

Robust Rydberg facilitation via rapid adiabatic passage

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Rydberg antiblockade, or facilitation, governs interaction-assisted excitation dynamics that generate a wide variety of collective phenomena and correlated many-body behavior, while also forming the basis of protocols for entangling gates, state transfer, and quantum sensing. Despite its rich physics, experimental realization of facilitation has been fragile because the resonance condition requires precise matching between interaction shifts and laser detuning, making the process highly sensitive to atomic position disorder and parameter fluctuations. To overcome this limitation, we develop a robust implementation of Rydberg facilitation based on rapid adiabatic passage. By adiabatically sweeping the laser detuning across the interaction-shifted resonance, our approach is inherently insensitive to typical experimental imperfections. As a direct demonstration, we show that this adiabatically protected facilitation naturally gives rise to avalanche excitation growth in both one- and two-dimensional arrays. This avalanche process yields high gain with exceptionally low background, making it promising for rare-event detection. More broadly, by completely eliminating disorder sensitivity, our approach transforms Rydberg facilitation from an experimentally elusive regime into a robust mechanism implementable with standard cooling techniques and existing optical tweezer platforms, bridging the gap between theoretical proposals and disorder-free experimental realizations of antiblockade-based quantum control.

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I. INTRODUCTION

The ability to engineer strong and controllable interactions among neutral atoms has made Rydberg arrays a leading platform for quantum information processing and quantum simulation [1–3]. Central to this interaction engineering is the Rydberg blockade, where strong interactions prevent simultaneous excitation of nearby atoms, enabling high-fidelity multiqubit control. Complementary to this blockade regime, Rydberg antiblockade, or facilitation, has been proposed as a route to faster quantum logic gates free from blockade errors and to implement swap operations without the overhead of decomposing them into multiple elementary gates [4–6]. Moreover, antiblockade enables entanglement generation [7], provides avalanche amplification for sensing weak signals [8], and constitutes a pivotal ingredient for kinetically constrained lattice models that test the limits of our understanding of quantum thermalization and quantum matter [9–14].

Despite its potential to unlock quantum-device functionalities and many-body physics regimes unattainable with Rydberg blockade alone, experimental realization of antiblockade remains fragile because it requires precise control of the interaction strength and laser detuning. Finite atomic temperature results in a distribution of atomic positions, lead-

ing to variations in interatomic separations, which translates into disorder in the interactions due to their strong distance dependence. Such disorder is generally inconsequential in the blockade regime, although exceptions exist [15]. In contrast, antiblockade relies on a finely tuned resonance condition in which the interaction shift is compensated by the laser detuning [16,17]. Small deviations in interatomic spacing or drive parameters therefore destroy the facilitation resonance, making experimental operation challenging.

Consequently, high-fidelity antiblockade protocols demand advanced cooling techniques beyond the capability of typical tweezer-array experiments. Even with extreme cooling techniques and ultraprecise controls, the achievable fidelities often remain noncompetitive due to residual position disorder. For example, a recently proposed high-tolerance antiblockade swap gate [4] predicts a fidelity of 0.955, whereas a compiled swap gate on a neutral-atom platform, constructed from experimentally demonstrated [18,19] *cz* gates (0.995) and single-qubit rotations (0.9997), is expected to reach a higher fidelity (0.982), despite its greater circuit overhead. The same sensitivity to position disorder that limits the performance of antiblockade logic operations also constrains its use for many-body quantum simulation. Rydberg antiblockade is a compelling route to realizing the transverse-field Ising model in atom-array systems. However, position disorder transforms the expected ballistic propagation of correlations into sub-ballistic dynamics and changes the growth of entanglement entropy from linear to logarithmic in time, preventing faithful realization of the target Hamiltonian [20].

Stimulated Raman adiabatic passage has been proposed to implement robust unconditional excitations via antiblockade [7,21]. Here, we propose to realize robust facilitated

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excitations, where existing Rydberg atoms catalyze further excitations, via rapid adiabatic passage (RAP). We compare with a resonant Rabi protocol to showcase its robustness to variations in driving parameters and position disorders. Although adiabaticity is typically associated with longer evolution times, our study shows that under realistic constraints this is not necessarily the case. As an example application, we show that this mechanism can be leveraged to trigger avalanche Rydberg excitations, enabling amplified readout of a single initial Rydberg excitation. If such an initial excitation is generated as part of a sensing process, for example, as a consequence of detecting a microwave or terahertz photon [8], this low-dark-count, high-gain amplifier may assist a variety of rare-event detection, including dark-matter searches [22,23].

II. MODEL

Consider a one dimensional (1D) array of N sites. Each site is occupied by a single atom. A driving laser field with time-dependent detuning $\Delta(t)$ and site-specific, time-dependent Rabi frequency $\Omega_j(t)$ couples the ground state $|0\rangle_j$ to the Rydberg state $|1\rangle_j$ at site j . The single-body contribution to the Hamiltonian is given by

$$H_{\text{site}}(t) = \frac{1}{2} \sum_{j=1}^N \Omega_j(t) \sigma_j^x - \frac{\Delta(t)}{2} \sum_{j=1}^N \sigma_j^z, \quad (1)$$

where $\sigma_j^x = |1\rangle_{jj}\langle 0| + |0\rangle_{jj}\langle 1|$ and $\sigma_j^z = |1\rangle_{jj}\langle 1| - |0\rangle_{jj}\langle 0|$.

The van der Waals interaction between Rydberg atoms scales as $V_r = C_6/|r|^6$, where C_6 is the van der Waals interaction coefficient (along the array direction if anisotropic). The resulting nearest-neighbor interaction in a chain with spacing r is

$$H_{\text{NN}} = V_r \sum_{j=1}^{N-1} n_j n_{j+1}, \quad (2)$$

where $n_i = |1\rangle_{ii}\langle 1|$. We operate in the “facilitated” or “antiblockade” regime: The ground-to-Rydberg transition is driven off-resonantly. The kinetic constraint ensures that resonance is achieved if and only if one of the nearest neighbors is in the Rydberg state. To execute RAP, we slowly ramp the laser detuning $\Delta(t)$ through this interaction-shifted resonant condition, $\Delta_0 = V_r$.

III. PULSE SEQUENCE: TWO-ATOM DEMONSTRATION

First, we study a two-site case where the total Hamiltonian is $H = H_{\text{site}}(t) + H_{\text{NN}}$. Initially, a Rydberg atom is located in site 1, whereas site 2 is occupied by a ground-state atom. We alternatively drive the two atoms. When driving atom 2 with Rabi frequency Ω_2 , the initial state $|10\rangle$ is coupled to $|11\rangle$. The ramping step ① adiabatically follows the resulting upper dressed state [Fig. 1(a)]. The second ramping step drives atom 1, coupling $|11\rangle \leftrightarrow |01\rangle$. By adiabatically sweeping the detuning Δ from above V_r to below, the state is transferred to $|01\rangle$. In Fig. 1(c), the colored balls encode on-site populations: Red denotes the Rydberg state $|1\rangle$, and black denotes the ground state $|0\rangle$. A red $\leftrightarrow$ black change marks an RAP. The rectangles indicate which atom is driven, with the color gradient showing the sweep direction of Δ from top to bottom as time evolves.

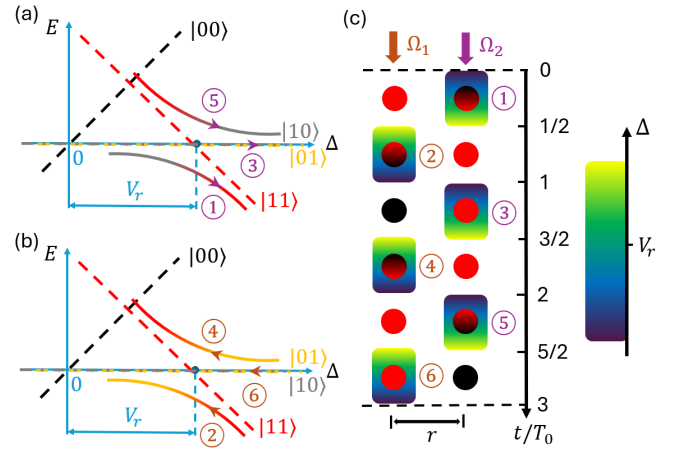


FIG. 1. Robust Rydberg facilitation in a two-site array. (a) Sketch of energy diagrams as a function of laser detuning Δ for driving atom 2. (b) Same for driving atom 1. The dashed lines indicate bare states (red $|11\rangle$, gray $|10\rangle$, orange $|01\rangle$, and black $|00\rangle$). (c) Evolution of the Rydberg excitation under the periodic driving. ① $|10\rangle \rightarrow |11\rangle$, ② $|11\rangle \rightarrow |01\rangle$, ③ stays in $|01\rangle$, ④ $|01\rangle \rightarrow |11\rangle$, ⑤ $|11\rangle \rightarrow |10\rangle$, and ⑥ stays in $|10\rangle$.

Steps ③ and ⑥ do not result in any population transfer due to the energy penalty of interaction energy V_r . These steps would lead to population transfer if the detuning Δ is swept across 0 where the state $|10\rangle$ or $|01\rangle$ crosses state $|00\rangle$ in the energy diagram. The driving parameters complete one cycle T_0 in two steps, whereas the excitation hops between sites 1 and 2 with a period of $3T_0$.

We deploy the Allen-Eberly adiabatic passage scheme [24,25]. The Rabi frequencies and detuning evolve as follows within one cycle $0 \leq t < T_0$. For the first half-cycle ($0 \leq t < T_0/2$),

$$\Omega_1(t) = 0, \quad \Omega_2(t) = \Omega_0 \operatorname{sech} \left[\frac{\pi(t - T_0/4)}{T_0/8} \right], \quad (3a)$$

$$\Delta(t) = \Delta_0 + \left(\frac{\beta^2 T_0}{8\pi} \right) \tanh \left[\frac{\pi(t - T_0/4)}{T_0/8} \right]. \quad (3b)$$

For the second half-cycle ($T_0/2 \leq t < T_0$),

$$\Omega_1(t) = \Omega_0 \operatorname{sech} \left[\frac{\pi(t - 3T_0/4)}{T_0/8} \right], \quad \Omega_2(t) = 0, \quad (4a)$$

$$\Delta(t) = \Delta_0 - \left(\frac{\beta^2 T_0}{8\pi} \right) \tanh \left[\frac{\pi(t - 3T_0/4)}{T_0/8} \right]. \quad (4b)$$

The sweep direction of the detuning can be swapped between the first and the second half-cycle without affecting any population transfer. From an experimental perspective, this protocol requires site-resolved laser intensity ramping, but the detuning ramping is global, which is straightforward to implement.

IV. ROBUSTNESS TO IMPERFECTIONS

To evaluate the robustness of our scheme, we consider implementation in ^{87}Rb tweezer arrays and compare RAP with a resonant Rabi scheme based on sequential π pulses. We focus on one-photon excitation to a P Rydberg state. At

room temperature (300 K), the decay of Rydberg states is dominated by blackbody-radiation-induced transitions. In the more common two-photon excitation scheme, admixture of a short-lived intermediate state can significantly increase the decay rate. Because RAP strongly suppresses the impact of disorder and sensitivity to driving parameters, Rydberg-state decay becomes the primary limitation. We therefore adopt the one-photon excitation to minimize decay and highlight the strength of RAP. The case of a two-photon excitation to an S or D Rydberg state, where RAP remains advantageous, is discussed in the Supplemental Material [26].

To simulate the dynamics, we use QuTiP [27,28]. The Rydberg decay is described by the Lindblad superoperator:

$$\mathcal{L}_{\text{decay}}(\rho) = \frac{\Gamma}{2} \sum_{j=1}^N (2\sigma_j^- \rho \sigma_j^+ - \sigma_j^+ \sigma_j^- \rho - \rho \sigma_j^+ \sigma_j^-), \quad (5)$$

where $\sigma_i^+ = |1\rangle_{ii}\langle 0|$ and $\sigma_i^- = |0\rangle_{ii}\langle 1|$. The density matrix ρ evolves according to the Lindblad master equation:

$$\frac{d\rho}{dt} = i[\rho, H] + \mathcal{L}_{\text{decay}}(\rho). \quad (6)$$

For the Rabi scheme, the laser detuning is set at the interaction-shifted resonance Δ_0 . Each step in Fig. 1 is implemented by applying a π pulse of duration $T_0/2 = \pi/\Omega_0$ with constant Rabi frequency Ω_0 . The sequence of driven atoms and the intended state transfers are identical to those in the RAP scheme.

A key challenge for the Rabi scheme is position disorder, which introduces static deviations in the interaction energy [10,29–31], characterized by

$$\overline{\delta V_r} \approx 6|V_r|^{7/6} \frac{\sqrt{2}\sigma_r}{|C_6|^{1/6}}, \quad (7)$$

where σ_r denotes the root-mean-square radius of the atomic position distribution along the array direction. Although going to higher principal quantum number n and reducing interaction V_r alleviate the impact of position disorder, the required interatomic spacing can become too large to fit a reasonable number of sites within the imaging field of view. Moreover, high n states are more vulnerable to stray electric fields [32–34]. Taking $70P_{3/2}$ as a reasonable choice, we conducted a limited parameter search to find workable parameters that balance disorder sensitivity with decay. This yields $V_r/(2\pi) = 1.1$ MHz and $\Omega_0/(2\pi) = 0.34$ MHz, accounting for Rydberg decay rate $\Gamma/(2\pi) = 0.839$ kHz and position disorder corresponding to $\sigma_r = 34$ nm. The disorder is modeled by averaging over 50 realizations, in which each atom's displacement from its ideal position is randomly sampled from a Gaussian distribution with standard deviation σ_r .

This level of position distribution can be achieved by Raman sideband cooling [35,36] or narrow-line cooling [37], approaching the motional ground state [38]. To evaluate performance under these conditions, we simulate a two-atom array driven for $t = 4T_0$. Figure 2(a) shows the evolution of the Rydberg excitation, initially located at site 1, following the RAP scheme, without decay or disorder. An ideal Rabi scheme reproduces the same state transfers at the end of each step. Figures 2(b)–2(d) illustrate how position disorder, decay, and their combination affect the dynamics, while Fig. 2(e)

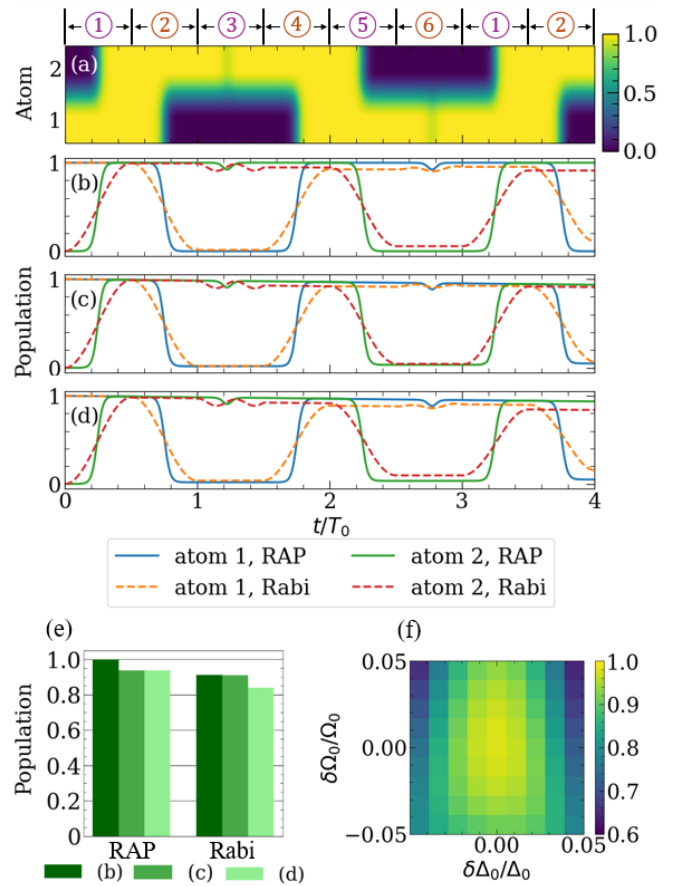


FIG. 2. Comparison of sensitivity to imperfections between RAP and Rabi schemes. (a) Rydberg population evolution under the pulse sequences of Fig. 1 and Eqs. (3) and (4), without position disorder or decay. Evolution subject to (b) position disorder, (c) Rydberg-state decay Γ , and (d) both effects. Parameters are $\Omega_0/(2\pi) = \{10.7, 0.34\}$ MHz, $\sigma_r = \{76, 34\}$ nm, $\overline{\delta V_r}/(2\pi) = \{1.37, 0.02\}$ MHz, $V_r/(2\pi) = \{20, 1.1\}$ MHz, and $r = \{6.4, 10.4\}$ μm for the RAP and Rabi schemes, respectively. (e) Final Rydberg population of atom 2 at $t = 4T_0$, extracted from panels (b)–(d). (f) Rydberg population of atom 2 at $t = 4T_0$ for the Rabi scheme as a function of detuning error $\delta\Delta_0$ and Rabi-frequency error $\delta\Omega_0$, without disorder or decay.

presents the corresponding final Rydberg population at site 2. Under our parameter choice, both effects comparably reduce its value from the ideal value of unity. We emphasize that we have assumed extreme cooling and very large atomic spacing to mitigate position disorder; a naive Rabi implementation would suffer much more severe degradation from position disorder.

For RAP, the protocol speed is set by Ω_0 , which is technically limited by available laser power. Currently, 3 W laser at the transition ($5S_{1/2} \leftrightarrow 70P_{3/2}$) wavelength 297 nm is available, and future experiments may further enhance the Rabi frequency using a build-up cavity. In any case, laser power is a less fundamental limitation than the spread of the ground-state wave function. We choose $\Omega_0/(2\pi) = 10.7$ MHz, which requires 2.6 mW per site for a beam waist ($1/e^2$ radius) of 1.5 μm . With $T_0 = 3$ μs and $\beta/(2\pi) = 7.6$ MHz, determined by the Allen-Eberly pulse shape and the chosen Ω_0 [25,39], we find that the final Rydberg population at site 2

is reduced only by the blackbody-radiation-limited Rydberg lifetime [Figs. 2(b)–2(e)]. Position disorder has no noticeable effect at a much larger atomic position spread. The value $\sigma_r = 76$ nm used in our RAP simulation can be achieved by a radial trap frequency of 91 kHz at a temperature of 20 μ K, readily accessible by a simple and efficient cooling method: Λ -enhanced gray molasses [18,40]. The error resulting from adiabaticity breakdown is negligible, although there is a minor violation in steps ③ and ⑥ [see Figs. 2(b)–2(d)]. Detailed discussion about adiabaticity is provided in the Supplemental Material [26].

In addition to position disorder and decay, we also consider imperfections in the driving parameters, namely, deviations of the detuning Δ_0 by $\delta\Delta_0$ and of the Rabi frequency Ω_0 by $\delta\Omega_0$. As shown in Fig. 2(f), for the Rabi scheme, a $\pm 5\%$ change in both can reduce the final Rydberg population from unity to about 0.6 ($\delta\Omega_0 > 0$, 0.8 for $\delta\Omega_0 < 0$). The apparent trend that larger Ω_0 performs worse than its smaller counterpart is coincidental: If the sequence included one more or one fewer step of excitation transfer, the comparison could reverse and the larger Ω_0 would appear more favorable. For the RAP scheme, we do not show a plot because it is essentially insensitive to any realistic driving parameter error: Scanning the same relative variations $\delta\Delta_0/\Delta_0$ and $\delta\Omega_0/\Omega_0$ changes the final population by less than 2×10^{-4} .

V. AVALANCHE EXCITATIONS: 1D CHAIN

Next, we extend the system from the two-atom array to an N -atom one-dimensional chain. We continue using the same parameters for the chain as for the two-atom case. We do not consider position disorder or imperfections of driving parameters as they do not affect the Rydberg population evolution. The driving of sites 1 and 2 naturally generalizes to all odd and even sites, respectively, while all other elements of the RAP protocol, including the Allen-Eberly pulse shapes, the cyclic sequence structure, the overall mechanism, and parameter choices, remain unchanged. In this case, the site-dependent Rabi frequency takes the form

$$\Omega_j(t) = \begin{cases} \Omega_1(t), & \text{if } j \bmod 2 = 1, \\ \Omega_2(t), & \text{if } j \bmod 2 = 0. \end{cases} \quad (8)$$

Because of the sharp distance dependence of the interaction, the next-nearest-neighbor (NNN) coupling is $2^6 = 64$ times weaker than the nearest-neighbor (NN) term. Numerical checks confirm that including NNN interactions alters the results by less than one part in 10^4 and can thus be neglected.

To illustrate how avalanche growth proceeds, we simulate a five-atom chain and track how an initial seed excitation expands step by step (Fig. 3). Unlike in the two-atom case, here each half period leads to a state transfer. The excitation spreads symmetrically outward, adding one atom at each end in every step. Once an atom inside the cluster is excited, it is no longer driven down, because having two nearest-neighbor excitations shifts its resonance outside the RAP sweeping range. If the odd sites were driven in the very first step, no state transfer would occur in that step; from the second step onward, the growth is identical to the case where even sites are driven first. The cluster expands until it reaches the array boundary, after which it begins to shrink. To achieve a net gain

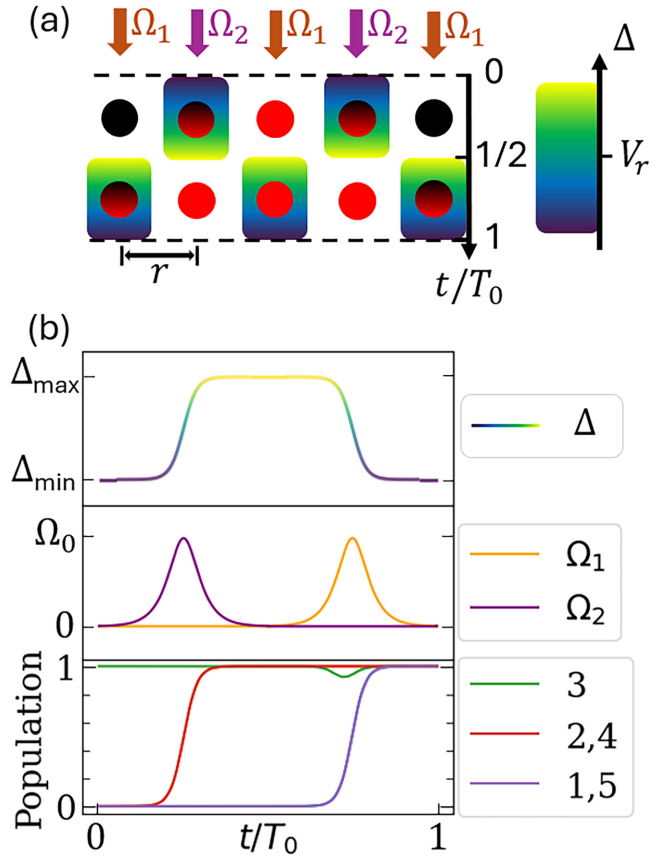


FIG. 3. Mechanism of avalanche excitation growth. After one driving period ($t = T_0$), all atoms are transferred to the Rydberg state. The detuning Δ is swept between $\Delta_{\max, \min} = \Delta_0 \pm \beta^2 T_0 / (8\pi)$. The bottom panel of plot (b) shows the Rydberg population for each site: site 3 (green), sites 2 and 4 (red), and sites 1 and 5 (purple).

of $2S + 1$, S steps are required, and the initial seed must have at least S available sites on each side of the chain.

Having established the avalanche mechanism, we now quantify its gain and noise performance. Figure 4 shows avalanche excitations starting from a single seed in the middle of a nine-atom chain. At the end of two driving periods $t = 2T_0$, the total number of excitations (gain) reaches 8.54. The reduction of gain from its ideal value $2S + 1 = 9$ stems almost entirely from decay: Without decay, the gain falls short of 9 only at the 10^{-3} level. By comparison, when starting with no excitation, the total number of excitations at this time is 7×10^{-4} . This background can be further suppressed by

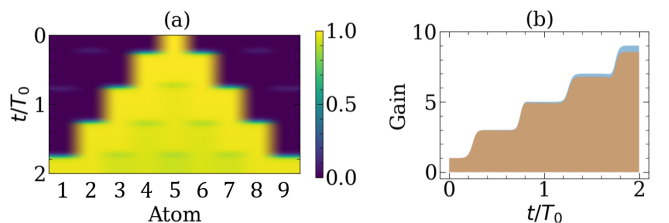


FIG. 4. Avalanche gain. (a) Rydberg population evolution. (b) Gain vs time: At $t = 2T_0$, the no-decay value (blue) is short of the ideal value 9 by 0.003, while with decay (brown) it is 8.54.

increasing V_r . For example, when $V_r = 80$ MHz, the noise signal drops to 5×10^{-6} . This behavior is directly analogous to a single-photon detector, such as a photomultiplier tube or a single-photon avalanche diode: The system exhibits a high-gain, avalanche response triggered by a single seed excitation, while maintaining an exceptionally low dark count in its absence. Although the simulation uses a nine-atom chain, the result reflects the bulk behavior of longer arrays. In larger systems, the avalanche can propagate across many more sites, leading to correspondingly greater gain, with the ultimate limit set by decay.

VI. GUTZWILLER MEAN-FIELD APPROACH

However, we can only numerically solve the master equation [Eq. (6)] for up to approximately ten atoms. One common approach to this problem is to use Gutzwiller mean-field theory [41], which ignores the effects of correlations and assumes a product state: $\psi = \prod_j \psi_j$. Here, we combine the relative small effective Hilbert space of the Gutzwiller ansatz with quantum Monte Carlo wave function (QMCWF) simulation [42] to recover stochastic correlations generated by decay events, which are otherwise omitted if using a mean-field master equation. Within this framework, the dynamics of each atom j is governed by an effective single-site non-Hermitian Hamiltonian:

$$H_j^{\text{MF}}(t) = \frac{\Omega_j(t)}{2} \sigma_j^x - \frac{\tilde{\Delta}_j(t)}{2} \sigma_j^z - i \frac{\Gamma}{2} \sigma_j^+ \sigma_j^- \quad (9)$$

with effective detuning

$$\tilde{\Delta}_j(t) = \Delta(t) - \sum_{k \in \mathcal{N}_j} V_{jk} \langle n_k \rangle(t), \quad (10)$$

where $\langle n_k \rangle(t) = \langle \psi_k(t) | n_k | \psi_k(t) \rangle$, $V_{jk} = \frac{C_6}{r_{jk}^6}$, r_{jk} is the separation between sites j and k , and $\mathcal{N}_j$ denotes neighbors of site j that interact with it. The instantaneous rate for a quantum jump to occur on site j is

$$\frac{dp_j}{dt} = \Gamma \langle \psi_j(t) | \sigma_j^+ \sigma_j^- | \psi_j(t) \rangle. \quad (11)$$

We first validate this mean field with QMCWF approach by reproducing the 1D avalanche dynamics from Fig. 4, finding agreement within 0.3% using 700 quantum trajectories. Figure 5(a) compares the three approaches under the same conditions as Fig. 4: full master equation (Full ME), mean field with QMCWF (MF-QMC), and mean-field master equation (MF-ME). The mean-field master equation underestimates the impact of decay. It shows that earlier excitations are more likely to have decayed, leading to the lowest Rydberg population in the middle of the chain, but it does not capture the correlations. In this approach, decay only slightly reduces the average Rydberg population and hence the interaction with its neighbors, leaving the dynamics largely unchanged because it is robust to small imperfections. In reality, quantum correlations mean that (1) a decay event at the edge of a cluster can terminate excitation growth from that site and (2) a decay event in the bulk can seed the avalanche growth of a hole (ground-state atom) whenever only one excited neighbor is present. In other words, holes obey the same kinetic constraint

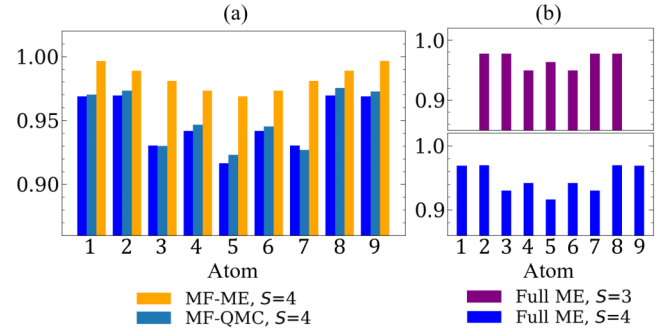


FIG. 5. Comparison of different simulation methods: full master equation (Full ME, blue), mean field with QMCWF (MF-QMC, sky blue), and mean-field master equation (MF-ME, orange). The blue bars in panels (a) and (b) show the same result as $t = 2T_0$ in Fig. 4. Panel (b) shows the Rydberg population pattern after $S = 3$ steps (top) and after $S = 4$ steps (bottom), illustrating two representative examples of the spatially oscillatory profile.

as excitations, so the dynamics involves the correlated growth of both. For example, a decay at site 5 during step 2 would lead to hole growth to sites 4 and 6 by the end of step 3.

Assuming small Γ and restricting to at most one Rydberg decay event, we can explicitly include all possible hole-growth trajectories. This treatment reproduces the spatially oscillatory profiles seen in Fig. 4(b): At a fixed step, the site populations within the excitation cluster exhibit an up-down modulation across neighboring sites. From one step to the next, the phase of this oscillation shifts by one site, so that the weakest-population sites switch parity (even sites at odd S , odd sites at even S). The same treatment further yields the analytical prediction for the gain after S steps:

$$G = 2S + 1 - \frac{\Gamma T_0}{2} \left(S^2 + \frac{2}{3} S^3 - \frac{2}{3} S - 1 \right), \quad (12)$$

as derived in the Supplemental Material [26], which also traces the oscillatory profiles back to hole-growth trajectories. The final gain for Fig. 4 from this equation is $G = 8.57$, in good agreement with the full-master-equation result. This equation also predicts the maximum gain of $G = 15.1$ [Fig. 6(a)], reached at a finite step before decay reduces the excitation, assuming the array is sufficiently large that boundary effects do not constrain the dynamics. To test this prediction, we extend the array size to 33 atoms and use MF-QMC to

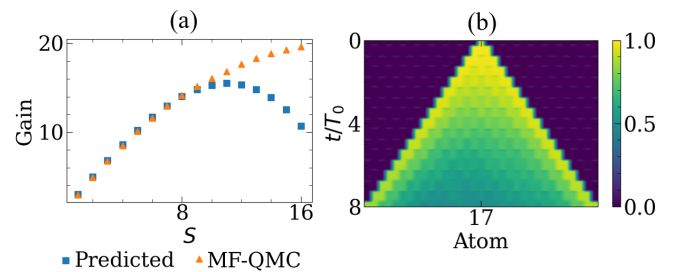


FIG. 6. Gain vs step number predicted by Eq. (12) (blue) and obtained from MF-QMC simulations (orange, averaged over 500 quantum trajectories). (b) Site-resolved population dynamics corresponding to the orange points in panel (a).

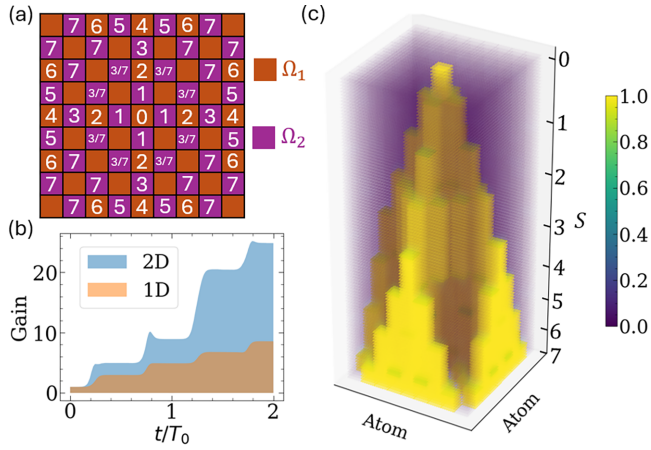


FIG. 7. Avalanche excitation in a 2D square lattice. (a) Atoms are located at the centers of plaquettes instead of at the vertices. Sites driven during the first (second) half-cycle are drawn in purple (brown). The site numbered “0” is the location of the initial seed. Other numbers denote the step number during which that atom is excited. The sites labeled as “3/7” are excited in step 3, de-excited in step 5, and re-excited in step 7. (b) Gain vs time for 1D [taken from Fig. 4(b)] and 2D (MF-QMC simulations, averaged over 100 quantum trajectories). (c) Stacked images of Rydberg population evolution up to seven steps for one trajectory that does not encounter any decay event.

obtain gain versus step number S (Fig. 6). Initially, the gain follows Eq. (12) closely. For larger S , however, the analytical expression underestimates the gain, because it neglects multiple decay events. In particular, a second decay event at the boundary of a hole can terminate its avalanche growth (see the Supplemental Material [26] for further discussion of dynamics associated with multiple decay events).

VII. AVALANCHE EXCITATIONS: 2D SQUARE LATTICE

Motivated by experimental feasibility, we next consider a two dimensional (2D) square lattice. A 2D geometry is advantageous because (1) it is more practical to scale up atom numbers compared to 1D chains, (2) the initial sensing signal has a larger receiving area, (3) the avalanche gain per unit time is greater, and (4) Rydberg decay is less detrimental, leading to a higher maximum gain. In our 2D simulations, each atom interacts with its eight nearest and next-nearest neighbors. We apply driving lasers in alternating checkerboard patterns [Fig. 7(a)], analogous to the even-odd site driving in 1D. The parameters follow the same ramping curves [Eqs. (3) and (4)], but with a shifted center detuning of the ramp: $\Delta_0 = V_r + V_r/(\sqrt{2})^6 = 22.5$ MHz, chosen as a middle-ground value to balance the different possible NNN interaction shifts. If the initial seed is not driven in the first step, the excitation spreads to its four nearest neighbors. If instead the seed is driven in the first step, no population transfer occurs in that step; from the second step onward, the growth is identical to the case where the seed is not driven first, just as in the 1D dynamics.

As shown in Figs. 7(a) and 7(b), the gain per step in 2D is at least twice as large as in 1D. Moreover, the 2D geometry is more resilient to decay. In 1D, an excitation adjacent to a

hole has only one neighboring excitation. This one neighbor provides the interaction shift that moves the transition into the RAP sweeping range, enabling de-excitation and hole growth. In 2D, however, an excitation adjacent to a hole may have multiple excited neighbors, whose combined interaction shifts move the transition outside the RAP sweeping range, thereby suppressing avalanche hole growth. Consequently, at the end of four steps, the gain with decay reaches 99.5% of the corresponding no-decay gain, as compared to 95% for 1D. More accurately, this high resilience is not only due to suppressed hole growth: Frozen sites can become flippable through decay. This process enhances gain and further inhibits hole expansion. From a many-body physics perspective, these dynamics may offer insights into driven, dissipative systems. We briefly discuss this in the Supplemental Material [26] and leave a full exploration to future work.

Setting aside decay, the coherent dynamics alone exhibits richer spatial-temporal patterns in 2D. The cluster growth along the main axes (the vertical and horizontal line “432101234”) is compact, filling all sites along the line. Along other directions, sites get excited and de-excited multiple times because the neighboring sites are not all filled. If the array is big enough so that the excitation patterns do not change due to boundary effects, all the sites labeled “5”, “6”, and “7” will be de-excited and re-excited in the following steps. Within the seven steps evolved here, only the sites labeled “3/7” undergo excitation-deexcitation cycles. The complete Rydberg population evolution pattern presented in Fig. 7(a) is verified by evolving a single decay-free quantum trajectory [Fig. 7(c)]. If the goal is to excite every site, one may consider ditching our generic pulse sequence that naturally extends from the 1D geometry. Careful engineered sequences, such as step-specific ramping curves, could further enhance gain while also suppressing correlated hole growth and excitation-deexcitation cycles.

VIII. DISCUSSION

In conclusion, we have demonstrated robust Rydberg antiblockade in neutral-atom arrays. By sweeping through the interaction-shifted resonance, our approach circumvents the sensitivity to position disorder and laser parameter variations that hinder antiblockade-based protocols. Numerical simulations confirm that excitation transfer remains unaffected under realistic experimental imperfections, with performance ultimately limited only by Rydberg decay. Our simulations adopt a blackbody-radiation-limited Rydberg decay at room temperature, which can be readily suppressed using cryogenic environments now available in state-of-the-art atom-array platforms [43,44]. We compare our RAP scheme with a resonant Rabi protocol and find similar timescales after preliminary optimization, resulting in comparable Rydberg decay for both. The timescale of RAP is constrained by available laser power, while that of Rabi is constrained by position disorder. Importantly, the Rabi scheme suffers substantial degradation from position disorder even when cooled close to the quantum limit, whereas RAP is immune to these imperfections at temperatures achievable through simple cooling techniques.

As an application, we demonstrate avalanche excitation growth as a high-gain, low-dark-count amplification mechanism for sensing. In 2D geometries, the excitations grow more efficiently, enhancing amplification. Quantum-trajectory analysis reveals rich correlated excitation and hole dynamics, highlighting the complex driven-dissipative behavior of this kinetically constrained system. Future work will explore emergent functionalities of such nonequilibrium dynamics and the opportunities offered by other lattice geometries, particularly frustrated ones. More broadly, RAP-enabled antiblockade is a general framework that can be applied beyond avalanche excitation. Adiabatic protocols have been proposed to accumulate dynamical phases for two-qubit gates while navigating complex level structure at small interatomic spacings [45]. It may be possible to accumulate geometric phases by echoing two identical RAP pulses near the

interaction-shifted resonance [46], thereby creating entangling gates that are robust to position disorder.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

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- [1] L. Henriet, L. Beguin, A. Signoles, T. Lahaye, A. Browaeys, G.-O. Reymond, and C. Jurczak, Quantum computing with neutral atoms, *Quantum* **4**, 327 (2020).
 - [2] T. Graham, Y. Song, J. Scott, C. Poole, L. Phuttitarn, K. Jooya, P. Eichler, X. Jiang, A. Marra, B. Grinkemeyer, *et al.*, Multi-qubit entanglement and algorithms on a neutral-atom quantum computer, *Nature (London)* **604**, 457 (2022).
 - [3] J. Muniz, M. Stone, D. Stack, M. Jaffe, J. Kindem, L. Wadleigh, E. Zalsys-Geller, X. Zhang, C.-A. Chen, M. Norcia, *et al.*, High-fidelity universal gates in the ^{171}Yb ground-state nuclear-spin qubit, *PRX Quantum* **6**, 020334 (2025).
 - [4] W.-X. Li, J.-L. Wu, S.-L. Su, and J. Qian, High-tolerance antiblockade swap gates using optimal pulse drivings, *Phys. Rev. A* **109**, 012608 (2024).
 - [5] X.-F. Shi, Rydberg quantum gates free from blockade error, *Phys. Rev. Appl.* **7**, 064017 (2017).
 - [6] J.-L. Wu, Y. Wang, J.-X. Han, Y.-K. Feng, S.-L. Su, Y. Xia, Y. Jiang, and J. Song, One-step implementation of Rydberg-antiblockade swap and controlled-swap gates with modified robustness, *Photon. Res.* **9**, 814 (2021).
 - [7] Y. Zhao, B. Liu, Y. Ji, S. Tang, and X. Shao, Robust generation of entangled state via ground-state antiblockade of Rydberg atoms, *Sci. Rep.* **7**, 16489 (2017).
 - [8] C. Nill, A. Cabot, A. Trautmann, C. Groß, and I. Lesanovsky, Avalanche terahertz photon detection in a Rydberg tweezer array, *Phys. Rev. Lett.* **133**, 073603 (2024).
 - [9] P. R. Datla, L. Zhao, W. W. Ho, N. Klco, and H. Loh, Statistical localization in a Rydberg simulator of $U(1)$ lattice gauge theory, *arXiv:2505.18143*.
 - [10] L. Zhao, P. R. Datla, W. Tian, M. M. Aliyu, and H. Loh, Observation of quantum thermalization restricted to Hilbert space fragments and $\mathbb{Z}_{2k}$ scars, *Phys. Rev. X* **15**, 011035 (2025).
 - [11] T. Zhang and Z. Cai, Quantum slush state in Rydberg atom arrays, *Phys. Rev. Lett.* **132**, 060503 (2024).
 - [12] F. Liu, Z.-C. Yang, P. Bienias, T. Iadecola, and A. V. Gorshkov, Localization and criticality in antiblockaded two-dimensional Rydberg atom arrays, *Phys. Rev. Lett.* **128**, 013603 (2022).
 - [13] L. Causer, I. Lesanovsky, M. C. Bañuls, and J. P. Garrahan, Dynamics and large deviation transitions of the XOR-Fredrickson-Andersen kinetically constrained model, *Phys. Rev. E* **102**, 052132 (2020).
 - [14] L. Zadnik and J. P. Garrahan, Slow heterogeneous relaxation due to constraints in dual XXZ models, *Phys. Rev. B* **108**, L100304 (2023).
 - [15] X. Li, S. Zhou, X. Chen, C. Li, and H. Wang, Random displacements in critical Rydberg atom arrays, *Phys. Rev. B* **112**, L180202 (2025).
 - [16] S.-L. Su, F.-Q. Guo, J.-L. Wu, Z. Jin, X. Shao, and S. Zhang, Rydberg antiblockade regimes: Dynamics and applications, *Europhys. Lett.* **131**, 53001 (2020).
 - [17] Y. Chew, T. Tomita, T. P. Mahesh, S. Sugawa, S. de Léséleuc, and K. Ohmori, Ultrafast energy exchange between two single Rydberg atoms on a nanosecond timescale, *Nat. Photon.* **16**, 724 (2022).
 - [18] S. J. Evered, D. Bluvstein, M. Kalinowski, S. Ebadi, T. Manovitz, H. Zhou, S. H. Li, A. A. Geim, T. T. Wang, N. Maskara, *et al.*, High-fidelity parallel entangling gates on a neutral-atom quantum computer, *Nature (London)* **622**, 268 (2023).
 - [19] A. Radnaev, W. Chung, D. Cole, D. Mason, T. Ballance, M. Bedalov, D. Belknap, M. Berman, M. Blakely, I. Bloomfield, *et al.*, Universal neutral-atom quantum computer with individual optical addressing and nondestructive readout, *PRX Quantum* **6**, 030334 (2025).
 - [20] C. B. Dag, H. Ma, P. M. Eugenio, F. Fang, and S. F. Yelin, Emergent disorder and sub-ballistic dynamics in quantum simulations of the Ising model using Rydberg atom arrays, *Phys. Rev. Lett.* **135**, 250403 (2025).
 - [21] L. Zhao, M. D. K. Lee, M. M. Aliyu, and H. Loh, Floquet-tailored Rydberg interactions, *Nat. Commun.* **14**, 7128 (2023).
 - [22] E. Graham, S. Ghosh, Y. Zhu, X. Bai, S. B. Cahn, E. Durcan, M. J. Jewell, D. H. Speller, S. M. Zacarias, L. T. Zhou, *et al.*, Rydberg-atom-based single-photon detection for haloscope axion searches, *Phys. Rev. D* **109**, 032009 (2024).
 - [23] G. Engelhardt, A. Bhonah, and W. V. Liu, Detecting axion dark matter with Rydberg atoms via induced electric dipole transitions, *Phys. Rev. Res.* **6**, 023017 (2024).
 - [24] L. Allen and J. H. Eberly, *Optical Resonance and Two-Level Atoms* (John Wiley and Sons, Inc., New York, 1974).

- [25] J. da Silva Andrade, Â. F. da Silva França, and B. de Lima Bernardo, Shortcuts to adiabatic population inversion via time-rescaling: stability and thermodynamic cost, *Sci. Rep.* **12**, 11538 (2022).
- [26] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/23x7-g2h6> for additional information.
- [27] J. R. Johansson, P. D. Nation, and F. Nori, QuTiP: An open-source python framework for the dynamics of open quantum systems, *Comput. Phys. Commun.* **183**, 1760 (2012).
- [28] N. Lambert, E. Giguère, P. Menczel, B. Li, P. Hopf, G. Suárez, M. Gali, J. Lishman, R. Gadhvi, R. Agarwal, *et al.*, QuTiP 5: The quantum toolbox in Python, *Phys. Rep.* **1153**, 1 (2026).
- [29] Y. Wang, J. Wang, A. Panja, X. Wang, and Q.-Y. Liang, Directional transport in Rydberg atom arrays via kinetic constraints and temporal modulation, *Phys. Rev. Res.* **7**, L022035 (2025).
- [30] M. Marcuzzi, J. Minář, D. Barredo, S. De Léséleuc, H. Labuhn, T. Lahaye, A. Browaeys, E. Levi, and I. Lesanovsky, Facilitation dynamics and localization phenomena in Rydberg lattice gases with position disorder, *Phys. Rev. Lett.* **118**, 063606 (2017).
- [31] R. J. Valencia-Tortora, N. Pancotti, M. Fleischhauer, H. Bernien, and J. Marino, Rydberg platform for nonergodic chiral quantum dynamics, *Phys. Rev. Lett.* **132**, 223201 (2024).
- [32] A. Panja, Y. Wang, X. Wang, J. Wang, S. Subhankar, and Q.-Y. Liang, Electric field control for experiments with atoms in Rydberg states, *AIP Adv.* **14**, 125013 (2024).
- [33] S. Anand, C. E. Bradley, R. White, V. Ramesh, K. Singh, and H. Bernien, A dual-species Rydberg array, *Nat. Phys.* **20**, 1744 (2024).
- [34] J. T. Wilson, S. Saskin, Y. Meng, S. Ma, R. Dilip, A. P. Burgers, and J. D. Thompson, Trapping alkaline earth Rydberg atoms optical tweezer arrays, *Phys. Rev. Lett.* **128**, 033201 (2022).
- [35] N. Lorenz, L. Festa, L.-M. Steinert, and C. Gross, Raman side-band cooling in optical tweezer arrays for Rydberg dressing, *SciPost Phys.* **10**, 052 (2021).
- [36] A. Jenkins, J. W. Lis, A. Senoo, W. F. McGrew, and A. M. Kaufman, Ytterbium nuclear-spin qubits in an optical tweezer array, *Phys. Rev. X* **12**, 021027 (2022).
- [37] K. N. Blodgett, S. S. Phatak, M. R. Chen, D. Peana, C. Pritts, and J. D. Hood, Narrow-line electric quadrupole cooling and background-free imaging of a single Cs atom with spatially structured light, *Phys. Rev. A* **112**, 043109 (2025).
- [38] Y. T. Chew, M. Poitrinal, T. Tomita, S. Kitade, J. Mauricio, K. Ohmori, and S. de Léséleuc, Ultraprecise holographic optical tweezer array, *Phys. Rev. A* **110**, 053518 (2024).
- [39] N. V. Vitanov and B. M. Garraway, Landau-Zener model: Effects of finite coupling duration, *Phys. Rev. A* **53**, 4288 (1996).
- [40] S. Rosi, A. Burchianti, S. Conclave, D. S. Naik, G. Roati, C. Fort, and F. Minardi, Λ -enhanced grey molasses on the D_2 transition of rubidium-87 atoms, *Sci. Rep.* **8**, 1301 (2018).
- [41] D. Jaksch, C. Bruder, J. I. Cirac, C. W. Gardiner, and P. Zoller, Cold bosonic atoms in optical lattices, *Phys. Rev. Lett.* **81**, 3108 (1998).
- [42] W. Casteels, R. M. Wilson, and M. Wouters, Gutzwiller Monte Carlo approach for a critical dissipative spin model, *Phys. Rev. A* **97**, 062107 (2018).
- [43] Z. Zhang, T.-W. Hsu, T. Y. Tan, D. H. Slichter, A. M. Kaufman, M. Marinelli, and C. A. Regal, High optical access cryogenic system for Rydberg atom arrays with a 3000-second trap lifetime, *PRX Quantum* **6**, 020337 (2025).
- [44] G. Pichard, D. Lim, E. Bloch, J. Vaneeckloo, L. Bourachot, G.-J. Both, G. Mériaux, S. Dutartre, R. Hostein, J. Paris, B. Ximenez, A. Signoles, A. Browaeys, T. Lahaye, and D. Dreon, Rearrangement of individual atoms in a 2000-site optical-tweezer array at cryogenic temperatures, *Phys. Rev. Appl.* **22**, 024073 (2024).
- [45] A. Mitra, S. Omanakuttan, M. J. Martin, G. W. Biedermann, and I. H. Deutsch, Neutral-atom entanglement using adiabatic Rydberg dressing, *Phys. Rev. A* **107**, 062609 (2023).
- [46] M. Xue, S. Xu, X. Li, and X. Li, High-fidelity and robust controlled-Z gates implemented with Rydberg atoms via echoing rapid adiabatic passage, *Phys. Rev. A* **110**, 032619 (2024).