivalent to tuning the perturbative parameter $\xi = |\Omega_l/\delta_0|$. For the $L = 16$ example in Fig. 3, the detuning sweep covers $\xi \in [0.075, 0.20]$. The optimal performance occurs at $\xi_{\text{opt}} \approx 0.089$, reflecting the balance between the perturbation error, which mainly manifests as

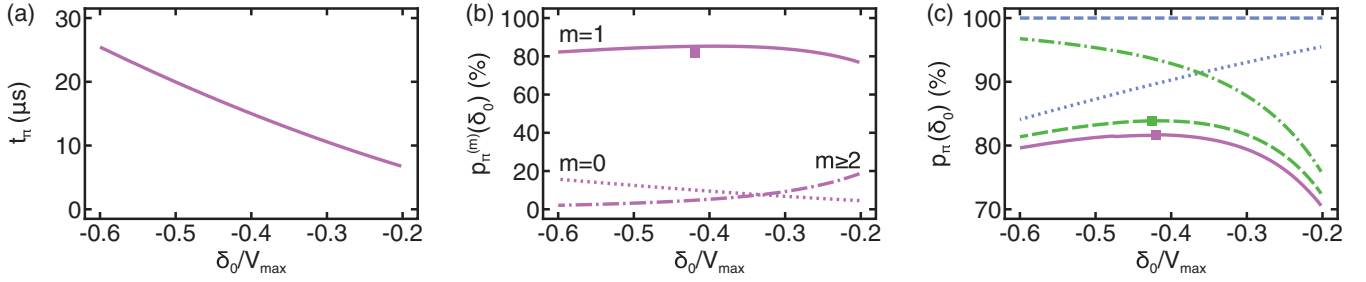


FIG. 3. Optimizing over the global detuning. (a) The transport time for the Rydberg model, which approximately overlaps for both NN interactions and long-range interactions, decreases as the detuning approaches resonance. (b) Fractional populations in the excitation subspaces containing $m = 0$ (dot), $m = 1$ (solid), and $m \geq 2$ (dash-dot) spin excitations computed for $\hat{\mathcal{L}}_{\text{Ryd}}^{\text{LRI}}$ at t_π . The maximum transport probability (purple square) achieved is less than the total population in the $m = 1$ subspace due to imperfect refocusing. (c) Transport probabilities computed in the absence and presence of radiative decay for the Heisenberg XX model with NN interactions (dashed blue and dotted blue, respectively) and the Rydberg model with NN interactions (dash-dot green and long-dash green, respectively); long-range interactions (LRIs; solid purple) are plotted only with radiative decay. The maximum probabilities (squares) are achieved at distinct optimal values of δ_0^* . All data are computed for $L = 16$, $n = 70$, and $V_{\text{max}} = 188.3$ MHz for $\Delta x_{\text{min}} = 3$ μm .

population leakage into the $m \geq 2$ excitation manifolds, and state decay, which mainly manifests as leakage into the zero-excitation ground state. Across all simulated chain lengths, this optimal operating point falls consistently within $\xi \in [0.088, 0.104]$.

C. Optimizing over the principal quantum number

The second trade-off involves a competition between longer lifetimes, stronger interactions, and weaker driving strengths when accessing larger principal quantum numbers. As n increases, the van der Waals interaction strength increases as $V_{\text{max}}(n) \sim C_6(n)/\Delta x_{\text{min}}^6 \sim n^{11}$, and the radiative lifetime increases as $\tau(n) \sim n^3$, while the effective Rabi frequency decreases as $\Omega(n) \sim d_{\text{er}}(n) \sim n^{-3/2}$ when the maximum accessible optical power is fixed.

We numerically compute the optimal transport probability $p_\pi^*(n)$ obtained for different n [see Fig. 4(a)]. The maximum transport probability monotonically increases with n for both the evolution generated by $\hat{\mathcal{L}}_{XX}^{\text{NN}}$ (dotted blue line) and $\hat{\mathcal{L}}_{\text{Ryd}}^{\text{NN}}$ (long-dash green line). This monotonic increase is caused by longer lifetimes at higher n , despite the decrease in transport speed caused by lower Rabi frequencies. The scaling of the transport speed with n can be understood by examining the effective interaction strength $J_{\text{max}}(n)$, whose amplitude can be expressed as the ratio of two terms,

$$|J_{\text{max}}(n)| = \frac{\Omega(n)^2/4|\delta_0(n)|}{|\delta_0(n)|/V_{\text{max}}(n) + 1}, \quad (11)$$

where $\delta_0(n) = -|\delta_0(n)| < 0$. For small values of n , $|\delta_0| \gg V_{\text{max}}$, such that J_{max} is dominated by the denominator [see Fig. 4(b)]. For larger values of n , $|\delta_0| \ll V_{\text{max}}$, such that the denominator goes to one, resulting in $|J_{\text{max}}(n)| \sim \Omega^2(n) \sim n^{-3}$. Therefore, there exists a critical value n_c at which J_{max} reaches a maximum, after which it decays according to a power law. Despite slower transport at large n , the corresponding increase in lifetime, $\tau(n) \sim n^3$, leads to a flattening of the transport probability, reflecting its asymptotic behavior. This plateau

is reflected in the ratio between the transport time and the Rydberg state lifetime [see Fig. 4(c)].

The monotonic increase in transport probability is interrupted by introducing long-range interactions to the Rydberg

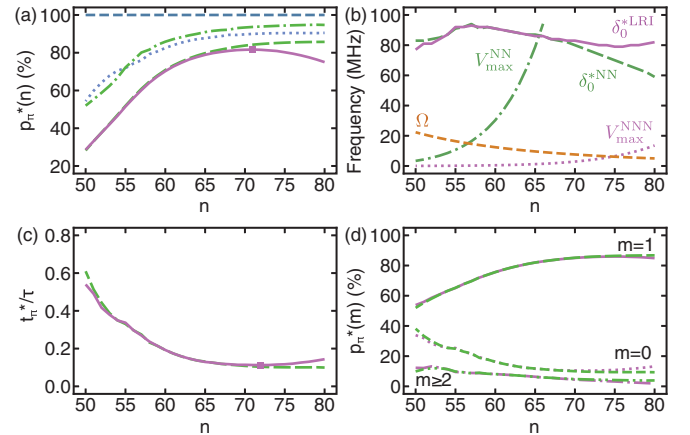


FIG. 4. Optimizing over the principal quantum number. (a) Maximum transport probability, $p_\pi^*(n)$, achieved for each principal quantum number n in the absence and presence of radiative decay for the Heisenberg XX model (dashed blue and dotted blue, respectively) and the Rydberg model with NN interactions (dash-dot green and long-dash green, respectively); long-range interactions (solid purple) are plotted only with radiative decay. The Rydberg LRI model achieves its maximum transport probability at $n^* = 71$ (purple square). (b) Maximum interaction strength, $V_{\text{max}}(n)$, between NN (dash-dot green) and next-nearest-neighbor (NNN; dotted purple) atoms. These values are compared against the optimal detuning values, $|\delta_0^*(n)| = -\delta_0^*(n)$, for the Rydberg NN (long-dash green) and LRI (solid purple) models, and the Rabi frequency (dashed orange), which is constrained by the maximum achievable optical power. (c) Rescaled transport time $t_\pi^*(n)/\tau(n)$, where $\tau(n)$ is the radiative lifetime of the Rydberg state, for the NN (long-dash green) and LRI (solid purple) models. (d) Fractional populations in the subspaces containing $m = 0, 1$, and $(m \geq 2)$ Rydberg excitations for the Rydberg NN (dashed green, long-dash green, dash-dot-dot green, respectively) and LRI (dotted purple, solid purple, dash-dot purple, respectively) models. All data are reported for $L = 16$.

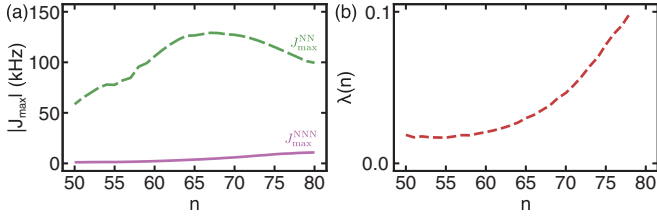


FIG. 5. Relative contribution of NNN interactions. (a) Maximum effective interaction strength, $|J_{\max}(n)|$, between NN (long-dash green) and NNN (solid purple) interactions for the Rydberg model. (b) The ratio $\lambda(n) = J_{\max}^{\text{NNN}}(n)/J_{\max}^{\text{NN}}(n)$ increases with n due to the growing relative contribution of $J_{\max}^{\text{NNN}}$ at larger n .

model, i.e., when considering the transport generated by $\hat{\mathcal{L}}_{\text{Ryd}}^{\text{LRI}}$ [see solid purple line in Fig. 4(a)]. In this case, the transport probability attains its maximum at $n^* = 71$ (purple square) for $L = 16$. The relative loss in transport probability beyond n^* is dominated by the existence of next-nearest-neighbor (NNN) interactions, which modify the energy spectrum and eigenstates of the chain [41]. While the dynamics remains restricted to the $m = 1$ excitation subspace [see Fig. 4(d)], the NNN interactions generate coherences among distant spins, hindering the constructive refocusing of the spin excitation at the target site. The imperfect refocusing also produces a time delay, such that maximal refocusing occurs slightly later than the QSL. This delay increases monotonically with chain length and reaches its largest value in this work at $L = 20$, where the transport is 1.2% slower than the QSL bound. For the n considered in this work, interactions beyond NNN contribute negligibly to the dynamics.

The loss in transport probability in the LRI model due to NNN interactions can be explained by the competition between the maximum effective interaction strengths $J_{\max}$ achieved for NN and NNN interactions [see Fig. 5(a)]. While $J_{\max}^{\text{NN}}$ reaches a local maximum at some critical n_c near the point where $V_{\max}(n)$ overtakes $|\delta_0^*(n)|$, $J_{\max}^{\text{NNN}}$ monotonically increases. Their ratio, which is given by

$$\lambda(n) = \frac{J_{\max}^{\text{NNN}}(n)}{J_{\max}^{\text{NN}}(n)} \quad (12)$$

$$\approx \frac{|\delta_0^{\text{NN}}(n)|/V_{\max}^{\text{NN}}(n) + 1}{|\delta_0^{\text{LRI}}(n)|/V_{\max}^{\text{NNN}}(n) + 1}, \quad (13)$$

is approximately constant until n_c , after which it starts increasing [see Fig. 5(b)]. This growth reflects the growing influence of NNN interactions relative to NN interactions, which hinders the realization of perfect transport. At the cost of having to solve an inverse eigenvalue problem, this loss in transport probability may be suppressed by including NNN interactions when solving the inversion problem, effectively choosing a different solution from the linear-spectrum one [41]. However, in this case the peak transport probability becomes narrowly localized in time and highly sensitive to small parameter variations, making it difficult to reproduce experimentally. For this reason, we focus on the linear spectrum solution, which provides high-fidelity transport with reduced sensitivity to experimental imperfections.

IV. QUANTUM CHANNEL FOR ENTANGLEMENT DISTRIBUTION

By enabling the transport of single-spin excitations, Rydberg chains can be used as coherent quantum channels to distribute entanglement between distant quantum nodes [3,29]. The entanglement distribution protocol involves adding an auxiliary spin at the beginning of the chain, entangling it with the first spin of the chain, decoupling it from the dynamics by nulling its Rabi frequency, storing its coherence in some long-lived clock state, and then letting the chain propagate the state for t_π until it reaches the other end of the chain [3,29]. This protocol results in, e.g., the Bell state $\Phi_+^{0,1} = (|0, 0\rangle + |1, 1\rangle)/\sqrt{2}$ being transported to $\Phi_\phi^{0,L} = (|0, 0\rangle + e^{i\phi(L)}|1, L\rangle)/\sqrt{2}$, where the length-dependent phase factor can be accounted for by rotating the local measurement basis of the end spin. The entanglement between the auxiliary and end spins can be detected using an entanglement witness [46–48], as well as quantified using an entanglement measure, for which we choose the concurrence [49], $C_\pi^{0,L} = \max(0, 2p_\pi(L) - 1)$. The concurrence for the single-spin transport protocol depends on the transport probability, and it vanishes when $p_\pi(L) \leq 1/2$. A nonzero concurrence indicates that the shared state contains distillable entanglement, enabling the implementation of teleportation protocols [8].

Because $p_\pi(L)$ decreases with L due to radiative decay over longer transport times, there exists a critical chain length, L_c , beyond which the Rydberg chain can no longer distribute quantum entanglement. Beyond this cutoff, the two-qubit joint state behaves effectively as a classical mixed state, rendering the quantum channel ineffective. To estimate L_c , we numerically compute the set of parameters that achieves the optimal transport probability $p_\pi^*(L)$ for different chain lengths. Because Rydberg state decay depends on temperature, we compute L_c for both ambient conditions (300 K) and cryogenic temperatures (4 K), where coupling to blackbody radiation is strongly suppressed for cryogenic temperatures [50].

As the chain length is increased, we find that the optimal principal quantum number n^* has a slowly increasing piecewise-continuous pattern [see Fig. 6(a)], while the optimal detuning values remain relatively constant [see Fig. 6(b)]. The stability of the control parameters results in a relatively stable $J_{\max}(n)$, such that the transport time grows approximately linearly with L for chains up to our largest simulated chain length $L = 20$ [see Fig. 6(c)]. This increase in transport time results in a decrease in $p_\pi^*(L) \sim \exp(-t_\pi^*(L)/\tau(n^*))$ [see Fig. 6(d)], which decays approximately exponentially. By extrapolating our simulation data to larger chain lengths, we find that the maximum chain length for distributing entanglement is $L_c = 54$ at 300 K and $L_c = 90$ at 4 K. Reducing the temperature thus provides an additional practical benefit beyond increasing vacuum lifetime and reducing Rydberg decay rates.

While population transfer is a key figure of merit for spin transport, the retention of coherence is crucial for entanglement transfer. The dominant sources of decoherence for entanglement transfer are dephasing due to atomic Doppler shifts, laser phase noise, and stray electric fields, as well as disorder effects arising from the finite precision of

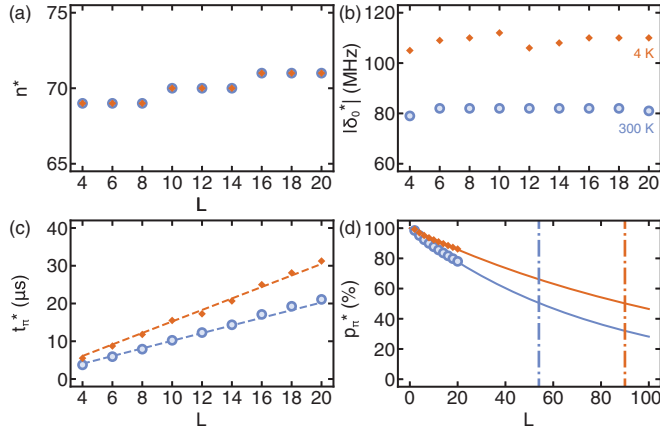


FIG. 6. Scaling with chain length. (a), (b) Optimal principal quantum number n^* and detuning $|\delta_0^*|$ for the Rydberg model with long-range interactions and decay, computed as a function of the chain length L at temperatures of 300 K (blue disks) and 4 K (orange diamonds). (c) The optimal transport time increases approximately linearly with L . (d) The transport probability decreases approximately exponentially (solid line) with L , with a smaller decay rate at 4 K due to reduced Rydberg state decay. The maximum chain lengths over which entanglement can be distributed are estimated to be $L_c = 54$ and $L_c = 90$ at 300 K and 4 K, respectively.

tweezer placement, the finite temperature of atoms in optical traps, and local addressing noise [51–55]. First, we consider dephasing effects, which we compute using Lindbladian evolution. With a sufficiently small dephasing rate γ , spin transport remains ballistic over the full perfect state transfer duration, where the characteristic crossover time from ballistic to diffusive transport occurs at $t_c = 1/\gamma$ [29]. A typical crossover time is $t_c \approx 50 \mu\text{s}$, considering a dephasing rate of $\gamma = 0.02 \text{ MHz}$ reported for alkali atoms in the absence of sideband cooling to the motional ground state, where technical noise contributed significantly to the observed value [52]. With $\gamma = 0.02 \text{ MHz}$, we compute less than 8.5% coherence loss for $L \leq 20$.

Next, we consider the effects of disorder using Monte Carlo simulations sampled over 100 iterations. We first evaluate positional disorder. Positional uncertainty manifests as disorder in the van der Waals interaction term of the Rydberg Hamiltonian, which in turn affects both the local potential and the exchange interaction in the effective picture. The figure of merit $\delta J_{II'}/J_{II'} \sim \delta V_{II'}/V_{II'} \ll 1$ quantifies robustness against positional disorder [29]. We evaluate this by combining thermal effects with the finite resolution of tweezer placement. Using a thermal wave packet standard deviation computed from experimentally measured atom-in-trap temperatures after Raman sideband cooling [53,54], alongside a tweezer placement precision of 65 nm [55], we incorporate the total spatial variance into the interactions. This positional disorder results in up to an 18% loss in transport probability for $L \leq 20$. Even in the motional ground state, the finite spatial extent of the atomic wave packet leads to motional dephasing, which is further exacerbated by time-dependent effects due to repulsive interactions. We have studied these effects in an accompanying paper, in which we find that mo-

tional degrees of freedom can lead to a reduction in L_c by half under suboptimal conditions [56]. The dephasing depends on trap parameters (e.g., trapping frequency) and control parameters (e.g., detuning and interatomic distance), offering an opportunity to achieve higher transport fidelity by optimizing over a larger parameter space that couples internal and external degrees of freedom.

Finally, we evaluate the effects of detuning disorder. Because the global detuning is implemented by a cavity-stabilized narrow-linewidth laser, the dominant detuning noise arises from intensity fluctuations in the local addressing beams, which have an error of less than 2% [34]. Applying this relative local addressing noise results in less than a 0.5% loss in transport probability for $L \leq 20$.

V. CONCLUSION

In conclusion, we have shown that Rydberg chains can act as a coherent quantum channel to distribute quantum information near the QSL, between distant atoms separated by hundreds of microns. These chains provide an alternative approach to increasing the connectivity of quantum processors, complementing existing methods based on mediating quantum gates, shared bosonic modes, photonic links, and the displacement of quantum particles.

Through numerical simulations, we have shown that finding the set of optimal control parameters for realizing the perfect transport condition in Rydberg chains requires balancing competing trade-offs imposed by approximation errors, physical errors, and control limitations. These trade-offs set an upper bound on the ultimate performance achievable by the spin chain as a coherent quantum channel. In practical settings, this performance will be further degraded by spatial inhomogeneities and temporal instabilities in both control and physical parameters, e.g., resulting from spatial disorder and motional dephasing. We estimate that these effects are not prohibitive, with further room for improvement. When considering further scaling of system size, these effects can be quantified using the transport probability, providing a figure of merit to quantify detrimental effects and optimize the performance of robust control protocols.

While our approach relies on mapping the native Rydberg spin model onto a Heisenberg XX model using second-order perturbation theory, our analysis framework can readily be applied to quantify the performance of alternative protocols based on resonant dipole-dipole interactions [57–60], exchange between Rydberg and Rydberg-dressed ground states [61,62], time-dependent driving [25], and second- and fourth-order exchange of ground hyperfine states coupled to Rydberg states [63–66]. This framework could also be extended to topological quantum phenomena capable of distributing quantum information, such as anomalous chiral edge flow in the Kitaev honeycomb model [67,68]. Beyond its experimental realization in Rydberg chains, this framework could be extended to arbitrary spin networks with multiple Rydberg excitations to explore transport and scattering phenomena, as well as applications such as quantum state routing.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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