

Adiabatic Echo Protocols for Robust Quantum Many-Body State Preparation

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Entangled many-body states are a key resource for quantum technologies. Yet their preparation through analog control of interacting quantum systems is often hindered by experimental imperfections. Here, we introduce the adiabatic echo protocol, a general approach to state preparation designed to suppress the effect of static perturbations. We provide an analytical understanding of its robustness in terms of dynamically engineered destructive interference. By applying quantum optimal control methods, we demonstrate that such a protocol emerges naturally in a variety of settings, without requiring assumptions on the form of the control fields. Examples include Greenberger-Horne-Zeilinger state preparation in Ising spin chains and two-dimensional Rydberg atom arrays, as well as the generation of quantum spin liquid states in frustrated Rydberg lattices. Our results highlight the broad applicability of this protocol, providing a practical framework for reliable many-body state preparation in present-day quantum platforms.

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Introduction—High-fidelity preparation of many-body quantum states is essential for advancing quantum technologies, with applications ranging from quantum computation and simulation [1–7] to precision sensing [8–10]. Analog approaches provide a natural and well-established route via continuous steering of the system through time-dependent control fields [11–15]. However, their effectiveness is often limited by experimental imperfections, such as static perturbations that alter the system dynamics in an uncontrolled fashion, thereby reducing the fidelity with the target state [16]. This motivates the development of control protocols that not only enable accurate state preparation in an ideal setting, but also ensure robustness against such imperfections under realistic conditions.

Given a time-dependent Hamiltonian $\hat{H}_0[s(t)]$, where $s(t)$ is a control parameter, we focus on the problem of preparing a strongly correlated ground state at $s(T) = s_T$ by evolving the system from a trivial ground state at $s(0) = s_0$. The buildup of quantum correlations in the target state typically requires crossing a continuous phase transition, a paradigm at the core of recent analog quantum simulation experiments [17–23]. Our goal is to devise

control protocols that are robust against static perturbations $\epsilon\hat{V}$ (cf. Fig. 1), i.e., insensitive to small variations in ϵ ,

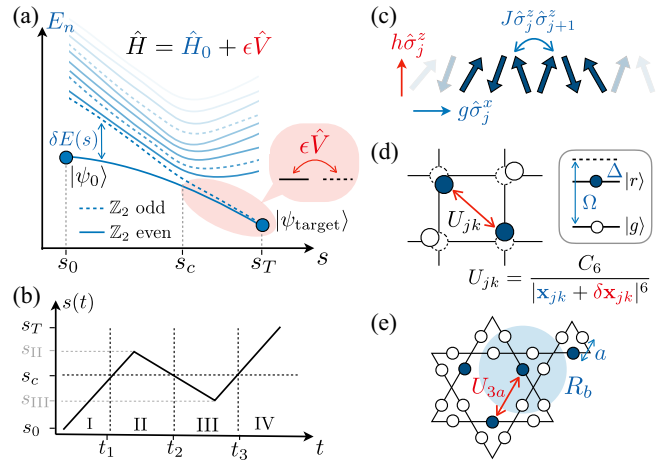


FIG. 1. Many-body adiabatic echo protocol. (a) Spectrum of a many-body Hamiltonian $\hat{H}_0(s)$ with $\mathbb{Z}_2$ symmetry breaking. The initial and target states lie in the trivial and ordered phases, separated by a critical point s_c . Solid (dashed) lines indicate eigenenergies in the even (odd) $\mathbb{Z}_2$ sectors. A perturbation $\hat{V}$ breaks the symmetry and couples these sectors, affecting the dynamics for $s > s_c$. (b) Illustration of the adiabatic echo protocol. Contributions to the infidelity $\mathcal{I}$ from regions I and III vanish, while those from regions II and IV destructively interfere. (c)–(e) Applications where the echo protocol mitigates the effect of the perturbation $\hat{V}$: (c) Ising chain with transverse and longitudinal fields, (d) square lattice Rydberg array with positional disorder, and (e) ruby lattice Rydberg array with long-range interactions.

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which is fixed during the evolution but may vary across experimental runs.

To this end, we introduce a generic preparation strategy that emerges from optimal control [24] and can be implemented across a broad class of many-body Hamiltonians. We dub this protocol *adiabatic echo protocol*, as it operates by dynamically “echoing out” the effect of static perturbations during an adiabatic evolution. After outlining the general framework in which the protocol is applicable, we develop an analytical description of the mechanism behind its robustness and validate it numerically in a concrete many-body setting. We then show how it naturally arises in several interacting models using the gradient ascent pulse engineering (GRAPE) algorithm [25]—a method widely used in few-qubit gate design [26–30] but rarely applied in the context of many-body state preparation. In particular, we showcase its versatility in various contexts: the preparation of Greenberger-Horne-Zeilinger (GHZ) states via crossing a symmetry-breaking phase transition in the quantum Ising chain [31] and in square Rydberg lattices [19] and the generation of quantum spin liquid states via crossing a topological phase transition in frustrated Rydberg arrays [22]. We further confirm the benefits of the adiabatic echo protocol in a companion experimental work by preparing GHZ states in Rydberg atom ladders [32].

Adiabatic echo protocol—Dynamical preparation of many-body states often relies on the adiabatic theorem, which ensures that a system governed by a Hamiltonian $\hat{H}_0[s(t)]$ remains close to its instantaneous ground state when $s(t)$ is varied slowly in time, relative to the inverse energy gap [15]. Therefore, driving the system from an easy-to-prepare ground state $|\psi_0\rangle$ at s_0 to an entangled target ground state $|\psi_{\text{target}}\rangle$ at s_T through a quantum phase transition, where energy gaps close polynomially with system size L , requires preparation times that scale algebraically with L to achieve a fixed fidelity [15].

In the simplest scenario, the phase transition occurs via spontaneous symmetry breaking, with a critical point s_c separating a trivial phase ($s < s_c$) from an ordered phase ($s > s_c$). The initial and final values of the control function, s_0 and s_T , lie in the trivial and ordered phases, respectively. Below, we focus on $\mathbb{Z}_2$ symmetry breaking and refer to [33] for a discussion of higher-order symmetry groups. The typical energy spectrum of a many-body Hamiltonian undergoing such a transition is sketched in Fig. 1(a): For $s < s_c$, the ground state $|E_0(s)\rangle$ is well separated from the excited states, while for $s > s_c$, spontaneous symmetry breaking leads to an exponentially small gap between the ground state and the first excited state $|E_1(s)\rangle$. We assume, without loss of generality, that these two states belong to different $\mathbb{Z}_2$ symmetry sectors and that no level crossing occurs for any value of s [43].

In the absence of perturbations, symmetry forbids mixing between $|E_0(s)\rangle$ and $|E_1(s)\rangle$. However, static

symmetry-breaking terms $\epsilon\hat{V}$ —arising from experimental imperfections or theoretical approximations—induce a nonzero coupling $V_{10}(s) = \langle E_1(s)|\hat{V}|E_0(s)\rangle$. In the adiabatic limit, couplings to other higher excited states can be neglected, and the leading-order contribution to the preparation infidelity $\mathcal{I} = 1 - |\langle\psi_{\text{target}}|\psi(T)\rangle|^2$ is given by time-dependent perturbation theory as [33]

$$\mathcal{I} = \epsilon^2 \left| \int_0^T dt V_{10}[s(t)] e^{-i \int_t^T d\tau \delta E[s(\tau)]} \right|^2, \quad (1)$$

where $\delta E(s) = E_1(s) - E_0(s)$ is the lowest energy gap.

The dominant contributions to the infidelity generically arise from the region where $\delta E(s)$ is exponentially small with the system size L ($s > s_c$), while contributions from the finite-gap region ($s < s_c$) are subleading and parametrically suppressed with L . We neglect the latter contributions in the following and refer to [33] for a formal justification of this approximation.

In a standard adiabatic protocol, $s(t)$ is monotonic and the critical point s_c is crossed only once at $t = t_c$. Because of the exponentially small gap, the phase of the integrand in Eq. (1) simplifies, and the infidelity reduces to $\mathcal{I} \simeq \epsilon^2 |V_{10}|^2 (T - t_c)^2$, where we assumed $V_{10}(s) \simeq V_{10}$. Hence, while increasing the total evolution time T improves adiabaticity, it amplifies the infidelity in the presence of symmetry-breaking perturbations.

This trade-off can be avoided by protocols that sweep across the phase transition multiple times, such as the one in Fig. 1(b), which consists of four segments. In segments I and III, $s < s_c$, while in II and IV, $s > s_c$. Neglecting the contribution to the infidelity for $s < s_c$, we have

$$\mathcal{I} \simeq \epsilon^2 \left| e^{-i\alpha} \int_{t_1}^{t_2} dt V_{10}[s(t)] + \int_{t_3}^T dt V_{10}[s(t)] \right|^2, \quad (2)$$

where $\alpha = \int_{t_2}^{t_3} d\tau \delta E[s(\tau)]$. By designing the protocol $s(t)$ such that $\alpha = \pi$ and $\int_{t_1}^{t_2} dt V_{10} = \int_{t_3}^T dt V_{10}$, these two amplitudes interfere destructively, thereby eliminating preparation infidelities at leading order in ϵ . We define a protocol satisfying these two conditions as an *adiabatic echo protocol*, where the name stems from its analogy with spin echo sequences [44]. While in that context dynamical decoupling cancels coherent errors at the Hamiltonian level [45], in our many-body setting the cancellation occurs in the instantaneous adiabatic eigenbasis.

The analytical argument above identifies a family of control pulses with a universal structure that ensures robustness against static perturbations but does not uniquely fix the control profile. In what follows, we consider concrete model Hamiltonians for the preparation of selected many-body target states and show how adiabatic echo protocols can be designed in practice. We first

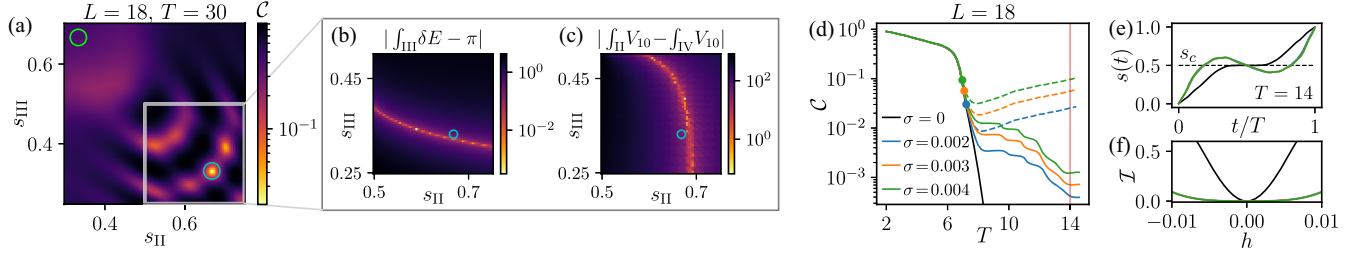


FIG. 2. Adiabatic echo in the quantum Ising chain. (a) Average infidelity Eq. (4) for the protocol in Fig. 1(b) in the $(s_{\text{II}}, s_{\text{III}})$ plane, for $\sigma = 0.003$, $N_s = 30$, and $T = 30$. The light blue circle indicates the optimal solution, while the green circle denotes the standard linear ramp. (b),(c) Interference conditions defining the adiabatic echo protocol in the $(s_{\text{II}}, s_{\text{III}})$ plane; the minimum in (a) (light blue circle) lies close to their intersection. (d) Optimal cost function $\mathcal{C}$ versus total preparation time T for different σ ($N_s = 30$). Dashed lines show the averaged infidelity evaluated on time-optimal protocols obtained for $\sigma = 0$. Colored dots indicate the time T^* beyond which the $\sigma > 0$ optimization deviates from the time-optimal solution. (e) Optimized control profiles $s(t)$ for $T = 14$ [vertical red line in (d)] and different σ . (f) Preparation infidelity for the profiles in (e) as a function of the longitudinal field h .

demonstrate that, within a simple manifold spanned by two parameters that set the relative phase and amplitude for destructive interference, the adiabatic echo protocol coincides with the optimal robust solution. We then use GRAPE to show how adiabatic echo protocols emerge in unbiased optimal control across multiple interacting models and target states. We further show in [33] that these results are independent of the optimal control method employed.

GHZ state preparation in the quantum Ising chain—The quantum Ising chain Hamiltonian reads [46]

$$\hat{H}_{\text{Ising}} = -J \sum_j \hat{\sigma}_j^z \hat{\sigma}_{j+1}^z - g \sum_j \hat{\sigma}_j^x + h \sum_j \hat{\sigma}_j^z, \quad (3)$$

where $\hat{\sigma}_j^{x/z}$ are Pauli matrices and periodic boundary conditions are assumed. We parametrize the couplings J and g with a single control parameter s as $J = J_0 \sin(\pi s/2)$ and $g = J_0 \cos(\pi s/2)$ and fix the overall energy scale by setting $J_0 = 1$. For $h = 0$, $\hat{H}_{\text{Ising}}$ represents the unperturbed Hamiltonian $\hat{H}_0$, which has a $\mathbb{Z}_2$ symmetry generated by $\prod_j \hat{\sigma}_j^z$. The critical point $s_c = 1/2$ separates the disordered phase ($s < s_c$) from the ordered one ($s > s_c$). The last term in Eq. (3) breaks the symmetry and plays the role of the static perturbation $\hat{V}$.

The dynamical preparation is carried out by evolving the system under a time-dependent $s(t)$, starting from the product state $|+\cdots+\rangle$, i.e., the ground state of $\hat{H}_0(s=0)$. The target state is $|\psi_{\text{GHZ}}^+\rangle = \frac{1}{\sqrt{2}}(|\uparrow\uparrow\cdots\uparrow\rangle + |\downarrow\downarrow\cdots\downarrow\rangle)$, i.e., the ground state of $\hat{H}_0(s=1)$, exactly degenerate with $|\psi_{\text{GHZ}}^-\rangle = \frac{1}{\sqrt{2}}(|\uparrow\uparrow\cdots\uparrow\rangle - |\downarrow\downarrow\cdots\downarrow\rangle)$. Both the initial and target states belong to the even $\mathbb{Z}_2$ symmetry sector. The static field h couples these two sectors, inducing leakage from $|\psi_{\text{GHZ}}^+\rangle$ to $|\psi_{\text{GHZ}}^-\rangle$ when $s(t) \gtrsim s_c$, resulting in a loss of fidelity for the prepared GHZ state.

We begin by considering a preparation protocol of the form depicted in Fig. 1(b), i.e., a piecewise linear $s(t)$ satisfying $s(0) = 0$, $s(T) = 1$, $s(T/3) = s_{\text{II}}$, and

$s(2T/3) = s_{\text{III}}$, where s_{II} and s_{III} are two control parameters. To quantify robustness against symmetry breaking, we use the preparation infidelity averaged over N_s values of the perturbation strength ϵ_ℓ :

$$\mathcal{C} = 1 - \frac{1}{N_s} \sum_{\ell=1}^{N_s} \left| \langle \psi_{\text{target}} | \prod_{j=0}^{N-1} e^{-idt[\hat{H}_0(s_j) + \epsilon_\ell \hat{V}]} | \psi_0 \rangle \right|^2, \quad (4)$$

where $s_j = s(t_j)$, $t_j = (j + \frac{1}{2})dt$ with $j = 0, \dots, N-1$, $dt = T/N$ with T the total evolution time, and ϵ is the longitudinal field h in Eq. (3), sampled uniformly in the interval $[-\sigma, \sigma]$. The case $\sigma = 0$ ($N_s = 1$) corresponds to the unperturbed control problem.

In Fig. 2(a), we plot $\mathcal{C}$ in the $(s_{\text{II}}, s_{\text{III}})$ plane for $L = 18$ spins and $T = 30$, revealing a clear global minimum for $(s_{\text{II}}, s_{\text{III}}) \simeq (0.67, 0.33)$. Remarkably, this point lies at the intersection of the two interference conditions defining the adiabatic echo protocol: the accumulation of a dynamical phase equal to π in region III [Fig. 2(b)] and equal symmetry-breaking transition amplitudes in regions II and IV [Fig. 2(c)].

We now apply GRAPE to optimize the function $s(t)$ in the full control space for preparing $|\psi_{\text{GHZ}}^+\rangle$. The variational parameters in this case are the N time-discretized values s_j . The optimization is performed using standard routines that rely on the gradient $\partial \mathcal{C} / \partial s_j$, which can be efficiently computed numerically [33]. Moreover, we add a term to $\mathcal{C}$ that penalizes sharp variations of the optimal control function between consecutive time steps [33]. In Fig. 2(d), we plot the optimal $\mathcal{C}$ as a function of the preparation time T for various values of the maximal perturbation strength σ , for $L = 18$. The system-size scaling of these results is discussed in [33]. The solid black line represents purely time-optimal protocols obtained for $\sigma = 0$, while the dashed lines represent the corresponding averaged infidelities Eq. (4) evaluated for $\sigma > 0$.

For small T , the infidelity is dominated by nonadiabatic transitions within the even $\mathbb{Z}_2$ sector, and the optimization

for $\sigma > 0$ leads to the same results as for $\sigma = 0$. As T increases and the evolution turns adiabatic, the leakage to the odd $\mathbb{Z}_2$ sector due to the perturbation becomes the main source of infidelity. In this regime, the optimal solution for $\sigma > 0$ deviates from the $\sigma = 0$ case and takes the form of an echo protocol. This is illustrated in Fig. 2(e), where a typical time-optimal solution (black line), in which $s(t)$ crosses s_c only once, is compared with an echolike solution which resembles a smoothed version of Fig. 1(b). Notably, the optimal control solution is effective at shorter preparation times than those considered in Fig. 2(a).

Although the evolution time T^* at which the echo protocol becomes beneficial decreases as σ increases [cf. the colored markers in Fig. 2(d)], the control profile $s(t)$ appears to be insensitive to the value of σ . In Fig. 2(f), we plot the infidelity as a function of the longitudinal field h for the echo protocol (colored lines), demonstrating its significantly enhanced robustness compared to the time-optimal case (black line). We stress that this property is not limited to the particular perturbation considered here but extends to more general symmetry-breaking perturbations and remains effective even for time-dependent noise in the quasistatic limit [33].

GHZ state preparation in Rydberg atom arrays—As a second, experimentally relevant example, we turn to Rydberg atom arrays [47], a paradigmatic quantum platform that naturally features coherent control and strong interactions with static imperfections [48]. Each atom is modeled as a two-level system with ground state $|g\rangle$ and Rydberg state $|r\rangle$, and the dynamics is described by the Hamiltonian

$$\hat{H}_{\text{Ryd}} = \sum_j \left(\frac{\Omega}{2} \hat{\sigma}_j^x - \Delta \hat{n}_j \right) + \sum_{j < k} U_{jk} \hat{n}_j \hat{n}_k, \quad (5)$$

where $\hat{\sigma}_j^x = |g\rangle_j \langle r| + |r\rangle_j \langle g|$ and $\hat{n}_j = |r\rangle_j \langle r|$. Here, Ω is the Rabi frequency, Δ is the detuning, and $U_{jk} \sim 1/|\mathbf{x}_j - \mathbf{x}_k|^6$ is the van der Waals interaction between atoms at positions $\mathbf{x}_j$ and $\mathbf{x}_k$. The blockade radius R_b is defined as the characteristic length scale at which $U_{R_b} = \Omega$, marking the distance below which simultaneous Rydberg excitations are suppressed [49–51].

We consider arrays of 16 atoms on a square lattice [20,52] with open boundary conditions, since this geometry is readily realizable experimentally and avoids the need for additional local control required in open chains [19]. Similarly to the Ising model, $\hat{H}_{\text{Ryd}}$ possesses a $\mathbb{Z}_2$ symmetry encoded as a lattice reflection exchanging the two sublattices and, for suitable R_b , exhibits a disordered and an ordered phases, separated by a critical point. Below, we set $R_b = 1.15a$, where a is the lattice spacing.

We aim to prepare $|\psi_{\text{GHZ}}^+\rangle = \frac{1}{\sqrt{2}}(|g\rangle_A |r\rangle_B + |r\rangle_A |g\rangle_B)$, where $|g\rangle_A |r\rangle_B$ denotes the configuration with all atoms in sublattice A (B) in state $|g\rangle$ ($|r\rangle$), and vice versa.

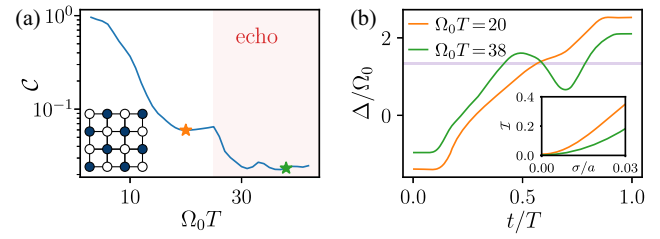


FIG. 3. GHZ state preparation in Rydberg arrays. (a) Optimal cost function Eq. (4) vs total preparation time, for $N_s = 30$ and $\sigma = 0.01a$. In the shaded red region, the optimal protocols exhibit the features of the echo mechanism. (b) Optimal profiles for $\Omega_0 T = 20$ (standard adiabatic) and $\Omega_0 T = 38$ (adiabatic echo). The insets compare the robustness of the two against the strength of the positional disorder. The shaded purple region marks the finite-size critical regime, estimated from two independent ground state diagnostics [33].

We initialize the system in the product state $|\psi_0\rangle = |gg \cdots g\rangle$ and evolve it under the time-dependent Rydberg Hamiltonian in Eq. (5), with $\Omega(t) = \Omega_0 w(t)$, where $w(t)$ is a window function satisfying $w(0) = w(T) = 0$ [33]. We use the detuning $\Delta(t)$ as the control function. Provided that $\Delta(0) < 0$, $\Delta(T) > 0$ and that $\mathbb{Z}_2$ symmetry is not explicitly broken, $|\psi_0\rangle$ is the ground state for $t = 0$, while $|\psi_{\text{GHZ}}^+\rangle$ is the ground state for $t \rightarrow T$.

A relevant source of symmetry breaking in this setting arises from static but spatially inhomogeneous atomic displacements, which vary from shot to shot and explicitly break the $\mathbb{Z}_2$ lattice reflection symmetry [cf. Fig. 1(d)]. These experimental imperfections are particularly significant due to the diagonal interaction $U_{\sqrt{2}a} \approx 0.3\Omega$, which is comparable to the Rabi frequency.

To identify the echo protocol for robust GHZ state preparation, we apply GRAPE to the cost function Eq. (4) (cf. Fig. 3), where each sample consists of one displacement $\delta \mathbf{x}_j$ per atom, independently drawn from a Gaussian distribution with zero mean and standard deviation $\sigma = 0.01a$. As in the Ising model, the optimal control solution differs qualitatively from a standard adiabatic protocol beyond a timescale $\Omega_0 T^* \simeq 20$ [colored region in Fig. 3(a)]. The corresponding optimal profiles are shown in Fig. 3(b): While for $T < T^*$ the solution features a monotonic ramp-up of $\Delta(t)$, for $T > T^*$ the optimal protocol exhibits the characteristic structure of the echo scheme, featuring three sweeps across the phase transition [cf. Fig. 3(b)]. The inset in Fig. 3(b) further highlights the enhanced robustness of the echo protocol, which was experimentally validated in a parallel study on Rydberg ladders [32,33], enabling the preparation of GHZ states nearly twice as large as those achievable with standard adiabatic protocols.

Quantum spin liquid state preparation—As a final example, we consider Rydberg atoms on the ruby lattice and set $R_b = 2.3a$ [22,53]. Similarly to the GHZ preparation discussed above, we fix $\Omega(t) = \Omega_0 w(t)$, with

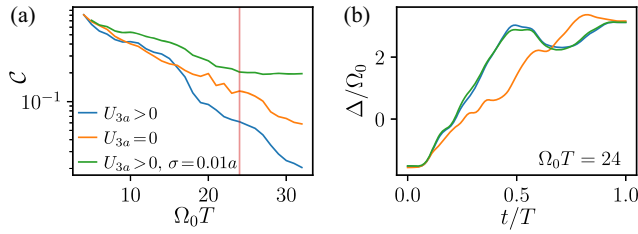


FIG. 4. RVB state preparation in Rydberg arrays. (a) Optimal RVB preparation fidelity vs total preparation time on the ruby lattice with 24 atoms, truncating ($U_{3a} = 0$) and including ($U_{3a} > 0$) the van der Waals potential. For $U_{3a} > 0$ the optimization is also performed including positional disorder via the cost function Eq. (4), averaged over $N_s = 30$ Gaussian distributed samples (green line). (b) Optimal control profiles for $\Omega_0 T = 24$ [vertical red line in (a)]. The optimal protocol exhibits the structure of the adiabatic echo only when $U_{3a} > 0$.

$w(0) = w(T) = 0$ [33], take $\Delta(t)$ as the control function, with $\Delta(0) < 0$ and $\Delta(T) > 0$, and initialize the system in the product state $|gg \cdots g\rangle$. For $t = T$, and with the van der Waals potential truncated as $U_R = 0$ for $R > \sqrt{7}a$, the classical ground states correspond to all fully packed dimer coverings of the kagome lattice [53]. Their equal-weight superposition—the resonating valence bond (RVB) state [54]—is a $\mathbb{Z}_2$ quantum spin liquid [55,56] and is the target of our preparation. Previous works have demonstrated that this state can be prepared with high fidelity in this setting [57–59]. However, the inclusion of longer-range interaction tails U_{3a} lifts the degeneracy of the RVB components in the classical limit, imprinting relative phases during the evolution and thereby degrading the fidelity of the prepared state [57,60].

While this perturbation does not explicitly break any microscopic symmetry of the Rydberg Hamiltonian on the ruby lattice, it plays a role analogous to the static perturbation $\hat{V}$ in the GHZ preparation discussed earlier, with the two GHZ components replaced by exponentially many RVB configurations. To determine the optimal protocol for preparing the RVB state in this setting, we apply GRAPE to the full Hamiltonian equation (5) on a periodic cluster of 24 atoms, including all van der Waals interaction tails. Remarkably, as shown in Fig. 4(b), the resulting optimal control profile displays the characteristic nonmonotonic shape of the adiabatic echo. This solution is also robust, as the optimal protocol remains essentially unchanged when positional disorder is included in the optimization [cf. blue and green lines in Fig. 4(b)]. In contrast, when the interaction is truncated at distance $R = \sqrt{7}a$, the optimal solution is monotonic [cf. orange line in Fig. 4(b)], confirming that the echo protocol arises specifically to mitigate the effect of the long-range tails.

Conclusions—We introduced the adiabatic echo protocol as a general approach to many-body quantum state preparation in the presence of static perturbations. Its key feature is the suppression of leading-order errors via dynamically

engineered interference, which we analyzed both analytically and numerically. Going beyond standard approaches to many-body optimal control, which rely on restrictive parametrizations of the control fields [19,33,61,62], we applied GRAPE to various interacting models and uncovered robust solutions that naturally exhibit the echo structure.

This Letter demonstrates that unbiased optimal control methods such as GRAPE can be effectively applied to strongly interacting quantum systems, highlighting their effectiveness in this complex setting. The adiabatic echo protocol introduced here offers a novel approach to robust quantum state preparation, with potential applicability across a wide range of platforms beyond those explicitly considered. Taken together, these findings open new directions for theoretical and numerical exploration of robust many-body quantum dynamics.

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Data availability—The data that support the findings of this article are openly available [63].

- [1] M. A. Nielsen and I. L. Chuang, *Quantum Computation and Quantum Information*, 10th ed. (Cambridge University Press, Cambridge, England, 2000).
- [2] E. Farhi, J. Goldstone, S. Gutmann, and M. Sipser, Quantum computation by adiabatic evolution, [arXiv:quant-ph/0001106](https://arxiv.org/abs/quant-ph/0001106).
- [3] R. Raussendorf and H. J. Briegel, A one-way quantum computer, *Phys. Rev. Lett.* **86**, 5188 (2001).
- [4] J. I. Cirac and P. Zoller, Goals and opportunities in quantum simulation, *Nat. Phys.* **8**, 264 (2012).
- [5] I. M. Georgescu, S. Ashhab, and F. Nori, Quantum simulation, *Rev. Mod. Phys.* **86**, 153 (2014).
- [6] C. Gross and I. Bloch, Quantum simulations with ultracold atoms in optical lattices, *Science* **357**, 995 (2017).
- [7] E. Altman *et al.*, Quantum simulators: Architectures and opportunities, *PRX Quantum* **2**, 017003 (2021).
- [8] V. Giovannetti, S. Lloyd, and L. Maccone, Advances in quantum metrology, *Nat. Photonics* **5**, 222 (2011).
- [9] C. L. Degen, F. Reinhard, and P. Cappellaro, Quantum sensing, *Rev. Mod. Phys.* **89**, 035002 (2017).

- [10] J. Ye and P. Zoller, Essay: Quantum sensing with atomic, molecular, and optical platforms for fundamental physics, *Phys. Rev. Lett.* **132**, 190001 (2024).
- [11] J. Dziarmaga, Dynamics of a quantum phase transition and relaxation to a steady state, *Adv. Phys.* **59**, 1063 (2010).
- [12] A. Polkovnikov, K. Sengupta, A. Silva, and M. Vengalattore, Colloquium: Nonequilibrium dynamics of closed interacting quantum systems, *Rev. Mod. Phys.* **83**, 863 (2011).
- [13] M. Kolodrubetz, D. Sels, P. Mehta, and A. Polkovnikov, Geometry and non-adiabatic response in quantum and classical systems, *Phys. Rep.* **697**, 1 (2017).
- [14] D. Sels and A. Polkovnikov, Minimizing irreversible losses in quantum systems by local counterdiabatic driving, *Proc. Natl. Acad. Sci. U.S.A.* **114**, E3909 (2017).
- [15] T. Albash and D. A. Lidar, Adiabatic quantum computation, *Rev. Mod. Phys.* **90**, 015002 (2018).
- [16] P. Hauke, F. M. Cucchietti, L. Tagliacozzo, I. Deutsch, and M. Lewenstein, Can one trust quantum simulators?, *Rep. Prog. Phys.* **75**, 082401 (2012).
- [17] H. Bernien, S. Schwartz, A. Keesling, H. Levine, A. Omran, H. Pichler, S. Choi, A. S. Zibrov, M. Endres, M. Greiner, V. Vuletić, and M. D. Lukin, Probing many-body dynamics on a 51-atom quantum simulator, *Nature (London)* **551**, 579–584 (2017).
- [18] A. Keesling, A. Omran, H. Levine, H. Bernien, H. Pichler, S. Choi, R. Samajdar, S. Schwartz, P. Silvi, S. Sachdev, P. Zoller, M. Endres, M. Greiner, V. Vuletić, and M. D. Lukin, Quantum Kibble–Zurek mechanism and critical dynamics on a programmable Rydberg simulator, *Nature (London)* **568**, 207 (2019).
- [19] A. Omran, H. Levine, A. Keesling, G. Semeghini, T. T. Wang, S. Ebadi, H. Bernien, A. S. Zibrov, H. Pichler, S. Choi, J. Cui, M. Rossignolo, P. Rembold, S. Montangero, T. Calarco, M. Endres, M. Greiner, V. Vuletić, and M. D. Lukin, Generation and manipulation of Schrödinger cat states in Rydberg atom arrays, *Science* **365**, 570 (2019).
- [20] S. Ebadi, T. T. Wang, H. Levine, A. Keesling, G. Semeghini, A. Omran, D. Bluvstein, R. Samajdar, H. Pichler, W. W. Ho, S. Choi, S. Sachdev, M. Greiner, V. Vuletić, and M. D. Lukin, Quantum phases of matter on a 256-atom programmable quantum simulator, *Nature (London)* **595**, 227 (2021).
- [21] P. Scholl, M. Schuler, H. J. Williams, A. A. Eberharter, D. Barredo, K.-N. Schymik, V. Lienhard, L.-P. Henry, T. C. Lang, T. Lahaye, A. M. Läuchli, and A. Browaeys, Quantum simulation of 2D antiferromagnets with hundreds of Rydberg atoms, *Nature (London)* **595**, 233 (2021).
- [22] G. Semeghini, H. Levine, A. Keesling, S. Ebadi, T. T. Wang, D. Bluvstein, R. Verresen, H. Pichler, M. Kalinowski, R. Samajdar, A. Omran, S. Sachdev, A. Vishwanath, M. Greiner, V. Vuletić, and M. D. Lukin, Probing topological spin liquids on a programmable quantum simulator, *Science* **374**, 1242 (2021).
- [23] S.-A. Guo, Y.-K. Wu, J. Ye, L. Zhang, W.-Q. Lian, R. Yao, Y. Wang, R.-Y. Yan, Y.-J. Yi, Y.-L. Xu, B.-W. Li, Y.-H. Hou, Y.-Z. Xu, W.-X. Guo, C. Zhang, B.-X. Qi, Z.-C. Zhou, L. He, and L.-M. Duan, A site-resolved two-dimensional quantum simulator with hundreds of trapped ions, *Nature (London)* **630**, 613–618 (2024).
- [24] S. J. Glaser, U. Boscain, T. Calarco, C. P. Koch, W. Köckenberger, R. Kosloff, I. Kuprov, B. Luy, S. Schirmer, T. Schulte-Herbrüggen, D. Sugny, and F. K. Wilhelm, Training Schrödinger’s cat: Quantum optimal control. Strategic report on current status, visions and goals for research in Europe, *Eur. Phys. J. D* **69**, 279 (2015).
- [25] N. Khaneja, T. Reiss, C. Kehlet, T. Schulte-Herbrüggen, and S. J. Glaser, Optimal control of coupled spin dynamics: Design of NMR pulse sequences by gradient ascent algorithms, *J. Magn. Reson.* **172**, 296 (2005).
- [26] A. Garon, S. J. Glaser, and D. Sugny, Time-optimal control of SU(2) quantum operations, *Phys. Rev. A* **88**, 043422 (2013).
- [27] D. J. Egger and F. K. Wilhelm, Optimized controlled-Z gates for two superconducting qubits coupled through a resonator, *Supercond. Sci. Technol.* **27**, 014001 (2014).
- [28] F. Dolde, I. Jakobi, B. Naydenov, N. Zhao, S. Pezzagna, C. Trautmann, J. Meijer, P. Neumann, F. Jelezko, and J. Wrachtrup, High-fidelity spin entanglement using optimal control, *Nat. Commun.* **5**, 3371 (2014).
- [29] S. Jandura and G. Pupillo, Time-optimal two- and three-qubit gates for Rydberg atoms, *Quantum* **6**, 712 (2022).
- [30] G. Giudici, S. Veroni, G. Giudice, H. Pichler, and J. Zeiher, Fast entangling gates for Rydberg atoms via resonant dipole-dipole interaction, *PRX Quantum* **6**, 030308 (2025).
- [31] X. Wang, A. Bayat, S. Bose, and S. G. Schirmer, Global control methods for GHZ state generation on 1-D Ising chain, *Phys. Rev. A* **82**, 012330 (2010).
- [32] A. Senoo, A. Baumgärtner, J. W. Lis, G. M. Vaidya, Z. Zeng, G. Giudici, H. Pichler, and A. M. Kaufman, High-fidelity entanglement and coherent multi-qubit mapping in an atom array, *arXiv:2506.13632*.
- [33] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/z2tj-cwzb> for perturbative analyses of the adiabatic echo protocol, robustness to additional perturbations, details of the optimal-control methods, finite-size scaling of GHZ state preparation, and additional results for Rydberg atom arrays and higher-symmetry cat states, which includes Refs. [17,25,32,34–42].
- [34] A. H. Al-Mohy and N. J. Higham, Computing the Fréchet derivative of the matrix exponential, with an application to condition number estimation, *SIAM J. Matrix Anal. Appl.* **30**, 1639 (2009).
- [35] J. Haegeman, C. Lubich, I. Oseledets, B. Vandereycken, and F. Verstraete, Unifying time evolution and optimization with matrix product states, *Phys. Rev. B* **94**, 165116 (2016).
- [36] P. Virtanen *et al.*, SciPy 1.0: Fundamental algorithms for scientific computing in Python, *Nat. Methods* **17**, 261 (2020).
- [37] S. Machnes, E. Assémat, D. Tannor, and F. K. Wilhelm, Tunable, flexible, and efficient optimization of control pulses for practical qubits, *Phys. Rev. Lett.* **120**, 150401 (2018).
- [38] L. Eck and P. Fendley, Critical lines and ordered phases in a Rydberg-blockade ladder, *Phys. Rev. B* **108**, 125135 (2023).
- [39] J. Soto-Garcia and N. Chepiga, Numerical investigation of quantum phases and phase transitions in a two-leg ladder of Rydberg atoms, *Phys. Rev. Res.* **7**, 013215 (2025).
- [40] J. Zhang, S. H. Cantú, F. Liu, A. Bylinskii, B. Braverman, F. Huber, J. Amato-Grill, A. Lukin, N. Gemelke, A. Keesling, S.-T. Wang, Y. Meurice, and S.-W. Tsai, Probing quantum

- floating phases in Rydberg atom arrays, *Nat. Commun.* **16**, 712 (2025).
- [41] P. Zanardi and N. Paunković, Ground state overlap and quantum phase transitions, *Phys. Rev. E* **74**, 031123 (2006).
- [42] S.-J. Gu, Fidelity approach to quantum phase transitions, *Int. J. Mod. Phys. B* **24**, 4371 (2010).
- [43] Level crossings between the two symmetry sectors occurring at isolated points would lead only to subleading contributions to the preparation infidelity induced by symmetry breaking. As shown in [33], the dominant contribution instead arises from the ordered phase, where hybridization between the ground and first excited states persists over a time interval that scales extensively with system size.
- [44] E. L. Hahn, Spin echoes, *Phys. Rev.* **80**, 580 (1950).
- [45] R. R. Ernst, G. Bodenhausen, and A. Wokaun, *Principles of Nuclear Magnetic Resonance in One and Two Dimensions* (Oxford University Press, New York, 1990).
- [46] S. Suzuki, J.-i. Inoue, and B. K. Chakrabarti, *Quantum Ising Phases and Transitions in Transverse Ising Models*, Lecture Notes in Physics Vol. 862 (Springer, Berlin, Heidelberg, 2013).
- [47] A. Browaeys and T. Lahaye, Many-body physics with individually controlled Rydberg atoms, *Nat. Phys.* **16**, 132 (2020).
- [48] J. Wurtz, A. Bylinskii, B. Braverman, J. Amato-Grill, S. H. Cantu, F. Huber, A. Lukin, F. Liu, P. Weinberg, J. Long, S.-T. Wang, N. Gemelke, and A. Keesling, Aquila: QuEra's 256-qubit neutral-atom quantum computer, [arXiv:2306.11727](https://arxiv.org/abs/2306.11727).
- [49] D. Jaksch, J. I. Cirac, P. Zoller, S. L. Rolston, R. Côté, and M. D. Lukin, Fast quantum gates for neutral atoms, *Phys. Rev. Lett.* **85**, 2208 (2000).
- [50] M. D. Lukin, M. Fleischhauer, R. Cote, L. M. Duan, D. Jaksch, J. I. Cirac, and P. Zoller, Dipole blockade and quantum information processing in mesoscopic atomic ensembles, *Phys. Rev. Lett.* **87**, 037901 (2001).
- [51] A. Gaëtan, Y. Miroshnychenko, T. Wilk, A. Chotia, M. Viteau, D. Comparat, P. Pillet, A. Browaeys, and P. Grangier, Observation of collective excitation of two individual atoms in the Rydberg blockade regime, *Nat. Phys.* **5**, 115 (2009).
- [52] R. Samajdar, W. W. Ho, H. Pichler, M. D. Lukin, and S. Sachdev, Complex density wave orders and quantum phase transitions in a model of square-lattice Rydberg atom arrays, *Phys. Rev. Lett.* **124**, 103601 (2020).
- [53] R. Verresen, M. D. Lukin, and A. Vishwanath, Prediction of Toric code topological order from Rydberg blockade, *Phys. Rev. X* **11**, 031005 (2021).
- [54] P. W. Anderson, Resonating valence bonds: A new kind of insulator?, *Mater. Res. Bull.* **8**, 153 (1973).
- [55] G. Misguich, D. Serban, and V. Pasquier, Quantum dimer model on the kagome lattice: Solvable dimer-liquid and Ising gauge theory, *Phys. Rev. Lett.* **89**, 137202 (2002).
- [56] R. Moessner and K. S. Raman, Quantum dimer models, in *Introduction to Frustrated Magnetism*, edited by C. Lacroix, P. Mendels, and F. Mila (Springer, New York, 2011), pp. 437–479.
- [57] G. Giudici, M. D. Lukin, and H. Pichler, Dynamical preparation of quantum spin liquids in Rydberg atom arrays, *Phys. Rev. Lett.* **129**, 090401 (2022).
- [58] R. Sahay, A. Vishwanath, and R. Verresen, Quantum spin puddles and lakes: NISQ-Era spin liquids from non-equilibrium dynamics, [arXiv:2211.01381](https://arxiv.org/abs/2211.01381).
- [59] N. O. Gjonbalaj, R. Sahay, and S. F. Yelin, Shortcuts to analog preparation of non-equilibrium quantum lakes, [arXiv:2502.03518](https://arxiv.org/abs/2502.03518).
- [60] Z. Zeng, G. Giudici, and H. Pichler, Quantum dimer models with Rydberg gadgets, *Phys. Rev. Res.* **7**, L012006 (2025).
- [61] T. Caneva, T. Calarco, and S. Montangero, Chopped random-basis quantum optimization, *Phys. Rev. A* **84**, 022326 (2011).
- [62] J. Cui, R. van Bijnen, T. Pohl, S. Montangero, and T. Calarco, Optimal control of Rydberg lattice gases, *Quantum Sci. Technol.* **2**, 035006 (2017).
- [63] Z. Zeng, G. Giudici, A. Senoo, A. Baumgärtner, A. Kaufmann, and H. Pichler, Adiabatic echo protocols for robust quantum many-body state preparation, 2026, [10.5281/zenodo.18865750](https://zenodo.org/record/18865750).