

Controlled-Z gates with giant atoms in structured waveguides

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Giant atoms are quantum emitters coupled to waveguides at multiple, spatially separated points, enabling interference effects that fundamentally change their light-matter interactions. A notable consequence of the interference is the emergence of decoherence-free interaction (DFI), which allows coherent excitation exchange between giant atoms via the waveguide without radiative loss. Leveraging DFI offers a promising route to implementing two-qubit quantum gates without the need for additional resources, positioning giant atoms as a versatile platform for scalable universal quantum simulators. However, existing work has focused primarily on continuous, Markovian waveguides; in structured waveguides, where non-Markovian effects become significant, only iSWAP gates have been explored. To address this gap, we introduce and analyze a protocol for implementing controlled-z (CZ) gates with giant atoms in structured waveguides. We first show that while a minimal two-point coupling scheme supports DFI, it also exhibits strong non-Markovian effects that substantially degrade gate fidelity. To overcome this limitation, we propose an extended design featuring a third coupling point. This configuration suppresses non-Markovian effects and enables CZ gates with average gate fidelities up to 97.7% (assuming typical values for experimental imperfections). Our results broaden the accessible gate set for giant atoms in structured waveguides to include both iSWAP and CZ gates, advancing these systems as a pathway toward universal quantum simulators operating in non-Markovian environments.

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I. INTRODUCTION

Universal quantum simulators are powerful tools for exploring complex open quantum many-body dynamics [1–3]. Yet, most existing quantum simulators are optimized for implementing only specific two-qubit gates [3–5]. Although this can be sufficient for realizing a universal gate set together with single-qubit gates, it is generally desirable to have more flexibility regarding which two-qubit gates one can implement, e.g., to reduce the depth of quantum circuits. However, enabling such flexibility typically requires additional coupling elements [6–19] or complicated drive schemes [15,20–24] that limit the feasibility of building large-scale quantum simulators.

Here, we explore how an extended set of two-qubit gates can be realized using giant atoms (GAs) [25–29], which have recently emerged as a versatile platform for scalable universal quantum simulators that can handle open quantum systems [30–33]. Unlike conventional small atoms that interact locally, GAs couple to waveguides at multiple, spatially separated points. This nonlocal coupling results in interference-driven light-matter dynamics, giving rise to qual-

itatively different phenomena such as frequency-dependent decay [27,29,34,35], directional emission [36–39], tunable interaction rates [28,29,35] between distant GAs, and much more [25,40–72]. Among the most striking consequences of this nonlocal coupling is the possibility of realizing decoherence-free interactions (DFIs), where two or more atoms exchange excitations coherently while completely suppressing radiative decay through destructive interference between coupling points [28,29]. These DFIs provide a natural mechanism for entangling gates without the need for additional parametric couplers, making GAs a versatile platform for quantum simulators [30–33].

Previous studies often focused on GAs coupled to continuous waveguides whose dispersion can be approximated as linear in the frequency regime where the atoms operate. In these setups, the Markovian approximation is typically valid and analytical treatments are tractable [30,32,33]. Additionally, in such setups, iSWAP gates have been experimentally demonstrated using superconducting qubits coupled to a microwave waveguide [29].

In contrast to continuous waveguides, structured waveguides [73–75], such as coupled-cavity arrays [39,76–78] or photonic crystals [79–81], have nonlinear dispersions with band edges and gaps in the frequency regime where the atoms operate. Exploring GAs coupled to structured waveguides thus opens up the possibility of exploiting band-edge effects, engineered dispersion, and strong non-Markovian effects [46,49,53,54,82–94]. Since the Markovian approximation does not hold, such studies require keeping track of the entire system's dynamics and have thus mainly focused on single-excitation dynamics, which only support iSWAP gates

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[53]. Establishing universal quantum simulators in GAs coupled to structured waveguides would benefit from additional entangling gates such as the controlled-z (CZ) gate, which remains to be explored as it relies on dynamics in the two-excitation sector.

In this article, we propose and analyze a protocol for implementing a CZ gate between two GAs coupled to a structured waveguide. We begin by studying the minimal two-point coupling configuration, which supports DFI, but suffers from non-Markovian effects that strongly degrade the DFI. We then demonstrate that introducing a third coupling point shifts the decoherence-free frequencies closer to the center of the photonic band of the structured waveguide, thereby mitigating these effects. For experimentally realistic parameters in superconducting qubits, we show that this design yields CZ gate fidelities up to 97.7%. These results equip GAs coupled to structured waveguides with an extended gate set involving both iSWAP and CZ gates, pushing these platforms forward toward universal quantum simulators in non-Markovian environments.

The rest of this article is organized as follows. In Sec. II, we introduce the model of two three-level GAs coupled to a structured waveguide. In Sec. III, we analyze the CZ gate protocol, beginning with the minimal two-point configuration and then extending to the improved three-point scheme. We further evaluate the resulting gate fidelity and discuss experimental feasibility. Finally, in Sec. IV, we summarize the main findings and outline their broader implications.

II. MODEL AND METHOD

A. System Hamiltonian

We consider a system of two GAs coupled to a one-dimensional structured waveguide, as shown in Fig. 1(a). The total Hamiltonian for this system is given by

$$H = H_A + H_B + H_{\text{int}}, \quad (1)$$

where H_A describes the bare GAs, H_B describes the bare bath (i.e., the waveguide), and H_{int} describes the interaction between the atoms and the bath. Since the CZ gate dynamics in our scheme involve the $|2\rangle$ level of the qubits [see Fig. 1(b)], we model each atom as a three-level system ($\hbar = 1$ here and in the rest of the article)

$$H_A = \sum_{m=1}^2 \left[(2\omega_m + \alpha_m) \frac{\sigma_m^{(2),+} \sigma_m^{(2),-}}{2} + \omega_m \sigma_m^{(1),+} \sigma_m^{(1),-} \right], \quad (2)$$

where ω_m denotes the transition frequency from the ground state to the first excited state for atom m , α_m denotes the anharmonicity of atom m , and the ladder operators are defined as

$$\sigma_m^{(j),+} |j-1\rangle_m = \sqrt{j} |j\rangle_m, \quad \sigma_m^{(j),-} |j\rangle_m = \sqrt{j} |j-1\rangle_m. \quad (3)$$

We model the structured bath as a one-dimensional array of N coupled cavities, with the center of the photonic band at the frequency ω_c :

$$H_B = \sum_{n=1}^N \omega_c a_n^\dagger a_n - J \sum_{n=1}^{N-1} (a_n a_{n+1}^\dagger + \text{H.c.}), \quad (4)$$

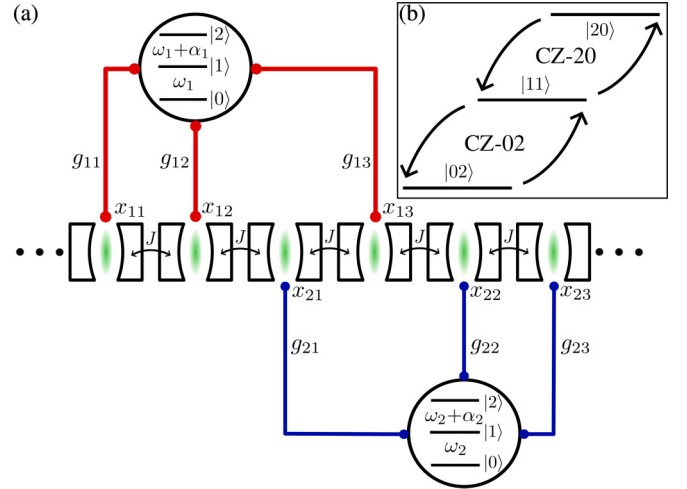


FIG. 1. Model and gate scheme. (a) Schematic of two giant atoms (GAs) coupled to a one-dimensional array of coupled cavities with nearest-neighbor hopping rate J . Each atom m is connected with coupling strength g_{mj} to the array at site x_{mj} . The atoms have distinct transition frequencies $\omega_{1(2)}$ and anharmonicities $\alpha_{1(2)}$. (b) The CZ gate in GAs can be implemented via the pathway CZ-02 (CZ-20), which coherently transfers the population of the $|11\rangle$ state to the $|02\rangle$ ($|20\rangle$) state and then brings it back, accumulating a π phase shift.

where J is the coupling rate between adjacent cavities, a_n and $a_n^\dagger$ are the annihilation and creation operators, respectively, for photons in cavity n , and H.c. denotes the Hermitian conjugate. This Hamiltonian yields the dispersion relation

$$\omega(k) = -2J \cos k + \omega_c, \quad (5)$$

where k is the wave number resulting in a cosine band of width $4J$. As ω_c is a constant, we can go to a frame rotating at that frequency in Eq. (1) such that the first term in Eq. (4) is zero and we redefine $\omega_m = \omega_m - \omega_c$ in Eq. (2).

Under the rotating-wave approximation, the interaction Hamiltonian between the atoms and the bath takes the form (see Appendix A for the full derivation)

$$H_{\text{int}} = g_{mj} \sum_{m=1}^2 \sum_{j=1}^{p_m} [\sigma_m^{(1),+} a_{x_{mj}} + \sigma_m^{(2),+} a_{x_{mj}} + \text{H.c.}], \quad (6)$$

where x_{mj} denotes the position of the j th coupling point of atom m and g_{mj} the corresponding coupling strength. Each atom is coupled to p_m cavities in total.

B. Numerical approach to dynamics

Solving the full Hamiltonian H in Eq. (1) is computationally challenging, as it acts on an infinite-dimensional Hilbert space. However, since H conserves the total excitation number of the atom-bath system, it can be projected onto subspaces with a fixed number of excitations. In particular, the dynamics of the iSWAP gate only involve single-excitation states, and the problem can therefore be solved by restricting the Hamiltonian to the single-excitation subspace. In contrast, the dynamics of the CZ gate involve two excitations, going beyond the single-excitation regime that is typically studied in giant-atom systems [53]. As a result, we must numerically simulate

the system in the two-excitation subspace. This subspace is spanned by the states

$$\begin{aligned} |20\rangle_A \otimes |0\rangle_B, & \quad |11\rangle_A \otimes |0\rangle_B, \\ |02\rangle_A \otimes |0\rangle_B, & \quad |10\rangle_A \otimes |j\rangle_B, \\ |01\rangle_A \otimes |j\rangle_B, & \quad |00\rangle_A \otimes |jk\rangle_B, \end{aligned} \quad (7)$$

where $|n_1, n_2\rangle_A$ denotes the state of the two atoms with excitation numbers $n_{1,2} \in \{0, 1, 2\}$, $|0\rangle_B$ is the vacuum state of the waveguide, $|j\rangle_B$ represents a single photon in cavity j , and $|jk\rangle_B$ corresponds to two photons in the waveguide, one in cavity j and one in cavity k ($j = k$ covers the case of two photons in one cavity). Since the photons are bosons, we impose the symmetry condition $|jk\rangle_B = |kj\rangle_B$. Using this basis, we project the full many-body description from the Hamiltonian in Eq. (1) onto the two-excitation subspace and simulate the system dynamics accordingly.

C. Computing gate fidelity

The average gate fidelity of the CZ gate $\bar{\mathcal{F}}_{\text{CZ}}$ is related to the process fidelity $\mathcal{F}_{\text{CZ}}$ by

$$\bar{\mathcal{F}}_{\text{CZ}} = \frac{d\mathcal{F}_{\text{CZ}} + 1}{d + 1}, \quad (8)$$

where $d = 4$ is the Hilbert-space dimension of the computational subspace for the two-qubit CZ gate [95].

To evaluate $\mathcal{F}_{\text{CZ}}$, we first describe the state of the total system as

$$|\psi\rangle = \sum_{m,n} \alpha_{mn} |m\rangle_A \otimes |n\rangle_B, \quad (9)$$

where $|m\rangle_A$ and $|n\rangle_B$ denote atomic and bath states, respectively [see Eq. (7)]. The corresponding density matrix is

$$\rho_{\text{tot}} = \sum_{m,n} \sum_{m',n'} \alpha_{mn} \alpha_{m'n'}^* |m\rangle_A \langle m'|_A \otimes |n\rangle_B \langle n'|_B. \quad (10)$$

Tracing out the bath degrees of freedom gives the reduced atomic density matrix

$$\rho_A = \text{Tr}_B[\rho_{\text{tot}}] = \sum_{m,m'} \alpha_{mn} \alpha_{m'n}^* |m\rangle_A \langle m'|_A. \quad (11)$$

Since the computational subspace only involves the atomic states $\mathcal{B}_{\text{comp}} = \{|11\rangle, |10\rangle, |01\rangle, |00\rangle\}$, we are interested in the further reduction of ρ_A to this subspace. We denote this further reduced density matrix as $\tilde{\rho}_A$; it can be computed as the product of a matrix M with its conjugate transpose

$$M_{mn} = \begin{cases} \alpha_{mn}, & \text{if } |m\rangle_A \in \mathcal{B}_{\text{comp}}, \\ 0, & \text{otherwise,} \end{cases} \quad \tilde{\rho}_A = MM^\dagger, \quad (12)$$

where n denotes the state in the bath $|n\rangle_B$.

Using this further reduced density matrix, a process $\mathcal{E}$ can be represented by its Choi matrix [96]

$$\Phi = \sum_{m,m'=11,10,01,00} |m\rangle \langle m'| \otimes \mathcal{E}(|m\rangle \langle m'|), \quad (13)$$

where $|m\rangle, |m'\rangle$ runs over the computational basis states and $\mathcal{E}(|m\rangle \langle m'|)$ denotes the further reduced density matrix after

evolving the corresponding input state. The process fidelity is then defined as [95]

$$\mathcal{F}_{\text{CZ}} = \left[\text{Tr} \left(\sqrt{\sqrt{\Phi} \Phi_{\text{true}} \sqrt{\Phi}} \right) \right]^2, \quad (14)$$

where Φ_{true} is the Choi matrix of the ideal CZ gate.

We note that the computational subspace contains states from the zero-, one-, and two-excitation sectors. The evolution of the single-excitation states $|10\rangle$ and $|01\rangle$ is obtained from standard single-excitation simulations [53], while the zero-excitation state $|00\rangle$ remains unchanged under the dynamics. Together with the two-excitation dynamics of $|11\rangle$, this fully determines the evolution of initial states from the computational subspace.

III. CONTROLLED-Z GATE

We now investigate how a CZ gate can be implemented using GAs coupled to a structured waveguide. The CZ gate applies a phase shift of π to the $|11\rangle$ state of a two-qubit system while leaving all other basis states unchanged. In our system, this operation can be achieved by coherently transferring the population of the $|11\rangle$ state to $|02\rangle$ or $|20\rangle$ and then bringing it back [97]; see Fig. 1(b). This process requires a resonant transition between these states. In our setup, we realize this resonance by setting $\omega_2 = \omega_1 + \alpha_1$, which enables a resonant coupling between $|11\rangle$ and $|20\rangle$.

To prevent radiative decay during the evolution, we require that transitions at both ω_1 and $\omega_1 + \alpha_1$ must be decoherence-free. Thus, implementing the CZ gate requires at least two decoherence-free frequencies within the photonic band. For a GA with coupling points x_j and coupling strengths g_j , the decoherence-free condition arises from destructive interference [27]

$$\sum_j g_j \exp(ik_{\text{DF}} x_j) = 0, \quad (15)$$

where k_{DF} is a wave number that fulfills this equation.

A. Minimal two-point implementation

The most straightforward configuration that enables fulfilling all the above conditions involves two coupling points separated by Δx and uniform coupling strength, as shown in Fig. 2(a). In this case, Eq. (15) reduces to

$$1 + e^{ik_{\text{DF}} \Delta x} = 0 \quad \Rightarrow \quad k_{\text{DF}} = \frac{\pi + 2n\pi}{\Delta x}, \quad n \in \mathbb{N}. \quad (16)$$

Using the dispersion relation from Eq. (5), we obtain the corresponding decoherence-free frequencies

$$\omega_{\text{DF}} = -2J \cos\left(\frac{\pi + 2n\pi}{\Delta x}\right), \quad n \in \mathbb{N}. \quad (17)$$

To obtain two distinct DF frequencies within the band, we require $\Delta x \geq 3$. However, for $\Delta x = 3$, one of the DF frequencies is located at the band edge ($\omega_{\text{DF}} = 2J$), where strong non-Markovian effects dominate [53]. We therefore choose $\Delta x = 4$, yielding DF frequencies at $\omega_{\text{DF}} = \pm\sqrt{2}J$, as shown in Fig. 2(b).

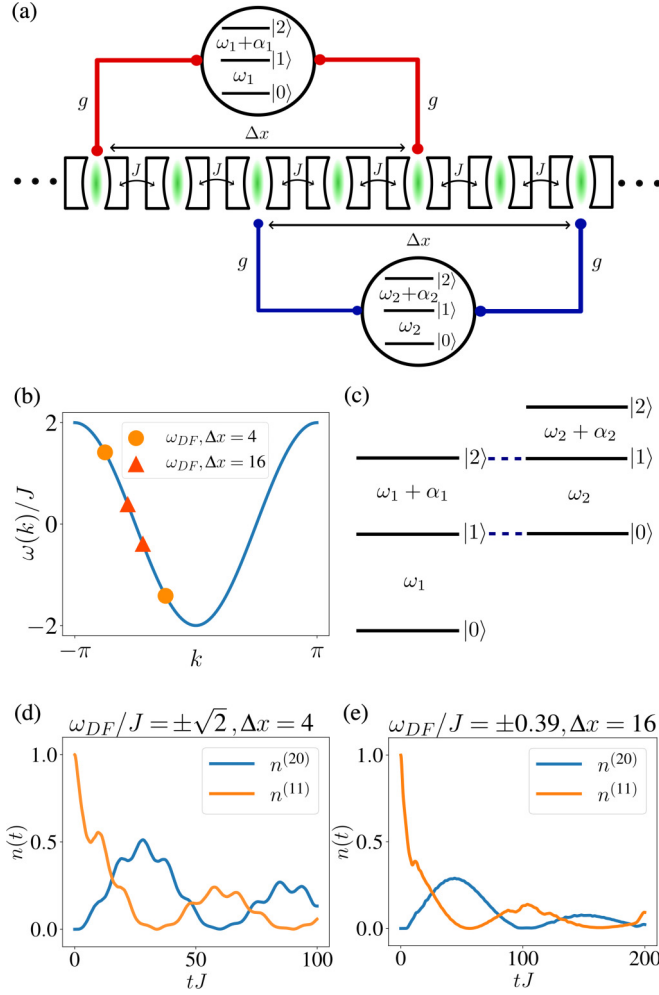


FIG. 2. Setup and results for performing a CZ gate between GAs with two coupling points each. (a) Schematic of the GA-based setup with two coupling points to perform a CZ gate. (b) Bath dispersion relation and the decoherence-free frequencies for $\Delta x = 4$ ($\omega_{DF}/J = \pm\sqrt{2}$) and $\Delta x = 16$ ($\omega_{DF}/J = \pm 0.39$) shown with circles and triangles, respectively. (c) Energy-level diagram of the two GAs for settings enabling the CZ gate. (d) Dynamics in the decoherence-free regime, with parameters $N = 100$, $g/J = 0.175$, $\omega_1/J = \sqrt{2}$, $\omega_2/J = -\sqrt{2}$, $\alpha_1/J = -2\sqrt{2}$, and $\alpha_2/J = -3$. (e) Same as in (d), but with increased separation $\Delta x = 16$ and parameters $\omega_1/J = 0.39$, $\omega_2/J = -0.39$, $\alpha_1/J = -0.78$, and $\alpha_2/J = -1$.

To perform the CZ gate, the level spacings of the two GAs need to satisfy $\omega_1 + \alpha_1 = \omega_2$ [Fig. 2(c)] and that both ω_1 and ω_2 are decoherence-free. These conditions are satisfied for $\Delta x = 4$ with $\omega_1/J = \sqrt{2}$, $\omega_2/J = -\sqrt{2}$, and $\alpha_1/J = -2\sqrt{2}$. The parameter α_2 does not qualitatively influence the dynamics; we set $\alpha_2/J = -3$ to have it similar to the anharmonicity of atom 1.

The dynamics with the system configured thus and initialized in $|11\rangle_A$ is shown in Fig. 2(d). The populations of the $|20\rangle$ and $|11\rangle$ states are computed as

$$n^{(20)}(t) = \frac{1}{2} \langle \psi(t) | \sigma_1^{(2)+} \sigma_1^{(1)+} \sigma_1^{(1)-} \sigma_1^{(2)-} | \psi(t) \rangle, \quad (18)$$

$$n^{(11)}(t) = \langle \psi(t) | \sigma_2^{(1)+} \sigma_1^{(1)+} \sigma_1^{(1)-} \sigma_2^{(1)-} | \psi(t) \rangle, \quad (19)$$

where $|\psi(t)\rangle$ is the state at time t . Although this setup formally satisfies the DF condition, the DF frequencies lie far from the band center, leading to strong non-Markovian effects [53].

The first consequence of non-Markovianity is leakage of excitations into the waveguide before the DFI is established. Since the group velocity v_g is maximal at the band center and decreases towards the band edge, the propagation delay $\tau = \Delta x/v_g$ increases as the atomic frequency approaches the band edge. As a result, the DFI is not established instantaneously, and part of the atomic excitation decays during this delay time.

The second consequence of non-Markovianity is a slower leakage of excitations into the waveguide even after the DFI has been established. Atoms tuned close to the band edge hybridize strongly with photonic bound states, which in turn hybridize with the continuum of propagating modes. This leads to a gradual loss of excitation from the atomic subsystem and a corresponding decay of the DFI.

As a result, population decays rapidly, the DFI breaks down, and the two-point configuration fails to support a high-fidelity CZ gate. We note that the failure of the two-point coupling configuration with large Δx to support high-fidelity population transfer is already present in the single-excitation sector (see Appendix B). Therefore, the poor performance observed in the two-excitation sector is not unique to the CZ-gate protocol, but rather originates from the same non-Markovian effects that already degrade the single-excitation dynamics.

A natural way to address this issue is to increase the separation Δx , which shifts the DF frequencies closer to the band center. For example, $\Delta x = 16$ yields $\omega_{DF}/J = \pm 0.39$, as shown in Fig. 2(b). However, as shown in Fig. 2(e), this strategy does not improve performance: the larger separation introduces a significant time delay for photons traveling between coupling points, enhancing non-Markovian effects leading to significant population loss into the waveguide.

B. Improvement with a three-point implementation

To overcome the limitations of the two-point configuration, we introduce a third coupling point placed in the middle between the original two, with a relative coupling strength ζ , as shown in Fig. 3(a). From here on, we use Δx to denote the distance between the middle coupling point and each of the outer coupling points, such that the distance between the outer coupling points is $2\Delta x$. In this setup, the decoherence-free condition in Eq. (15) is modified to

$$1 + \zeta e^{ik_{DF}\Delta x} + e^{2ik_{DF}\Delta x} = 0. \quad (20)$$

The corresponding decoherence-free wave vectors are

$$k_{DF} = \frac{\arccos\left(-\frac{\zeta}{2}\right) + 2n\pi}{\Delta x}, \quad n \in \mathbb{N}, \quad (21)$$

with associated frequencies

$$\omega_{DF} = -2J \cos\left[\frac{\arccos\left(-\frac{\zeta}{2}\right) + 2n\pi}{\Delta x}\right], \quad n \in \mathbb{N}. \quad (22)$$

The dependence of ω_{DF} on ζ for $\Delta x = 2$ is shown in Fig. 3(b). As ζ increases, the DF frequencies shift progressively toward the center of the photonic band. At $\zeta = 2$, when the coupling strength in the middle coupling point equals the

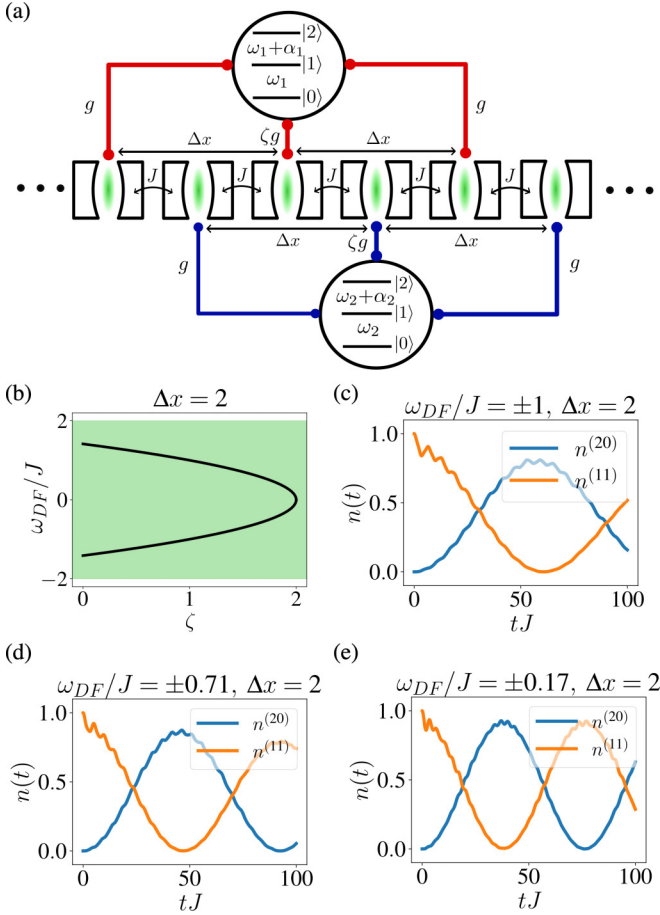


FIG. 3. Setup and results for performing a CZ gate between GAs with three coupling points each. (a) Schematic of the GA-based setup with three coupling points to perform a CZ gate, with $\Delta x = 2$ between adjacent coupling points. Compared to Fig. 2(a), a third, central coupling point with coupling strength ζg is added. This configuration yields two decoherence-free (DF) frequencies for $0 \leq \zeta < 2$. (b) Decoherence-free frequencies as a function of ζ , showing that the two frequencies merge at $\zeta = 2$ with the band of propagating frequencies between $[-2, 2]$ shown in light green. (c) Dynamics in the DFI regime for $N = 100$, $\Delta x = 2$, $g/J = 0.1$, and $\zeta = 1$, yielding DF frequencies $\omega_{DF}/J = \pm 1$. The frequencies and anharmonicities of the GAs are chosen to fit in the CZ-gate configuration: $\omega_1/J = 1$, $\omega_2/J = -0.98$, $\alpha_1/J = -2$, and $\alpha_2/J = -1.52$. (d) Same as (c), but with $\zeta = 1.5$, giving DF frequencies $\omega_{DF}/J = \pm 0.71$. Here, we take parameters $\omega_1/J = 0.71$, $\omega_2/J = -0.69$, $\alpha_1/J = -1.42$, and $\alpha_2/J = -1.31$. (e) Same as (c) and (d), but with $\zeta = 1.97$, resulting in DF frequencies $\omega_{DF}/J = \pm 0.17$. Here, we take parameters $\omega_1/J = 0.17$, $\omega_2/J = -0.17$, $\alpha_1/J = -0.34$, and $\alpha_2/J = -0.67$.

sum of the coupling strengths in the outer coupling points, the two solutions merge at $\omega_{DF} = 0$.

The dynamics for this configuration are illustrated in Figs. 3(c)–3(e) for three successively increasing values of ζ matched to the decoherence-free conditions required for the CZ gate. We have chosen $N = 100$ cavities, for which finite-size effects are already negligible (see Appendix C). As ζ increases, the quality of the DFI improves: after a full cycle, a larger fraction of the population returns to $|11\rangle$, which is essential for high-fidelity CZ gate operation. This enhancement

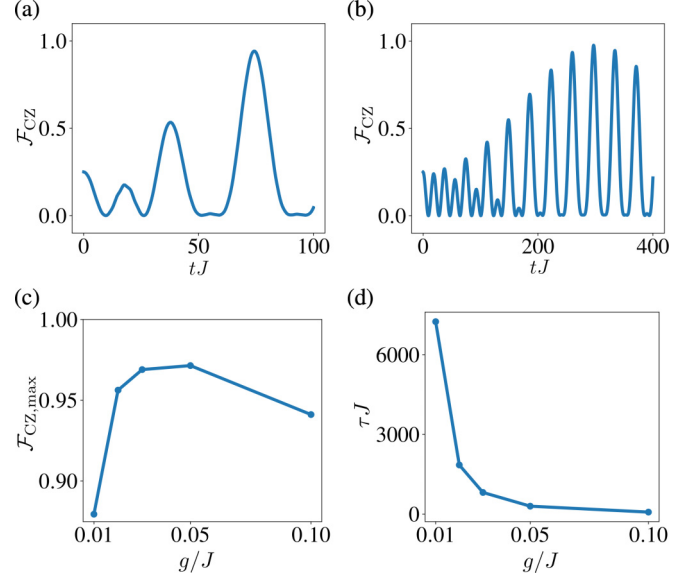


FIG. 4. Process fidelity $\mathcal{F}_{CZ}$ of CZ gates performed with the setup in Fig. 3(a). (a) Process fidelity as a function of time t for the same parameters as in Fig. 3(c). (b) Same as (a), but with a smaller $g = 0.05J$. (c) The best process fidelity of CZ gates, $\mathcal{F}_{CZ,max}$, for different values of g , considering realistic qubit and cavity decay rates of $\Gamma_q = 1.6 \times 10^{-5}J$ and $\Gamma_c = 8 \times 10^{-5}J$. (d) The gate time τ required for the best CZ gate in (c).

stems from the suppression of non-Markovian effects since the relevant DF frequencies are brought closer to the band center.

We note that the $|1\rangle \rightarrow |0\rangle$ transition of atom 2 is also off-resonantly coupled to the $|1\rangle \rightarrow |0\rangle$ transition in atom 1. As a consequence, to obtain a coherent population exchange between the atoms, ω_2 needs to deviate slightly from $\omega_1 + \alpha_1$ to take into account the Lamb shift [98] resulting from this coupling; hence the parameter choices in Figs. 3(c) and 3(d). We also note that the gate dynamics become faster for larger ζ , as a consequence of the increased effective coupling between the atom and the waveguide.

Finally, we note that introducing additional coupling points could provide further flexibility in engineering the decoherence-free frequencies. This additional flexibility would have to be weighed against both the increase in the complexity of the experimental implementation and the drawback that the distance between the outermost coupling points would be increased, thereby making non-Markovian effects stronger again. More coupling points become particularly relevant when one wishes to realize protocols requiring more than two decoherence-free frequencies [32,33].

C. Gate fidelity

Having analyzed the improved setup for performing a CZ gate with GAs in Fig. 3, we now compute the corresponding gate fidelity, as described in Sec. II C. The process fidelity $\mathcal{F}_{CZ}$ during the time evolution is shown in Fig. 4(a) for the parameters from Fig. 3(e). A maximum process fidelity of 94.2% is achieved at $tJ = 74.0$, corresponding to the time at which the population is transferred from $|11\rangle$ to $|20\rangle$ and back. The

corresponding average gate fidelity is $\bar{F}_{\text{CZ}} \approx 95.4\%$. For ideal qubits and cavities, the fidelity can be further improved by reducing the coupling strength g . For example, choosing $g = 0.05J$ increases the process fidelity to $\mathcal{F}_{\text{CZ}} \approx 97.6\%$, reached at $tJ = 296.8$ [Fig. 4(b)]. We note that the main oscillation in Fig. 4(b) corresponds to the population transfer $|11\rangle \rightarrow |20\rangle \rightarrow -|11\rangle$ that implements the CZ gate. The smaller oscillations and secondary peaks originate from a high-frequency phase oscillation within the computational subspace.

In realistic scenarios, where qubits and cavities exhibit intrinsic decay, a comprehensive optimization of the coupling strength g would require solving the full master equation. This significantly increases the computational cost, as the Liouvillian superoperator scales with the square of the Hilbert-space dimension [99]. To make the problem tractable, we instead analyze decay effects using the following effective non-Hermitian Hamiltonian, obtained by neglecting quantum jumps and collective decay processes

$$H_{\text{eff}} = H - \frac{i\Gamma_q}{2} \sum_{m=1}^2 [\sigma_m^{(2),+} \sigma_m^{(2),-} + \sigma_m^{(1),+} \sigma_m^{(1),-}] - \frac{i\Gamma_c}{2} \sum_{n=1}^N a_n^\dagger a_n. \quad (23)$$

Here, H denotes the ideal system Hamiltonian defined in Eq. (1) and $\Gamma_{q(c)}$ is the qubit (cavity) decay rate, assumed identical for all qubits (cavities). The higher decay rate of the $|2\rangle$ level compared to $|1\rangle$ is incorporated through the definition of $\sigma_m^{(2),\pm}$ in Eq. (3).

We now evaluate the CZ gate fidelity for experimentally relevant parameters in superconducting circuits, where GAs have already been demonstrated [29,34,37,39,70]. In particular, transmon qubits [100] coupled to microwave resonators provide a promising platform for implementing our proposal. In state-of-the-art devices, the transmon anharmonicity can be as low as $-\alpha/2\pi \approx 50$ MHz [101,102], while the intercavity coupling between microwave resonators can reach $J/2\pi \approx 200$ MHz [103], yielding a ratio $-\alpha/J \approx 1/4$. Our choice of $-\alpha/J = 1/3$ in Fig. 3(e) is therefore conservative. Furthermore, a coupling strength of $g/2\pi = 0.05J/2\pi \approx 10$ MHz is well within the reach of current superconducting circuit technology [29,39]. The required strength of the central coupling point with $0 \leq \zeta < 2$ is modest and experimentally accessible [29,39]. A conservative estimation of typical decay rates for transmons and microwave resonators are $\Gamma_q = 0.02$ MHz and $\Gamma_c = 0.1$ MHz, respectively (for resonators operating at 5 GHz; even lower rates are achievable at smaller resonator frequencies) [104–112]. For $J/2\pi = 200$ MHz, these values correspond to $\Gamma_q = 1.6 \times 10^{-5}J$ and $\Gamma_c = 8 \times 10^{-5}J$, which we adopt in our analysis.

For different values of g , the optimal process fidelity $\mathcal{F}_{\text{CZ,max}}$ and the corresponding gate time τ are shown in Figs. 4(c) and 4(d). For large g , the gate time τ is short, and decoherence from qubits and cavities does not constitute the primary limitation to fidelity [113,114]. In contrast, for small g , the increased gate time makes intrinsic decay the dominant source of infidelity. Consequently, the optimal performance is obtained at $g = 0.05J$ with gate time $\tau = 297/J \approx 236$ ns, yielding a process fidelity of $\mathcal{F}_{\text{CZ}} \approx 97.1\%$, corresponding to

an average gate fidelity of $\bar{F}_{\text{CZ}} \approx 97.7\%$. We note that since quantum jumps introduce additional decoherence and leakage, neglecting them by using the non-Hermitian Hamiltonian in Eq. (23) generally leads to a slight overestimation of the gate fidelity. Therefore, the fidelities reported here should be regarded as approximate upper bounds on the fidelities obtained from a full master-equation treatment.

IV. CONCLUSION

We demonstrated how a controlled-z (CZ) gate can be implemented using GAs coupled to a structured waveguide. Starting from a minimal two-point coupling configuration, we showed that although the required DFI can be realized, strong non-Markovian effects significantly limit the achievable gate fidelity. To address this limitation, we proposed an extended configuration that includes a third, intermediate coupling point. This modification shifts the decoherence-free frequencies closer to the center of the photonic band, thereby mitigating non-Markovian effects. With this design, we can achieve CZ gates with average gate fidelities up to 97.7% when including realistic imperfections like typical resonator and transmon decay rates; remaining non-Markovian effects from spacing between coupling points also contribute to reducing the gate fidelity.

Beyond the specific case of the CZ gate, our results establish a general strategy for engineering multiphoton processes in structured waveguide QED by tailoring interference conditions through multipoint coupling. This approach opens possibilities for realizing other entangling operations, exploiting strongly correlated photonic dynamics, and extending gate protocols to larger networks of GAs. Moreover, the protocol is compatible with state-of-the-art superconducting circuits, where transmon coherence times are orders of magnitude longer than the gate duration, suggesting that experimental realization is within reach.

Looking ahead, our findings point toward several promising directions. One is the integration of multiple GAs in extended resonator arrays, enabling scalable architectures for quantum computation and simulation, such as passive-leakage reset for high-fidelity gates [115]. Another is the exploration of richer structured environments, such as topological or chiral waveguides, where decoherence-free conditions could be further stabilized, or new types of protected interactions could emerge. More broadly, the interplay between photon-photon interactions [116], engineered dispersion, and non-Markovian dynamics in these systems defines a fertile platform for advancing fundamental studies of open quantum systems.

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DATA AVAILABILITY

The software that supports the findings of this article is openly available at [117].

APPENDIX A: ROTATING-WAVE APPROXIMATION

Before applying the rotating-wave approximation (RWA) to arrive at Eq. (6), the interaction between the giant atoms and the waveguide is described by

$$H_{\text{int}}^{\text{full}} = \sum_{m=1}^2 \sum_{j=1}^{p_m} g_{mj} (\sigma_m^{(1),+} + \sigma_m^{(2),+} + \text{H.c.}) (a_{x_{mj}} + a_{x_{mj}}^\dagger), \quad (\text{A1})$$

where the notation is defined after Eq. (6) in the main text. The free Hamiltonian is

$$H_0 = H_A + H_B. \quad (\text{A2})$$

The evolution of a state $|\psi(t)\rangle$ for the giant atoms and the waveguide is governed by the Schrödinger equation

$$i \frac{d}{dt} |\psi(t)\rangle = (H_0 + H_{\text{int}}^{\text{full}}) |\psi(t)\rangle. \quad (\text{A3})$$

It is convenient to move to the interaction picture, where the state vector is defined as

$$|\psi_I(t)\rangle = e^{iH_0 t} |\psi(t)\rangle. \quad (\text{A4})$$

The dynamics in this picture are governed by the interaction Hamiltonian

$$H_I(t) = e^{iH_0 t} H_{\text{int}}^{\text{full}} e^{-iH_0 t} \quad (\text{A5})$$

and the corresponding evolution operator

$$U_I(t) = \mathcal{T} \exp \left[-i \int_0^t H_I(t') dt' \right], \quad (\text{A6})$$

where $\mathcal{T}$ denotes the time-ordering operator. The interaction-picture state therefore evolves as

$$|\psi_I(t)\rangle = U_I(t) |\psi_I(0)\rangle. \quad (\text{A7})$$

The advantage of the interaction picture is that it separates the rapid oscillations generated by the free Hamiltonian from the slower interaction dynamics. Terms in $H_I(t)$ that oscillate slowly contribute coherently to the time integral and therefore strongly affect the evolution. In contrast, rapidly oscillating terms tend to average out and contribute only weakly. For the atomic operators, we obtain

$$\begin{aligned} e^{iH_0 t} \sigma_m^{(1),+} e^{-iH_0 t} &= e^{i\omega_m t} \sigma_m^{(1),+}, \\ e^{iH_0 t} \sigma_m^{(2),+} e^{-iH_0 t} &= e^{i(\omega_m + \alpha_m) t} \sigma_m^{(2),+}, \end{aligned} \quad (\text{A8})$$

with the Hermitian conjugates acquiring the opposite phases. Similarly, for the waveguide modes

$$e^{iH_0 t} a_k e^{-iH_0 t} = e^{-i\omega_k t} a_k, \quad (\text{A9})$$

and

$$e^{iH_0 t} a_k^\dagger e^{-iH_0 t} = e^{i\omega_k t} a_k^\dagger. \quad (\text{A10})$$

The excitation-conserving terms therefore acquire phases

$$\begin{aligned} e^{iH_0 t} \sigma_m^{(1),+} a_k e^{-iH_0 t} &= e^{i(\omega_m - \omega_k) t} \sigma_m^{(1),+} a_k, \\ e^{iH_0 t} \sigma_m^{(1),-} a_k^\dagger e^{-iH_0 t} &= e^{-i(\omega_m - \omega_k) t} \sigma_m^{(1),-} a_k^\dagger, \end{aligned} \quad (\text{A11})$$

and similarly for the $\sigma_m^{(2),\pm}$ terms. Near resonance, $\omega_m \approx \omega_k$ (or $\omega_m + \alpha_m \approx \omega_k$), these terms oscillate slowly and therefore accumulate coherently in the evolution operator.

In contrast, the counterrotating terms acquire phases

$$\begin{aligned} e^{iH_0 t} \sigma_m^{(1),+} a_k^\dagger e^{-iH_0 t} &= e^{i(\omega_m + \omega_k) t} \sigma_m^{(1),+} a_k^\dagger, \\ e^{iH_0 t} \sigma_m^{(1),-} a_k e^{-iH_0 t} &= e^{-i(\omega_m + \omega_k) t} \sigma_m^{(1),-} a_k, \end{aligned} \quad (\text{A12})$$

and similarly for the $\sigma_m^{(2),\pm}$ terms. Their contribution to the evolution operator is proportional to

$$g \int_0^t e^{i(\omega_m + \omega_k) t'} dt' = g \frac{e^{i(\omega_m + \omega_k) t} - 1}{i(\omega_m + \omega_k)}, \quad (\text{A13})$$

whose magnitude is bounded by

$$\left| g \int_0^t e^{i(\omega_m + \omega_k) t'} dt' \right| \leq \frac{2g}{\omega_m + \omega_k}. \quad (\text{A14})$$

Therefore, when

$$g \ll \omega_m, \omega_k, \quad (\text{A15})$$

the effect of the counterrotating terms is suppressed by the small parameter $g/(\omega_m + \omega_k)$ and can be neglected within the RWA. Retaining only the excitation-conserving terms yields

$$H_{\text{int}} = \sum_{m=1}^2 \sum_{j=1}^{p_m} g_{mj} [\sigma_m^{(1),+} a_{x_{mj}} + \sigma_m^{(2),+} a_{x_{mj}} + \text{H.c.}], \quad (\text{A16})$$

which is Eq. (6) in the main text.

APPENDIX B: NON-MARKOVIAN EFFECTS IN THE SINGLE-EXCITATION SECTOR

In this appendix, we show that the non-Markovian effects degrading the DFI in Fig. 2(d) are already present in the single-excitation sector. We also show that the mitigation strategy of adding a third coupling point, as in the two-excitation sector in Figs. 3(c)–3(e), works in the single-excitation sector as well.

Figure 5(a) shows the dynamics in the single-excitation sector for two GAs with two coupling points each. Here, we see that the excitation exchange between $|10\rangle$ and $|01\rangle$ at a supposedly decoherence-free frequency is strongly degraded by non-Markovian effects. This is the same behavior as we observed in the two-excitation sector [Fig. 2(d)], indicating that the reduced CZ-gate fidelity is not a phenomenon unique to the two-excitation manifold.

As shown in Fig. 5(b), the addition of a third coupling point substantially improves (in the sense of reducing leakage) the single-excitation dynamics by shifting the decoherence-free frequencies closer to the band center. The same strategy is em-

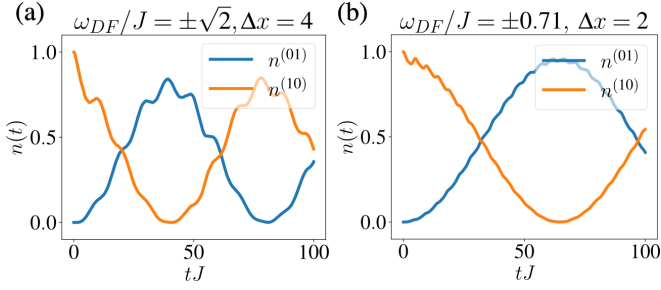


FIG. 5. Non-Markovian effects in the single-excitation sector. (a) Dynamics in the DFI regime for two GAs with two coupling points each and a single excitation. Parameters: $N = 100$, $g/J = 0.175$, $\omega_1/J = \sqrt{2}$, and $\omega_2/J = \sqrt{2}$. (b) Dynamics in the DFI regime for two GAs with three coupling points each and a single excitation. Parameters: $N = 100$, $\Delta x = 2$, $g/J = 0.1$, and $\zeta = 1.97$, yielding DF frequencies $\omega_{DF}/J = \pm 0.71$. The frequencies and anharmonicities of the GAs are chosen to fit in the CZ-gate configuration: $\omega_1/J = 0.71$, $\omega_2/J = 0.71$.

ployed in Figs. 3(c)–3(e) for the two-excitation sector, where it similarly suppressed non-Markovian effects and enhanced the gate fidelity. These results indicate that the trends in the single- and two-excitation sectors are consistent and that they are governed by the same underlying mechanisms.

APPENDIX C: FINITE-SIZE EFFECTS

To verify that our results are not affected by finite-size effects, we repeated the simulation from Fig. 3(e) for different waveguide lengths. As shown in Fig. 6(a), the gate dynamics for $N = 50$ are nearly indistinguishable from the case of $N = 100$ in Fig. 3(e), indicating that increasing the waveguide size does not lead to any noticeable difference in the gate performance. In contrast, for a much smaller waveguide with $N = 10$, pronounced oscillations appear in the dynamics [Fig. 6(b)], indicating significant finite-size effects. These results demonstrate that the waveguide sizes used in the article are sufficient to accurately capture the gate dynamics.

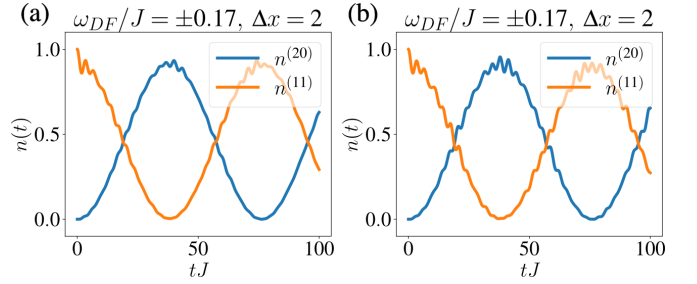


FIG. 6. Dynamics in the DFI regime for the same parameters as in Fig. 3(e), but with different waveguide lengths. In Fig. 3(e), $N = 100$ cavities were used in the numerical simulations. Here, we use (a) $N = 50$ and (b) $N = 10$.

- [1] I. M. Georgescu, S. Ashhab, and F. Nori, Quantum simulation, *Rev. Mod. Phys.* **86**, 153 (2014).
- [2] E. Altman, K. R. Brown, G. Carleo, L. D. Carr, E. Demler, C. Chin, B. DeMarco, S. E. Economou, M. A. Eriksson, K.-M. C. Fu, M. Greiner, K. R. A. Hazzard, R. G. Hulet, A. J. Kollár, B. L. Lev, M. D. Lukin, R. Ma, X. Mi, S. Misra, C. Monroe, *et al.*, Quantum simulators: Architectures and opportunities, *PRX Quantum* **2**, 017003 (2021).
- [3] B. Fauseweh, Quantum many-body simulations on digital quantum computers: State-of-the-art and future challenges, *Nat. Commun.* **15**, 2123 (2024).
- [4] K. Wintersperger, F. Dommert, T. Ehmer, A. Housanov, J. Klepsch, W. Maurer, G. Reuber, T. Strohm, M. Yin, and S. Luber, Neutral atom quantum computing hardware: Performance and end-user perspective, *EPJ Quantum Technol.* **10**, 32 (2023).
- [5] T. Strohm, K. Wintersperger, F. Dommert, D. Basilewitsch, G. Reuber, A. Housanov, T. Ehmer, D. Vodola, and S. Luber, Ion-based quantum computing hardware: Performance and end-user perspective, [arXiv:2405.11450](https://arxiv.org/abs/2405.11450).
- [6] S. A. Caldwell, N. Didier, C. A. Ryan, E. A. Sete, A. Hudson, P. Karalekas, R. Manenti, M. P. da Silva, R. Sinclair, E. Acala, N. Alidoust, J. Angeles, A. Bestwick, M. Block, B. Bloom, A. Bradley, C. Bui, L. Capelluto, R. Chilcott, J. Cordova, *et al.*, Parametrically activated entangling gates using transmon qubits, *Phys. Rev. Appl.* **10**, 034050 (2018).
- [7] R. Barends, C. M. Quintana, A. G. Petukhov, Y. Chen, D. Kafri, K. Kechedzhi, R. Collins, O. Naaman, S. Boixo, F. Arute, K. Arya, D. Buell, B. Burkett, Z. Chen, B. Chiaro, A. Dunsworth, B. Foxen, A. Fowler, C. Gidney, M. Giustina, *et al.*, Diabatic gates for frequency-tunable superconducting qubits, *Phys. Rev. Lett.* **123**, 210501 (2019).
- [8] A. Noguchi, A. Osada, S. Masuda, S. Kono, K. Heya, S. P. Wolski, H. Takahashi, T. Sugiyama, D. Lachance-Quirion, and Y. Nakamura, Fast parametric two-qubit gates with suppressed residual interaction using the second-order nonlinearity of a cubic transmon, *Phys. Rev. A* **102**, 062408 (2020).
- [9] D. M. Abrams, N. Didier, B. R. Johnson, M. P. d. Silva, and C. A. Ryan, Implementation of XY entangling gates with a single calibrated pulse, *Nat. Electron.* **3**, 744 (2020).
- [10] M. Ganzhorn, G. Salis, D. J. Egger, A. Fuhrer, M. Mergenthaler, C. Müller, P. Müller, S. Paredes, M. Pechal, M. Werninghaus, and S. Filipp, Benchmarking the noise sensitivity of different parametric two-qubit gates in a single superconducting quantum computing platform, *Phys. Rev. Res.* **2**, 033447 (2020).
- [11] N. Lacroix, C. Hellings, C. K. Andersen, A. Di Paolo, A. Remm, S. Lazar, S. Krinner, G. J. Norris, M. Gabureac, J. Heinsoo, A. Blais, C. Eichler, and A. Wallraff, Improving the performance of deep quantum optimization algorithms with continuous gate sets, *PRX Quantum* **1**, 020304 (2020).
- [12] E. A. Sete, N. Didier, A. Q. Chen, S. Kulshreshtha, R. Manenti, and S. Poletto, Parametric-resonance entangling gates with a tunable coupler, *Phys. Rev. Appl.* **16**, 024050 (2021).
- [13] Y. Sung, L. Ding, J. Braumüller, A. Vepsäläinen, B. Kannan, M. Kjaergaard, A. Greene, G. O. Samach, C. McNally, D. Kim, A. Melville, B. M. Niedzielski, M. E. Schwartz, J. L. Yoder, T. P. Orlando, S. Gustavsson, and W. D. Oliver, Re-

- alization of high-fidelity CZ and ZZ-free iSWAP gates with a tunable coupler, *Phys. Rev. X* **11**, 021058 (2021).
- [14] Rigetti QCS, *Aspen-M-3 quantum processor* (2023).
- [15] H. Zhang, C. Ding, D. K. Weiss, Z. Huang, Y. Ma, C. Guinn, S. Sussman, S. P. Chitta, D. Chen, A. A. Houck, J. Koch, and D. I. Schuster, Tunable inductive coupler for high-fidelity gates between fluxonium qubits, *PRX Quantum* **5**, 020326 (2024).
- [16] Z. Chen, W. Liu, Y. Ma, W. Sun, R. Wang, H. Wang, H. Xu, G. Xue, H. Yan, Z. Yang, J. Ding, Y. Gao, F. Li, Y. Zhang, Z. Zhang, Y. Jin, H. Yu, J. Chen, and F. Yan, Efficient implementation of arbitrary two-qubit gates via unified control, *Nat. Phys.* **21**, 1489 (2025).
- [17] T.-M. Li, J.-C. Zhang, B.-J. Chen, K. Huang, H.-T. Liu, Y.-X. Xiao, C.-L. Deng, G.-H. Liang, C.-T. Chen, Y. Liu, H. Li, Z.-T. Bao, K. Zhao, Y. Xu, L. Li, Y. He, Z.-H. Liu, Y.-H. Yu, S.-Y. Zhou, Y.-J. Liu, *et al.*, High-precision pulse calibration of tunable couplers for high-fidelity two-qubit gates in superconducting quantum processors, *Phys. Rev. Appl.* **23**, 024059 (2025).
- [18] C. Križan, J. Biznárová, L. Chen, E. Hogedal, A. Osman, C. W. Warren, S. Kosen, H.-X. Li, T. Abad, A. Aggarwal, M. Caputo, J. Fernández-Pendás, A. Gaikwad, L. Grönberg, A. Nylander, R. Rehammar, M. Rommel, O. I. Yuzepovich, A. Frisk Kockum, J. Govenius, *et al.*, Quantum swap gate realized with ca and iswap gates in a superconducting architecture, *New J. Phys.* **27**, 074507 (2025).
- [19] H.-T. Liu, B.-J. Chen, J.-C. Zhang, Y.-X. Xiao, T.-M. Li, K. Huang, Z. Wang, H. Li, K. Zhao, Y. Xu, C.-L. Deng, G.-H. Liang, Z.-H. Liu, S.-Y. Zhou, C.-P. Fang, X. Song, Z. Xiang, D. Zheng, Y.-H. Shi, K. Xu, *et al.*, Direct implementation of high-fidelity three-qubit gates for superconducting processor with tunable couplers, *Phys. Rev. Lett.* **135**, 050602 (2025).
- [20] K. X. Wei, E. Magesan, I. Lauer, S. Srinivasan, D. F. Bogorin, S. Carnevale, G. A. Keefe, Y. Kim, D. Klaus, W. Landers, N. Sundaresan, C. Wang, E. J. Zhang, M. Steffen, O. E. Dial, D. C. McKay, and A. Kandala, Hamiltonian engineering with multicolor drives for fast entangling gates and quantum crosstalk cancellation, *Phys. Rev. Lett.* **129**, 060501 (2022).
- [21] L. B. Nguyen, Y. Kim, A. Hashim, N. Goss, B. Marinelli, B. Bhandari, D. Das, R. K. Naik, J. M. Kreikebaum, A. N. Jordan, D. I. Santiago, and I. Siddiqi, Programmable Heisenberg interactions between Floquet qubits, *Nat. Phys.* **20**, 240 (2024).
- [22] S. J. Evered, D. Bluvstein, M. Kalinowski, S. Ebadi, T. Manovitz, H. Zhou, S. H. Li, A. A. Geim, T. T. Wang, N. Maskara, H. Levine, G. Semeghini, M. Greiner, V. Vuletić, and M. D. Lukin, High-fidelity parallel entangling gates on a neutral-atom quantum computer, *Nature (London)* **622**, 268 (2023).
- [23] B. Foxen, C. Neill, A. Dunsworth, P. Roushan, B. Chiaro, A. Megrant, J. Kelly, Z. Chen, K. Satzinger, R. Barends, F. Arute, K. Arya, R. Babbush, D. Bacon, J. C. Bardin, S. Boixo, D. Buell, B. Burkett, Y. Chen, R. Collins, *et al.*, Demonstrating a continuous set of two-qubit gates for near-term quantum algorithms, *Phys. Rev. Lett.* **125**, 120504 (2020).
- [24] K. X. Wei, I. Lauer, E. Pritchett, W. Shanks, D. C. McKay, and A. Javadi-Abhari, Native two-qubit gates in fixed-coupling, fixed-frequency transmons beyond cross-resonance interaction, *PRX Quantum* **5**, 020338 (2024).
- [25] A. F. Kockum, Quantum optics with giant atoms: The first five years, in *International Symposium on Mathematics, Quantum Theory, and Cryptography*, Mathematics for Industry, Vol. 33 (Springer, New York, 2021), pp. 125–146.
- [26] M. V. Gustafsson, T. Aref, A. F. Kockum, M. K. Ekström, G. Johansson, and P. Delsing, Propagating phonons coupled to an artificial atom, *Science* **346**, 207 (2014).
- [27] A. Frisk Kockum, P. Delsing, and G. Johansson, Designing frequency-dependent relaxation rates and Lamb shifts for a giant artificial atom, *Phys. Rev. A* **90**, 013837 (2014).
- [28] A. F. Kockum, G. Johansson, and F. Nori, Decoherence-free interaction between giant atoms in waveguide quantum electrodynamics, *Phys. Rev. Lett.* **120**, 140404 (2018).
- [29] B. Kannan, M. J. Ruckriegel, D. L. Campbell, A. Frisk Kockum, J. Braumüller, D. K. Kim, M. Kjaergaard, P. Krantz, A. Melville, B. M. Niedzielski, A. Vepsäläinen, R. Winik, J. L. Yoder, F. Nori, T. P. Orlando, S. Gustavsson, and W. D. Oliver, Waveguide quantum electrodynamics with superconducting artificial giant atoms, *Nature (London)* **583**, 775 (2020).
- [30] G. Chen and A. Frisk Kockum, Simulating open quantum systems with giant atoms, *Quantum Sci. Technol.* **10**, 025028 (2025).
- [31] C. L. Yang and W. Z. Jia, Simulating and probing many-body quantum states in waveguide-QED systems with giant atoms, *Phys. Rev. A* **111**, 023712 (2025).
- [32] G. Chen and A. F. Kockum, Scalable quantum simulator with an extended gate set in giant atoms, *Quantum* **10**, 1992 (2026).
- [33] G. Chen and A. F. Kockum, Efficient three-qubit gates with giant atoms, [arXiv:2510.04545](https://arxiv.org/abs/2510.04545).
- [34] A. M. Vadiraj, A. Ask, T. G. McConkey, I. Nsanzineza, C. W. S. Chang, A. F. Kockum, and C. M. Wilson, Engineering the level structure of a giant artificial atom in waveguide quantum electrodynamics, *Phys. Rev. A* **103**, 023710 (2021).
- [35] Z. Q. Wang, Y. P. Wang, J. Yao, R. C. Shen, W. J. Wu, J. Qian, J. Li, S. Y. Zhu, and J. Q. You, Giant spin ensembles in waveguide magnonics, *Nat. Commun.* **13**, 7580 (2022).
- [36] Y.-T. Chen, L. Du, L. Guo, Z. Wang, Y. Zhang, Y. Li, and J.-H. Wu, Nonreciprocal and chiral single-photon scattering for giant atoms, *Commun. Phys.* **5**, 215 (2022).
- [37] C. Joshi, F. Yang, and M. Mirhosseini, Resonance fluorescence of a chiral artificial atom, *Phys. Rev. X* **13**, 021039 (2023).
- [38] X. Wang, H.-B. Zhu, T. Liu, and F. Nori, Realizing quantum optics in structured environments with giant atoms, *Phys. Rev. Res.* **6**, 013279 (2024).
- [39] V. Jouanny, L. Peyruchat, M. Scigliuzzo, A. Mercurio, E. Di Benedetto, D. De Bernardis, D. Sbroggitto, S. Frasca, V. Savona, F. Ciccarello, and P. Scarlino, Superstrong dynamics and chiral emission of a giant atom in a structured bath, [arXiv:2509.01579](https://arxiv.org/abs/2509.01579).
- [40] L. Guo, A. L. Grimsmo, A. F. Kockum, M. Pletyukhov, and G. Johansson, Giant acoustic atom: A single quantum system with a deterministic time delay, *Phys. Rev. A* **95**, 053821 (2017).
- [41] A. González-Tudela, C. Sánchez Muñoz, and J. I. Cirac, Engineering and harnessing giant atoms in high-dimensional baths: A proposal for implementation with cold atoms, *Phys. Rev. Lett.* **122**, 203603 (2019).
- [42] L. Guo, A. F. Kockum, F. Marquardt, and G. Johansson, Oscillating bound states for a giant atom, *Phys. Rev. Res.* **2**, 043014 (2020).

- [43] P. O. Guimond, B. Vermersch, M. L. Juan, A. Sharafiev, G. Kirchmair, and P. Zoller, A unidirectional on-chip photonic interface for superconducting circuits, *npj Quantum Inf.* **6**, 32 (2020).
- [44] A. Ask, Y.-L. L. Fang, and A. F. Kockum, Synthesizing electromagnetically induced transparency without a control field in waveguide QED using small and giant atoms, *arXiv:2011.15077*.
- [45] D. Cilluffo, A. Carollo, S. Lorenzo, J. A. Gross, G. M. Palma, and F. Ciccarello, Collisional picture of quantum optics with giant emitters, *Phys. Rev. Res.* **2**, 043070 (2020).
- [46] X. Wang, T. Liu, A. F. Kockum, H.-R. Li, and F. Nori, Tunable chiral bound states with giant atoms, *Phys. Rev. Lett.* **126**, 043602 (2021).
- [47] L. Du and Y. Li, Single-photon frequency conversion via a giant Λ -type atom, *Phys. Rev. A* **104**, 023712 (2021).
- [48] A. Soro and A. F. Kockum, Chiral quantum optics with giant atoms, *Phys. Rev. A* **105**, 023712 (2022).
- [49] X. Wang and H.-R. Li, Chiral quantum network with giant atoms, *Quantum Sci. Technol.* **7**, 035007 (2022).
- [50] L. Du, Y. Zhang, J. H. Wu, A. F. Kockum, and Y. Li, Giant atoms in a synthetic frequency dimension, *Phys. Rev. Lett.* **128**, 223602 (2022).
- [51] L. Du, Y.-T. Chen, Y. Zhang, and Y. Li, Giant atoms with time-dependent couplings, *Phys. Rev. Res.* **4**, 023198 (2022).
- [52] S. Terradas-Briansó, C. A. González-Gutiérrez, F. Nori, L. Martín-Moreno, and D. Zueco, Ultrastrong waveguide QED with giant atoms, *Phys. Rev. A* **106**, 063717 (2022).
- [53] A. Soro, C. S. Muñoz, and A. F. Kockum, Interaction between giant atoms in a one-dimensional structured environment, *Phys. Rev. A* **107**, 013710 (2023).
- [54] L. Du, L. Guo, Y. Zhang, and A. F. Kockum, Giant emitters in a structured bath with non-Hermitian skin effect, *Phys. Rev. Res.* **5**, L042040 (2023).
- [55] E. Raaholt Ingelsten, A. F. Kockum, and A. Soro, Avoiding decoherence with giant atoms in a two-dimensional structured environment, *Phys. Rev. Res.* **6**, 043222 (2024).
- [56] L. Leonforte, X. Sun, D. Valenti, B. Spagnolo, F. Illuminati, A. Carollo, and F. Ciccarello, Quantum optics with giant atoms in a structured photonic bath, *Quantum Sci. Technol.* **10**, 015057 (2025).
- [57] F. Roccati and D. Cilluffo, Controlling Markovianity with chiral giant atoms, *Phys. Rev. Lett.* **133**, 063603 (2024).
- [58] R.-Y. Gong, Z.-Y. He, C.-H. Yu, G.-F. Zhang, F. Nori, and Z.-L. Xiang, Tunable quantum router with giant atoms, implementing quantum gates, teleportation, non-reciprocity, and circulators, *arXiv:2411.19307*.
- [59] L. Du and A. F. Kockum, Unconventional and robust light-matter interactions based on the non-Hermitian skin effect, *Phys. Rev. Res.* **7**, 013140 (2025).
- [60] L. Du, X. Wang, A. F. Kockum, and J. Splettstoesser, Dressed interference in giant superatoms: Entanglement generation and transfer, *Phys. Rev. Lett.* **135**, 223601 (2025).
- [61] T. Levy-Yeyati, T. Ramos, and A. González-Tudela, Engineering giant transmon molecules as mediators of conditional two-photon gates, *Phys. Rev. Appl.* **25**, 054060 (2026).
- [62] J.-Q. Li, A. F. Kockum, and X. Wang, Generating spatially separated correlated multiphoton states in nonlinear waveguide quantum electrodynamics, *arXiv:2511.14281*.
- [63] W. Rieck, A. F. Kockum, and G. Chen, Doublon bound states in the continuum through giant atoms, *arXiv:2511.18212*.
- [64] R. Manenti, A. F. Kockum, A. Patterson, T. Behrle, J. Rahamim, G. Tancredi, F. Nori, and P. J. Leek, Circuit quantum acoustodynamics with surface acoustic waves, *Nat. Commun.* **8**, 975 (2017).
- [65] K. J. Satzinger, Y. P. Zhong, H.-S. Chang, G. A. Peairs, A. Bienfait, M.-H. Chou, A. Y. Cleland, C. R. Conner, É. Dumur, J. Grebel, I. Gutierrez, B. H. November, R. G. Povey, S. J. Whiteley, D. D. Awschalom, D. I. Schuster, and A. N. Cleland, Quantum control of surface acoustic-wave phonons, *Nature (London)* **563**, 661 (2018).
- [66] A. Bienfait, K. J. Satzinger, Y. P. Zhong, H.-S. Chang, M.-H. Chou, C. R. Conner, É. Dumur, J. Grebel, G. A. Peairs, R. G. Povey, and A. N. Cleland, Phonon-mediated quantum state transfer and remote qubit entanglement, *Science* **364**, 368 (2019).
- [67] G. Andersson, B. Suri, L. Guo, T. Aref, and P. Delsing, Non-exponential decay of a giant artificial atom, *Nat. Phys.* **15**, 1123 (2019).
- [68] A. Bienfait, Y. P. Zhong, H.-S. Chang, M.-H. Chou, C. R. Conner, É. Dumur, J. Grebel, G. A. Peairs, R. G. Povey, K. J. Satzinger, and A. N. Cleland, Quantum erasure using entangled surface acoustic phonons, *Phys. Rev. X* **10**, 021055 (2020).
- [69] G. Andersson, M. K. Ekström, and P. Delsing, Electromagnetically induced acoustic transparency with a superconducting circuit, *Phys. Rev. Lett.* **124**, 240402 (2020).
- [70] J. Hu, D. Li, Y. Qie, Z. Yin, A. F. Kockum, F. Nori, and S. An, Engineering the environment of a superconducting qubit with an artificial giant atom, *arXiv:2410.15377*.
- [71] A. Almanakly, B. Yankelevich, M. Hays, B. Kannan, R. Assouly, A. Greene, M. Gingras, B. M. Niedzielski, H. Stickler, M. E. Schwartz, K. Serniak, J. Î.-j. Wang, T. P. Orlando, S. Gustavsson, J. A. Grover, and W. D. Oliver, Deterministic remote entanglement using a chiral quantum interconnect, *Nat. Phys.* **21**, 825 (2025).
- [72] L. Xiao, B. Zhang, Y. Zeng, X. Pan, J.-Q. Wang, Z. Hua, H. Huang, Y. Xu, G. Xue, H. Yu, X.-B. Xu, W. Wang, C.-L. Zou, and L. Sun, Giant-atom quantum acoustodynamics in hybrid superconducting-phononic integrated circuits, *arXiv:2512.16582*.
- [73] G. Calajó, F. Ciccarello, D. Chang, and P. Rabl, Atom-field dressed states in slow-light waveguide QED, *Phys. Rev. A* **93**, 033833 (2016).
- [74] D. E. Chang, J. S. Douglas, A. González-Tudela, C.-L. Hung, and H. J. Kimble, *Colloquium: Quantum matter built from nanoscopic lattices of atoms and photons*, *Rev. Mod. Phys.* **90**, 031002 (2018).
- [75] A. González-Tudela and J. I. Cirac, Quantum emitters in two-dimensional structured reservoirs in the nonperturbative regime, *Phys. Rev. Lett.* **119**, 143602 (2017).
- [76] M. Scigliuzzo, G. Calajó, F. Ciccarello, D. Perez Lozano, A. Bengtsson, P. Scarlino, A. Wallraff, D. Chang, P. Delsing, and S. Gasparinetti, Controlling atom-photon bound states in an array of Josephson-junction resonators, *Phys. Rev. X* **12**, 031036 (2022).
- [77] V. Jouanny, S. Frasca, V. J. Weibel, I. Peyruchat, M. Scigliuzzo, F. Oppliger, F. De Palma, D. Sbroggiò, G. Beaulieu, O. Zilberberg, and P. Scarlino, High kinetic in-

- ductance cavity arrays for compact band engineering and topology-based disorder meters, *Nat. Commun.* **16**, 3396 (2025).
- [78] C. Castillo-Moreno, K. R. Amin, I. Strandberg, M. Kervinen, A. Osman, and S. Gasparinetti, Dynamical excitation control and multimode emission of an atom-photon bound state, *Phys. Rev. Lett.* **134**, 133601 (2025).
- [79] Y. Liu and A. A. Houck, Quantum electrodynamics near a photonic bandgap, *Nat. Phys.* **13**, 48 (2017).
- [80] N. M. Sundaresan, R. Lundgren, G. Zhu, A. V. Gorshkov, and A. A. Houck, Interacting qubit-photon bound states with superconducting circuits, *Phys. Rev. X* **9**, 011021 (2019).
- [81] P. M. Harrington, M. Naghiloo, D. Tan, and K. W. Murch, Bath engineering of a fluorescing artificial atom with a photonic crystal, *Phys. Rev. A* **99**, 052126 (2019).
- [82] K. H. Lim, W.-K. Mok, and L.-C. Kwek, Oscillating bound states in non-Markovian photonic lattices, *Phys. Rev. A* **107**, 023716 (2023).
- [83] S. Longhi, Photonic simulation of giant atom decay, *Opt. Lett.* **45**, 3017 (2020).
- [84] W. Zhao and Z. Wang, Single-photon scattering and bound states in an atom-waveguide system with two or multiple coupling points, *Phys. Rev. A* **101**, 053855 (2020).
- [85] S. Longhi, Rabi oscillations of bound states in the continuum, *Opt. Lett.* **46**, 2091 (2021).
- [86] H. Yu, Z. Wang, and J.-H. Wu, Entanglement preparation and nonreciprocal excitation evolution in giant atoms by controllable dissipation and coupling, *Phys. Rev. A* **104**, 013720 (2021).
- [87] C. Vega, M. Bello, D. Porras, and A. González-Tudela, Qubit-photon bound states in topological waveguides with long-range hoppings, *Phys. Rev. A* **104**, 053522 (2021).
- [88] H. Xiao, L. Wang, Z. Li, X. Chen, and L. Yuan, Bound state in a giant atom-modulated resonators system, *npj Quantum Inf.* **8**, 80 (2022).
- [89] W. Cheng, Z. Wang, and Y. X. Liu, Topology and retardation effect of a giant atom in a topological waveguide, *Phys. Rev. A* **106**, 033522 (2022).
- [90] X. Zhang, W. Cheng, Z. Gong, T. Zheng, and Z. Wang, Superconducting giant atom waveguide QED: Quantum Zeno and anti-Zeno effects in ultrastrong coupling regime, *arXiv:2205.03674*.
- [91] L. Du, Y.-T. Chen, Y. Zhang, Y. Li, and J.-H. Wu, Decay dynamics of a giant atom in a structured bath with broken time-reversal symmetry, *Quantum Sci. Technol.* **8**, 045010 (2023).
- [92] R. Bag and D. Roy, Quantum light-matter interactions in structured waveguides, *Phys. Rev. A* **108**, 053717 (2023).
- [93] W. Z. Jia and M. T. Yu, Atom-photon dressed states in a waveguide-QED system with multiple giant atoms, *Opt. Express* **32**, 9495 (2024).
- [94] Z.-M. Gao, J.-Q. Li, Z.-W. Li, W.-X. Liu, and X. Wang, Circuit QED with giant atoms coupling to left-handed superlattice metamaterials, *Phys. Rev. A* **109**, 013716 (2024).
- [95] A. Gilchrist, N. K. Langford, and M. A. Nielsen, Distance measures to compare real and ideal quantum processes, *Phys. Rev. A* **71**, 062310 (2005).
- [96] M.-D. Choi, Completely positive linear maps on complex matrices, *Linear Algebra Appl.* **10**, 285 (1975).
- [97] F. W. Strauch, P. R. Johnson, A. J. Dragt, C. J. Lobb, J. R. Anderson, and F. C. Wellstood, Quantum logic gates for coupled superconducting phase qubits, *Phys. Rev. Lett.* **91**, 167005 (2003).
- [98] W. E. Lamb and R. C. Retherford, Fine structure of the hydrogen atom by a microwave method, *Phys. Rev.* **72**, 241 (1947).
- [99] For $N = 100$, the Hamiltonian Hilbert-space dimension is approximately 5500, implying a Liouvillian dimension of $5500^2 = 3.025 \times 10^7$, beyond the reach of current exact-diagonalization techniques.
- [100] J. Koch, T. M. Yu, J. Gambetta, A. A. Houck, D. I. Schuster, J. Majer, A. Blais, M. H. Devoret, S. M. Girvin, and R. J. Schoelkopf, Charge-insensitive qubit design derived from the Cooper pair box, *Phys. Rev. A* **76**, 042319 (2007).
- [101] C. Mori, V. Milchakov, F. D'Esposito, L. Ruela, S. Kumar, V. N. Suresh, W. Arditi, D. Nicolas, G. Cappelli, A. Ranadive, G. L. Gal, M. Esposito, Q. Ficheux, N. Roch, T. Ramos, and O. Buisson, High-power readout of a transmon qubit using a nonlinear coupling, *arXiv:2507.03642*.
- [102] Z. Wang, R. W. Parker, E. Champion, and M. S. Blok, High- E_J/E_C transmon qubits with up to 12 levels, *Phys. Rev. Appl.* **23**, 034046 (2025).
- [103] S. J. de Graaf, S. H. Xue, B. J. Chapman, J. D. Teoh, T. Tsunoda, P. Winkel, J. W. O. Garmon, K. M. Chang, L. Frunzio, S. Puri, and R. J. Schoelkopf, A mid-circuit erasure check on a dual-rail cavity qubit using the joint-photon number-splitting regime of circuit QED, *npj Quantum Inf.* **11**, 1 (2025).
- [104] M. Kjaergaard, M. E. Schwartz, J. Braumüller, P. Krantz, J. I.-J. Wang, S. Gustavsson, and W. D. Oliver, Superconducting qubits: Current state of play, *Annu. Rev. Condens. Matter Phys.* **11**, 369 (2020).
- [105] A. P. M. Place, L. V. H. Rodgers, P. Mundada, B. M. Smitham, M. Fitzpatrick, Z. Leng, A. Premkumar, J. Bryon, A. Vrajitoarea, S. Sussman, G. Cheng, T. Madhavan, H. K. Babla, X. H. Le, Y. Gang, B. Jäck, A. Gyenis, N. Yao, R. J. Cava, N. P. de Leon, *et al.*, New material platform for superconducting transmon qubits with coherence times exceeding 0.3 milliseconds, *Nat. Commun.* **12**, 1779 (2021).
- [106] A. Somoroff, Q. Ficheux, R. A. Mencia, H. Xiong, R. Kuzmin, and V. E. Manucharyan, Millisecond coherence in a superconducting qubit, *Phys. Rev. Lett.* **130**, 267001 (2023).
- [107] Y. Kim, A. Eddins, S. Anand, K. X. Wei, E. van den Berg, S. Rosenblatt, H. Nayfeh, Y. Wu, M. Zaletel, K. Temme, and A. Kandala, Evidence for the utility of quantum computing before fault tolerance, *Nature (London)* **618**, 500 (2023).
- [108] J. Biznárová, A. Osman, E. Rehnman, L. Chayanun, C. Križan, P. Malmberg, M. Rommel, C. Warren, P. Delsing, A. Yurgens, J. Bylander, and A. Fadavi Roudsari, Mitigation of interfacial dielectric loss in aluminum-on-silicon superconducting qubits, *npj Quantum Inf.* **10**, 78 (2024).
- [109] S. Kono, J. Pan, M. Chegnizadeh, X. Wang, A. Youssefi, M. Scigliuzzo, and T. J. Kippenberg, Mechanically induced correlated errors on superconducting qubits with relaxation times exceeding 0.4 ms, *Nat. Commun.* **15**, 3950 (2024).
- [110] M. Bal, A. A. Murthy, S. Zhu, F. Crisa, X. You, Z. Huang, T. Roy, J. Lee, D. v. Zanten, R. Pilipenko, I. Nekrashevich, A. Lunin, D. Bafia, Y. Krasnikova, C. J. Kopas, E. O. Lachman, D. Miller, J. Y. Mutus, M. J. Reagor, H. Cansizoglu, *et al.*, Systematic improvements in transmon qubit coherence enabled

- by niobium surface encapsulation, [npj Quantum Inf.](#) **10**, 43 (2024).
- [111] R. Acharya, D. A. Abanin, L. Aghababaie-Beni, I. Aleiner, T. I. Andersen, M. Ansmann, F. Arute, K. Arya, A. Asfaw, N. Astrakhantsev, J. Atalaya, R. Babbush, D. Bacon, B. Ballard, J. C. Bardin, J. Bausch, A. Bengtsson, A. Bilmes, S. Blackwell, S. Boixo, *et al.*, Quantum error correction below the surface code threshold, [Nature \(London\)](#) **638**, 920 (2025).
- [112] S. Poorgholam-Khanjari, V. Seferai, P. Foshat, C. Rose, H. Feng, R. H. Hadfield, M. Weides, and K. Delfanazari, Engineering high-Q superconducting tantalum microwave coplanar waveguide resonators for compact coherent quantum circuit, [Sci. Rep.](#) **15**, 27113 (2025).
- [113] T. Abad, J. Fernández-Pendás, A. Frisk Kockum, and G. Johansson, Universal fidelity reduction of quantum operations from weak dissipation, [Phys. Rev. Lett.](#) **129**, 150504 (2022).
- [114] T. Abad, Y. Schattner, A. F. Kockum, and G. Johansson, Impact of decoherence on the fidelity of quantum gates leaving the computational subspace, [Quantum](#) **9**, 1684 (2025).
- [115] T. Thorbeck, A. McDonald, O. Lanes, J. Blair, G. Keefe, A. A. Stabile, B. Royer, L. C. G. Govia, and A. Blais, High-fidelity gates in a transmon using bath engineering for passive leakage reset, [arXiv:2411.04101](#).
- [116] X. Wang, J.-Q. Li, Z. Wang, A. F. Kockum, L. Du, T. Liu, and F. Nori, Nonlinear cascaded quantum network with giant emitters, [Commun. Phys.](#) **9**, 205 (2026).
- [117] <https://github.com/Q-WAVES/GiantAtoms/>.