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Heralded quantum controlled phase gates with dissipative dynamics in macroscopically-distant resonators

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Heralded near-deterministic multi-qubit controlled phase gates with integrated error detection have recently been proposed by Borregaard *et al.* [Phys. Rev. Lett. **114**, 110502 (2015)]. This protocol is based on a single four-level atom (a heralding quartit) and N three-level atoms (operational qutrits) coupled to a single-resonator mode acting as a cavity bus. Here we generalize this method for two distant resonators without the cavity bus between the heralding and operational atoms. Specifically, we analyze the two-qubit controlled-Z gate and its multi-qubit-controlled generalization (i.e., a Toffoli-like gate) acting on the two-lowest levels of N qutrits inside one resonator, with their successful actions being heralded by an auxiliary microwave-driven quartit inside the other resonator. Moreover, we propose a circuit-quantum-electrodynamics realization of the protocol with flux and phase qutrits in linearly-coupled transmission-line resonators with dissipation. These methods offer a quadratic fidelity improvement compared to cavity-assisted deterministic gates.

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I. INTRODUCTION

The ability to carry out quantum gates is one of the central requirements for a functional quantum computer [1]. For this reason, quantum gates have been extensively studied in theoretical and experimental in various systems. These include (for reviews see [2, 3] and references therein): superconducting circuits, trapped ions, diamond nitrogen-vacancy centers, semiconductor nanostructures, or linear-optical setups. Despite of such substantial efforts, the environment-induced decoherence still represents a major hurdle in the quest for perfect quantum gates. To protect quantum systems from decoherence, three basic approaches have been explored; namely, quantum error correction, dynamical decoupling, and decoherence-free subspaces (for reviews see [4, 5] and references therein). Such approaches try to cope with decoherence, and thus, to prevent the leakage of quantum information from a quantum system into its environment.

Alternatively, there exists a very distinct way, where decoherence acts as a resource, rather than as a traditional noise source [6–10]. Recent progress in treating open quantum systems has yielded an effective operator formalism [11, 12]. As in Ref. [12], (i) when the interactions between the ground- and excited-state subspaces of an open system initially in its ground state are sufficiently weak (so it can be considered as perturbations of the two subspaces) and also (ii) when the interactions inside the ground-state subspace are much smaller than those inside the excited-state subspace, one can adiabatically eliminate these excited states in the presence of both unitary and dissipative dynamics, and obtain an effective master equation containing the effective Hamil-

tonian, as well as the effective Lindblad operators, associated only with the ground states. In addition to reducing the complexity of the time evolution for an open quantum system, this approximation treatment is applicable to the explorations of decay processes, hence may lead to a better performance than that in the case of relying upon coherent-unitary dynamics.

So far, such a formalism has been widely used for dissipative entanglement preparation [11, 13–18], quantum phase estimation [19] and other applications [20–23]. In particular, Borregaard *et al.* presented a heralded near-deterministic method for quantum gates in a single optical cavity [24], with a significant improvement in the error scaling, compared to deterministic cavity-based gates. However, in order to carry out scalable quantum information processing, a distant herald for quantum gates in coupled resonators is of a great importance concerning both experimental feasibility and fundamental tests of quantum mechanics.

The goal of this paper is to propose and analyze an approach to heralding controlled quantum gates in two macroscopically distant resonators, by generalizing the single-resonator method of Borregaard *et al.* [24]. We use an auxiliary quartit atom, which is located in one resonator and driven by two coherent fields, to distantly herald controlled phase gates acting on the two lowest levels of qutrit atoms, which are located in the other resonator. These controlled phase operations studied here include the two-qubit controlled-Z gate (controlled-sign gate) and its multi-qubit-controlled version referred to as a Toffoli-like gate. We also propose a realization of these gates using superconducting artificial atoms in a dissipative circuit quantum electrodynamics (QED) system.

Circuit-QED systems, which are composed of superconducting artificial atoms coupled to superconducting resonators, offer promising platforms for quantum engineering and quantum-information processing [2, 25–29]. Although artificial and natural atoms are similar in various properties including, for example, discrete anharmonic energy levels [30–32], superconducting atoms have some substantial advantages over natural atoms. These include: (i) The spacing between energy levels, decoherence rates, and coupling strengths between different circuit elements are tunable by adjusting external parameters, thus, providing flexibility for quantum information processing [2, 3]; (ii) The potential energies of natural atoms have an inversion symmetry, leading to a well-defined parity symmetry for each eigenstate. Thus, an allowable dipole transition requires a parity change, and can only occur between the two eigenstates having different parities. However, the potential energies of superconducting artificial atoms can be controlled [33] by external parameters and, in turn, the inversion symmetry can be broken or unbroken. When the symmetry is unbroken, each eigenstate of an artificial atom can have a well-defined parity symmetry and, thus, exhibit a transition behavior similar to natural atoms. But when the symmetry is broken then there is no well-defined parity symmetry for each eigenstate and, thus, the microwave-induced single-photon transition between any two eigenstates can be possible. That is, the selection rules can easily be modified [34]; (iii) The couplings between different circuit elements can be strong, ultrastrong (for reviews see [3, 32]), or even deep strong [35], which, in particular, enable efficient state preparation within a short time and with a high fidelity.

Specifically, based on circuit QED we realize a distant herald for a multi-qubit Toffoli-like gate, whose the fidelity scaling is quadratically improved as compared to unitary-dynamics-based gates in cavities [24, 36]; moreover, heralded controlled-Z gates even with more significant improvements can also be achieved by using single-qubit operations. Note that the conditional measurement on the heralding atom is performed to remove the detected errors from the quantum gate. Thus, these detected errors can only reduce the success probability but not the gate fidelity. When the gate is successful, the gate fidelity can still be very high as limited only by the undetected errors. In contrast to previous works, a macroscopically distant herald for quantum gates is proposed in superconducting circuits via the combined effect of the dissipative and unitary dynamics, thereby having the potential advantage of high efficiency in experimental scenarios. By considering only the d lowest energy levels of a superconducting artificial atom, one can use them as a logical qudit for performing quantum operations. In special cases, one can operate with a qubit (for $d = 2$), qutrit (for $d = 3$), or quartit (ququart, for $d = 4$). Here we will analyze these special qudits.

This paper is organized as follows. In Sec. II, we describe a physical model for heralded quantum phase gates

and propose a superconducting circuit for its realization. In Sec. III, we derive the effective master equation which is used in Sec. IV to realize a multi-qubit Toffoli-like gate. Working with single-qubit operations, Sec. V then presents a heralded controlled-Z gate. The last section is our summary. A detailed derivation of the effective master equation, which is applied in Sec. III, is given in the Appendix.

II. PHYSICAL MODEL IN SUPERCONDUCTING CIRCUITS

The basic idea underlying our protocol is schematically illustrated in Figs. 1 and 2. Figure 1 shows our proposal of the circuit-QED implementation of the protocol based on superconducting qudits, while the corresponding energy-level diagrams of the qudits are depicted in Fig. 2. Specifically, we consider two superconducting transmission-line resonators A and B , connected by a capacitor [39]. The coupling Hamiltonian, of strength J , for the two resonators can be expressed as (hereafter we set $\hbar = 1$)

$$H_{AB} = J(a_A a_B^\dagger + a_A^\dagger a_B), \quad (1)$$

where a_A (a_B) is the annihilation operator of the resonator A (B). We assume that superconducting artificial atoms, which are treated as qudits (i.e., d -level quantum systems) [38, 40], are coupled to the resonators. Specifically, a superconducting phase quartit is directly confined inside the resonator A , and is used as an auxiliary quartit atom to herald the success of quantum gates. Such a quartit consists of two ground levels $|g_1\rangle$ and $|g_2\rangle$, as well as two excited levels $|e_1\rangle$ and $|e_2\rangle$, depicted in Fig. 1(b). Because the potential energy of the phase quartit is always broken, the quartit levels have no well-defined parity symmetry [38, 41, 42]. Thus, we can couple any two levels by applied fields. We assume that the transition between $|g_2\rangle$ and $|e_2\rangle$ is coupled to the resonator A by an inductance, with a Hamiltonian

$$H_A = g_A (a_A |e_2\rangle \langle g_2| + \text{H.c.}), \quad (2)$$

where $g_A > 0$ is a coupling strength.

In a similar manner, we couple the resonator B to N Λ -type qutrits, for example, superconducting flux three-level atoms [33, 43]. Each qutrit consists of two ground levels $|0\rangle$ and $|1\rangle$, together with one excited level $|2\rangle$, depicted in Fig. 2(c). The lowest two levels of an atomic qutrit are treated as qubit states. With current technologies, superconducting atoms can be made almost identical. Thus, for simplicity, we can assume that these qutrits are identical and have the same coupling, of strength g_B , to the resonator B . Such a coupling can be ensured by adjusting the control signals on the qutrits and by tuning the separation between any two qutrits to be much smaller than the wavelength of the resonator B , respec-

TABLE I. Basic notations used in this paper. Here, $x = 1, 2$ and $z = 0, 1, 2$.

notation	Meaning
$\omega_{g(e),x}$	$ g_x\rangle$ -($ e_x\rangle$)- level frequency
ω_z	$ z\rangle$ -level frequency
ω_c	Common resonance frequency of resonators A and B
$\omega_{m,x}$	Microwave drive x frequency
Ω_x	Microwave x driving amplitude strength
g_A (g_B)	Coupling strength between the qutrit (qutrit) atom and resonator A (B)
J	Inter-resonator coupling strength
$\gamma_{g,x}$	Decay rate from level $ e_x\rangle$ to $ g_x\rangle$
γ_{x-1}	Decay rate from level $ 2\rangle$ to $ x-1\rangle$
γ	Total decay rate, $\gamma = \gamma_0 + \gamma_1$, of each qutrit atom
κ	Resonator decay rate
C_B	Atom-resonator cooperativity, $C_B = g_B^2 / (\kappa\gamma)$,
$\Delta_{e,1}$	Microwave detuning, $\Delta_{e,1} = \omega_{e,1} - \omega_{g,1} - \omega_{m,1}$
$\Delta_{e,2}$	Microwave detuning, $\Delta_{e,2} = \omega_{e,2} - \omega_{g,1} - \omega_{m,1} - \omega_{m,2}$
δ	Resonator detuning, $\delta = \omega_2 - \omega_1 - \omega_{m,1} - \omega_{m,2} - \omega_{g,1} + \omega_{g,2}$

tively. The corresponding Hamiltonian is

$$H_B = g_B \sum_{k=1}^N (a_B |2\rangle_k \langle 1| + \text{H.c.}). \quad (3)$$

where k labels the qutrits and $g_B > 0$ is the coupling strength of the resonator B . Note that the direct dipole-coupling between the qutrits has been neglected, owing to their large spatial separations. Nevertheless, these qutrits can interact indirectly via the common field a_B of the resonator B , analogously to the model of three-level quantum dots in a resonator studied for quantum-information processing in Refs. [44, 45]. Because the Josephson junctions are nonlinear circuit elements [32], and therefore the resulting levels are highly anharmonic compared to the driving strengths as well as to the atom-resonator coupling strengths, all transitions of the qutrit and the qutrits can be driven or coupled separately by the control lines, as shown in Fig. 2(a).

We also assume that a microwave field at the frequency $\omega_{m,1}$ drives the $|g_1\rangle \leftrightarrow |e_1\rangle$ transition, with a driving strength Ω_1 and at the same time, the excited states $|e_1\rangle$ and $|e_2\rangle$ are also coupled by means of a microwave field at the frequency $\omega_{m,2}$, with a coupling strength Ω_2 . The interaction Hamiltonian describing the effect of these external drives reads as

$$H_D = \frac{1}{2} (\Omega_1 e^{i\omega_{m,1}t} |g_1\rangle \langle e_1| + \Omega_2 e^{i\omega_{m,2}t} |e_1\rangle \langle e_2| + \text{H.c.}). \quad (4)$$

Define $\omega_{g,x}$, $\omega_{e,x}$, and ω_z as the frequencies of the atomic levels $|g_x\rangle$, $|e_x\rangle$, and $|z\rangle$, respectively, with $x = 1, 2$ and $z = 0, 1, 2$. Thus, the total Hamiltonian for our system is

$$H_T = H_0 + H_A + H_B + H_{AB} + H_D, \quad (5)$$

where

$$H_0 = \sum_{x=1,2} (\omega_{g,x} |g_x\rangle \langle g_x| + \omega_{e,x} |e_x\rangle \langle e_x|) + \sum_{k=1}^N \sum_{z=0,1,2} \omega_z |z\rangle_k \langle z| + \omega_c (a_A^\dagger a_A + a_B^\dagger a_B), \quad (6)$$

is the free Hamiltonian.

Upon introducing the symmetric and antisymmetric optical modes, $a_\pm = (a_A \pm a_B) / \sqrt{2}$, the total Hamiltonian in a proper rotating frame reads $H_T = H_e + V + V^\dagger$, with

$$H_e = \sum_{k=1}^N \left\{ \delta |2\rangle_k \langle 2| + \frac{g_B}{\sqrt{2}} [(a_+ - a_-) |2\rangle_k \langle 1| + \text{H.c.}] \right\} + \Delta_{e,1} |e_1\rangle \langle e_1| + \Delta_{e,2} |e_2\rangle \langle e_2| + 2J a_+^\dagger a_+ + \frac{g_A}{\sqrt{2}} [(a_+ + a_-) |e_2\rangle \langle g_2| + \text{H.c.}] + \frac{\Omega_2}{2} (|e_2\rangle \langle e_1| + \text{H.c.}), \quad (7)$$

and

$$V = \frac{\Omega_1}{2} |e_1\rangle \langle g_1|. \quad (8)$$

Note that we have applied the rotating-wave approximation (RWA), which directly drops the fast oscillating terms of the total Hamiltonian. The detunings are defined as (see Fig. 2):

$$\delta = \omega_2 - \omega_1 - \omega_{m,1} - \omega_{m,2} - \omega_{g,1} + \omega_{g,2}, \quad (9)$$

$$\Delta_{e,1} = \omega_{e,1} - \omega_{g,1} - \omega_{m,1}, \quad (10)$$

$$\Delta_{e,2} = \omega_{e,2} - \omega_{g,1} - \omega_{m,1} - \omega_{m,2}, \quad (11)$$

In Eq. (7) we have assumed

$$\omega_c = \omega_2 - \delta - \omega_1 + J, \quad (12)$$

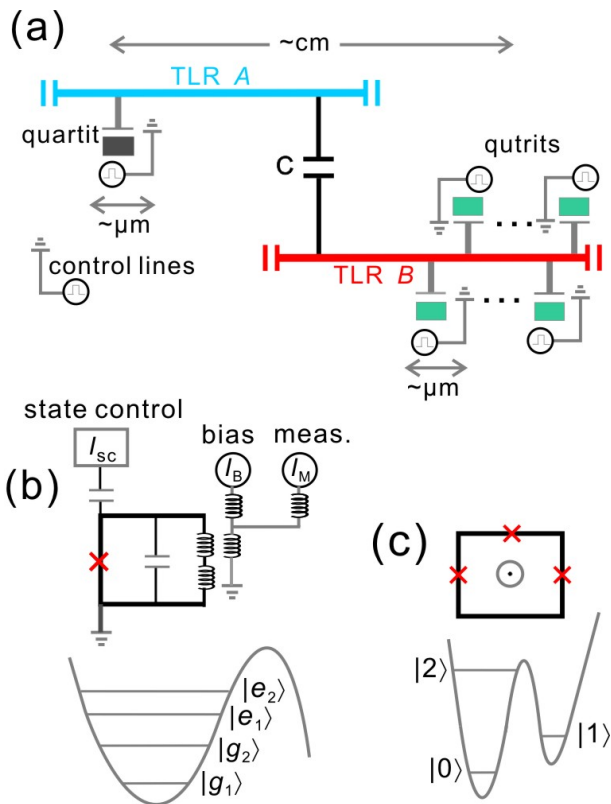


FIG. 1. (Color online) (a) Schematic diagram of a superconducting circuit layout, which shows our implementation of a heralded near-deterministic controlled multi-qubit Toffoli-like gate and, in a special case, the two-qubit controlled-Z (CZ) gate. Two transmission-line resonators (TLRs), labeled by A and B , are linearly coupled via a capacitor. Resonator A is coupled to a phase quartit (e.g., a four-level qudit) via a capacitance, while resonator B is coupled to N identical flux qutrits (e.g., three-level qudits). Such circuit elements can be controlled via ac and dc signals through the control lines. The distance between the qudit and qutrits can be of order of cm due to the macroscopic length of the resonators. Panel (b) shows an energy-level diagram for a prototypical phase quartit, while panel (c) depicts that diagram for a typical flux qutrit. Here we assume that the energy potential for the quartit is cubic and for the qutrit is an asymmetric double well. Note that both flux and phase qudits can be tuned (for example, by adjusting the flux bias in the qudit loop) to obtain exactly two, three, or four levels. For example, following the method of Refs. [37, 38], the phase qudit state can be controlled by an ac signal I_{sc} , while a dc signal I_B is used to bias the circuit. Moreover, a short measurement pulse I_M is applied to decrease the barrier in the energy potential well of panel (b) to enable given upper states to tunnel out of the well. These tunneling currents can be detected by another SQUID which, for brevity, is not plotted here. Thus, the proposed quantum gates can be realized, based solely on flux or phase qudits, or other superconducting qudits. The successful operation of this Toffoli-like gate is heralded by the detection of the quartit in its ground state $|g_1\rangle$.

such that the three-photon Raman transition between the levels $|g_1\rangle \leftrightarrow |e_1\rangle \leftrightarrow |e_2\rangle \leftrightarrow |g_2\rangle$, is resonant when mediated by the antisymmetric mode a_- , but is detuned by $2J$ when mediated by the symmetric mode a_+ . We further assume that $|e_1\rangle$ and $|e_2\rangle$ decay to $|g_1\rangle$ and $|g_2\rangle$, respectively, with rates $\gamma_{g,1}$ and $\gamma_{g,2}$, and for each qutrit atom, $|2\rangle$ can decay either to $|0\rangle$ with a rate γ_0 or to $|1\rangle$ with a rate γ_1 . In addition, both resonators are assumed to have the same decay rate κ . All basic symbols used in this paper are shown in Table I.

III. MASTER EQUATION

Here we present a standard approach based on the master equation in the Lindblad form to study the dissipative dynamics of our system. The master equation is valid under a few fundamental assumptions which include [46–48]: (A1) the approximation of the weak coupling between the analyzed system and its reservoir (environment), (A2) the Markov approximation, and (A3) the secular approximation. The approximations (A1) and (A2) are often referred to as the Born-Markov approximation. By applying (A2), one assumes that the environmental-memory effects are short-lived, such that the system evolution depends only on its present state. This approximation is valid if the environmental correlation time (which can be evaluated by the decay timescale of the two-time correlation functions of the environmental operators coupled to the system) is much shorter than a typical system-evolution timescale over which the system experiences a significant evolution. For example, the Born-Markov approximation is valid if an environment is large and weakly coupled to a system. The approximation (A3) is applied to cast a given Markovian master equation into the Lindblad form. This corresponds to ignoring fast-oscillating terms in the master equation based on (A1) and (A2). Thus, (A3) is sometimes called the RWA, although it is usually applied at the level of a given master equation, and not necessarily at the level of the system-reservoir interaction Hamiltonian. There are numerous references showing the excellent agreement between the experimental and theoretical results based upon the master equation in the Lindblad form describing the lossy interaction of quantum systems (including superconducting artificial atoms) and resonator fields (see, e.g., [49, 50] and references therein). The validity of these approximations for a single-qudit version of our system was also experimentally analyzed in Ref. [15].

The standard master equation in the Lindblad form for the system described by the Hamiltonian given in Eq. (5), assuming the zero-temperature environment (bath), can be given by [46–48]:

$$\dot{\rho}_T(t) = i[\rho_T(t), H_T] + \frac{1}{2} \sum_j \left[2L_j \rho_T(t) L_j^\dagger - \rho_T(t) L_j^\dagger L_j - L_j^\dagger L_j \rho_T(t) \right], \quad (13)$$

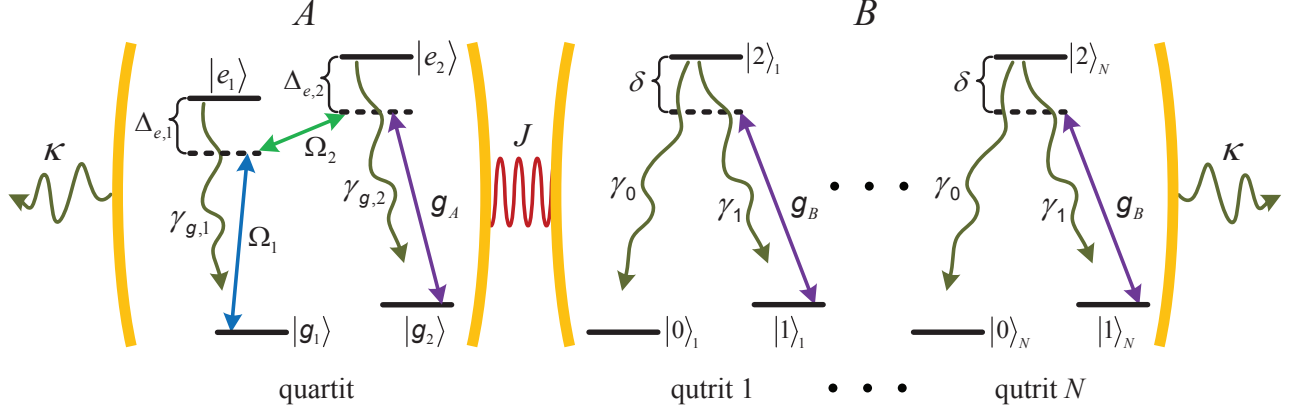


FIG. 2. (Color online) Energy-level diagram showing allowed transitions and couplings in the circuit-QED system depicted in Fig. 1. Two single-mode resonators, labeled by A and B are coupled with a strength J . An auxiliary, quartit atom, which acts as a heralding device, is confined in the resonator A . Two microwave fields drive off-resonantly the transitions $|g_1\rangle \rightarrow |e_1\rangle$ and $|e_1\rangle \rightarrow |e_2\rangle$ of the auxiliary quartit, with strengths Ω_1 and Ω_2 , respectively. Moreover, the states $|g_2\rangle$ and $|e_2\rangle$ are coupled by the resonator mode a_A with strength g_A . In the resonator B , there are confined N qutrit atoms, for which the two lowest-energy levels can be treated as qubits. For each qutrit atom, only the state $|1\rangle$ is coupled to $|2\rangle$ by the resonator mode a_B with strength g_B . Upon restricting our discussion to sufficiently weak microwave drive Ω_1 , we could adiabatically eliminate the excited states of the total system to yield an effective Lindblad-type master equation, which involves the ground states only. The resulting dynamics allows for the realization of the controlled-phase gates, which can successfully occur if the auxiliary quartit in the state $|g_1\rangle$ is measured.

where $\rho_T(t)$ is the density operator of the total system. The Lindblad operators associated with the resonator decay and atomic spontaneous emission are accordingly given by

$$\begin{aligned} L_{\pm} &= \sqrt{\kappa} a_{\pm}, \\ L_{g,x} &= \sqrt{\gamma_{g,x}} |g_x\rangle \langle e_x|, \\ L_{k,x-1} &= \sqrt{\gamma_{x-1}} |x-1\rangle_k \langle 2|, \end{aligned} \quad (14)$$

where k labels the qutrit atoms, and $x = 1, 2$.

We now consider a weak microwave drive Ω_1 , so that $\{\Omega_1/\Delta_{e,1}, \Omega_1/g_{A(B)}\} \ll 1$. In this situation, both the resonator modes and excited atomic states can be adiabatically eliminated if the system is initially in its ground state. Following the procedure in Ref. [12], the dynamics is therefore described by the effective Hamiltonian,

$$H_{\text{eff}} = -\frac{1}{2} V^\dagger \left[H_{\text{NH}}^{-1} + (H_{\text{NH}}^{-1})^\dagger \right] V, \quad (15)$$

and the effective Lindblad operators,

$$L_{\text{eff}}^j = L_j H_{\text{NH}}^{-1} V. \quad (16)$$

Here,

$$H_{\text{NH}} = H_e - \frac{i}{2} \sum_j L_j^\dagger L_j \quad (17)$$

accounts for the no-jump Hamiltonian, where the sum runs over all dissipative processes as mentioned in Eq. (14). The effective Lindblad master equation then

has the form

$$\begin{aligned} \dot{\rho}(t) &= i [\rho(t), H_{\text{eff}}] + \frac{1}{2} \sum_j \left\{ 2L_{\text{eff}}^j \rho(t) (L_{\text{eff}}^j)^\dagger \right. \\ &\quad \left. - \left[(L_{\text{eff}}^j)^\dagger L_{\text{eff}}^j \rho(t) + \rho(t) (L_{\text{eff}}^j)^\dagger L_{\text{eff}}^j \right] \right\}, \end{aligned} \quad (18)$$

assuming the reservoir at zero temperature. Here $\rho(t)$ is the density operator of the quartit and qutrit atoms.

As explained in detail in the Appendix, when working within the limits $\Omega_2 \ll \Delta_{e,1}$ and $\kappa \ll J$, we can more explicitly obtain the effective Hamiltonian

$$H_{\text{eff}} = |g_1\rangle \langle g_1| \otimes \sum_{n=0}^N \Delta_n \mathcal{P}_n, \quad (19)$$

where \mathcal{P}_n represents a projector onto a subspace characterized by n atomic qutrits in $|1\rangle$, and

$$\Delta_n = -\frac{\tilde{\Omega}^2}{4\gamma} \text{Re} \left\{ \frac{1}{\mathcal{Z}_n} \left(i\tilde{\delta} + nC_B \right) \right\} \quad (20)$$

refers to an AC stark shift with

$$\mathcal{Z}_n = i\tilde{\delta}\tilde{\Delta}_{e,2} + C_B \left(\alpha\tilde{\delta} + n\tilde{\Delta}_{e,2} \right) - n\alpha C_B^2/G. \quad (21)$$

Here, we have defined the overall decay rate $\gamma = \gamma_0 + \gamma_1$, of each qutrit atom, the effective drive $\tilde{\Omega} = \Omega_1\Omega_2/(2\Delta_{e,1})$, and the following dimensionless quantities: the atom-resonator cooperativity $C_B = g_B^2/(\kappa\gamma)$,

the strength $G = J/\kappa$, the complex detunings $\tilde{\delta} = \delta/\gamma - i/2$, $\tilde{\Delta}_{e,2} = \Delta_{e,2}/\gamma - i\beta/2$, and the parameters

$$\alpha = (g_A/g_B)^2, \quad \beta = \gamma_{g,2}/\gamma. \quad (22)$$

In all our numerical simulations we set $\alpha = \beta = 1$. The term, $-\Omega_1^2/(4\Delta_{e,1})$, of Δ_n has been removed since it only causes an overall phase. Equation (19) indicates that the time evolution under the effective Hamiltonian gives rise to a dynamical phase imprinted onto each state of the qutrits, while making the state of the quartit atom unchanged. Correspondingly, the effective Lindblad operators are found to be

$$L_{\text{eff}}^{\pm} = |g_2\rangle\langle g_1| \otimes \sum_{n=0}^N r_{\pm,n} \mathcal{P}_n, \quad (23)$$

$$L_{\text{eff}}^{g,x} = |g_x\rangle\langle g_1| \otimes \sum_{n=0}^N r_{g_x,n} \mathcal{P}_n, \quad (24)$$

$$L_{\text{eff}}^{k,x-1} = |g_2\rangle\langle g_1| \otimes \sum_{n=1}^N r_{x-1,n} |x-1\rangle_k \langle 1| \mathcal{P}_n, \quad (25)$$

with $x = 1, 2$. Here the effective decay rates $r_{\pm,n}$, $r_{g_x,n}$, and $r_{x-1,n}$ are expressed, respectively, as

$$\begin{aligned} r_{+,n} &= \frac{\tilde{\Omega}\sqrt{\alpha C_B}}{4G\mathcal{Z}_n\sqrt{2\gamma}} \left(i\tilde{\delta} + 2nC_B \right), \\ r_{-,n} &= -\frac{\tilde{\Omega}\sqrt{\alpha C_B}}{\mathcal{Z}_n\sqrt{2\gamma}} \left[i\tilde{\delta} - nC_B/(2G) \right], \\ r_{g_1,n} &= \frac{\Omega_1\sqrt{\gamma_{g,1}}}{2\Delta_{e,1}}, \\ r_{g_2,n} &= -\frac{\tilde{\Omega}\sqrt{\beta}}{2\mathcal{Z}_n\sqrt{\gamma}} \left(i\tilde{\delta} + nC_B \right), \\ r_{x-1,n} &= -\frac{\tilde{\Omega}\sqrt{\alpha\gamma_{x-1}}}{2\gamma\mathcal{Z}_n} C_B, \end{aligned} \quad (26)$$

as derived in the Appendix. According to these effective Lindblad operators, we find that all dissipative processes, except the one corresponding to $L_{\text{eff}}^{g,1}$ independent of n , cause the $|g_1\rangle \rightarrow |g_2\rangle$ decay. The errors induced by this decay are detectable by measuring the state of the quartit atom, and can be removed by conditioning on the quartit atom being measured in $|g_1\rangle$. Upon solving the effective master equation of Eq. (18), the probability of detecting the quartit atom in the state $|g_1\rangle$ is given by

$$P = \sum_{n=0}^N \text{Tr} [(|g_1\rangle\langle g_1| \otimes \mathcal{P}_n) \rho(t)], \quad (27)$$

where Tr is the trace operation over the subspace spanned by the ground states of the quartit and qutrit atoms, and $\sum_{n=0}^N \mathcal{P}_n = \mathcal{I}_G$ is the identity operator acting on the ground state manifold of qutrit atoms. Such a detection could be performed using the method developed in

Refs. [37, 38] for a phase qutrit, which we have briefly described in the caption of Fig. 1. We note that also other schemes for the detection of phase and flux qutrit states can be applied including the dispersive read-out methods [51–53]. After this measurement, the conditional density operator of the qutrits is then

$$\begin{aligned} \rho_{\text{qutrit}}(t) &= \frac{1}{P} \sum_{n,n'=0}^N e^{-i(\Delta_n - \Delta_{n'})t} e^{-(\Gamma_n + \Gamma_{n'})t/2} \\ &\quad \times \mathcal{P}_n [\langle g_1 | \rho(0) | g_1 \rangle] \mathcal{P}_{n'}, \end{aligned} \quad (28)$$

with the following total decay rate:

$$\Gamma_n = |r_{+,n}|^2 + |r_{-,n}|^2 + |r_{g_2,n}|^2 + n (|r_{0,n}|^2 + |r_{1,n}|^2). \quad (29)$$

Thus, by measuring the quartit atom and referring to the P as the success probability, we could realize heralded quantum controlled phase gates as discussed below. In this approach, the detectable errors only reduce the success probability, instead of reducing the gate fidelity. If we successfully measure the quartit in $|g_1\rangle$, then the resulting gate has a very high fidelity, which is limited only by the undetectable errors induced, for example, by the differences between the decay rates Γ_n for different states of the qutrit atoms. Thus, the underlying key idea is to remove the detectable errors by a heralding measurement and, then, to achieve quantum gates with very high fidelities in two macroscopically-distant resonators.

TABLE II. The action of a heralded multi-qubit Toffoli-like gate. Note that $\Delta_0 \neq \Delta_1 = \Delta_2 = \dots = \Delta_N$, and we have also ignored an overall phase, $\exp(-i\pi)$, of the output states.

n	Input states	Time evolution \rightarrow	Output states
0	$ 000\dots 0\rangle$	$\xrightarrow{\exp(-i\Delta_0 t_{\text{Toff}})}$	$- 000\dots 0\rangle$
	$ 100\dots 0\rangle$		$ 100\dots 0\rangle$
	$ 010\dots 0\rangle$		$ 010\dots 0\rangle$
1	\vdots	$\xrightarrow{\exp(-i\Delta_1 t_{\text{Toff}})}$	\vdots
	$ 0\dots 001\rangle$		$ 0\dots 001\rangle$
	$ 110\dots 0\rangle$		$ 110\dots 0\rangle$
2	$ 101\dots 0\rangle$	$\xrightarrow{\exp(-i\Delta_2 t_{\text{Toff}})}$	$ 101\dots 0\rangle$
	\vdots		\vdots
	$ 0\dots 011\rangle$		$ 0\dots 011\rangle$
\vdots	\vdots	\vdots	\vdots
N	$ 111\dots 1\rangle$	$\xrightarrow{\exp(-i\Delta_N t_{\text{Toff}})}$	$ 111\dots 1\rangle$

IV. HERALDED MULTI-QUBIT TOFFOLI-LIKE GATE

In this section we will demonstrate a heralded near-deterministic multi-qubit Toffoli-like gate, which is defined as the multi-qubit controlled-Z gate. Thus,

TABLE III. The action of the heralded controlled-Z gate. Note that Δ_0 , Δ_1 , and Δ_2 are not equal to each other. Thus, some unitary operations should be applied on qubits to achieve their proper phase evolution.

n	Input states	Time evolution \rightarrow	Output states
0	$ 00\rangle$	$\exp(-i\Delta_0 t_{CZ})$	$ 00\rangle$
1	$ 10\rangle$	$\exp(-i\Delta_1 t_{CZ})$	$ 10\rangle$
	$ 01\rangle$	$\exp(-i\Delta_1 t_{CZ})$	$ 01\rangle$
2	$ 11\rangle$	$\exp(-i\Delta_2 t_{CZ})$	$- 11\rangle$

the action of our Toffoli-like gate on N qubits in a state $|\psi\rangle = |q_1, q_2, \dots, q_N\rangle$ is given by $|\psi\rangle \rightarrow (-1)^{(q_1 \oplus 1)(q_2 \oplus 1) \dots (q_N \oplus 1)} |\psi\rangle$. In our case these logical qubits correspond to the two lowest states of the qutrits in the resonator B . The successful action of this Toffoli-like gate is heralded by the detection of the quartit in its ground state $|g_1\rangle$. Note that the standard three-qubit Toffoli gate is defined as the double-controlled NOT (CC-NOT) gate and given by the map [1]: $|q_1, q_2, q_3\rangle \rightarrow |q_1, q_2, q_3 \oplus q_1 q_2\rangle$, rather than the map $|q_1, q_2, q_3\rangle \rightarrow (-1)^{(q_1 \oplus 1)(q_2 \oplus 1)(q_3 \oplus 1)} |q_1, q_2, q_3\rangle$, which is applied here as in, e.g., Ref. [24].

In order to realize our multi-qubit Toffoli-like gate, we can make $\Delta_{n>0}$ independent of n , and ensure that $|\Delta_0| \ll |\Delta_{n>0}|$. To this end, we tune $\delta = 0$ and $\Delta_{e,2}/\gamma = \alpha C_B (R + 1/G)$, where

$$R = \sqrt{\frac{1}{2} \left(\frac{1}{G^2} + \frac{\beta}{\alpha C_B} + \frac{1}{C_B} \right)}. \quad (30)$$

In the limit of $\min\{G, C_B\} \gg 1$, this choice can lead to

$$\Delta_0 = -\frac{\tilde{\Omega}^2}{4\gamma} \frac{1}{\alpha C_B} (R + 1/G), \quad (31)$$

$$\Delta_{n>0} = -\frac{\tilde{\Omega}^2}{4\gamma} \frac{1}{\alpha C_B R}, \quad (32)$$

which satisfies the condition $|\Delta_0| \ll |\Delta_{n>0}|$. Therefore, our N -qubit Toffoli-like gate can be achieved by applying a driving pulse of duration $t_{\text{Toff}} = \pi/|\Delta_{n>0}|$. Up to an overall phase, such a gate flips the phase of the qubit state with all qubits in $|0\rangle$, but has no effect on the other qubit states. However, the particular detunings δ and $\Delta_{e,2}$ also yield

$$\Gamma \equiv \Gamma_0 = \Gamma_1 = \frac{\tilde{\Omega}^2}{2\gamma} \frac{1}{\alpha C_B}, \quad (33)$$

$$\Gamma_{n>1} = \frac{\tilde{\Omega}^2}{4\gamma} \frac{1}{\alpha C_B} \left(2 - \frac{1-1/n}{C_B R^2} \right), \quad (34)$$

again in the limit of $\min\{G, C_B\} \gg 1$. According to Eq. (27), the decay factor $\exp[-(\Gamma_n + \Gamma_{n'})/2]$ cannot be completely removed conditional on the quartit being measured in $|g_1\rangle$, thus, leading to gate errors.

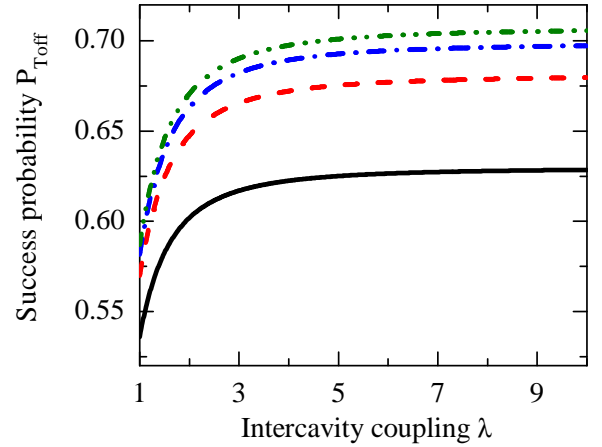
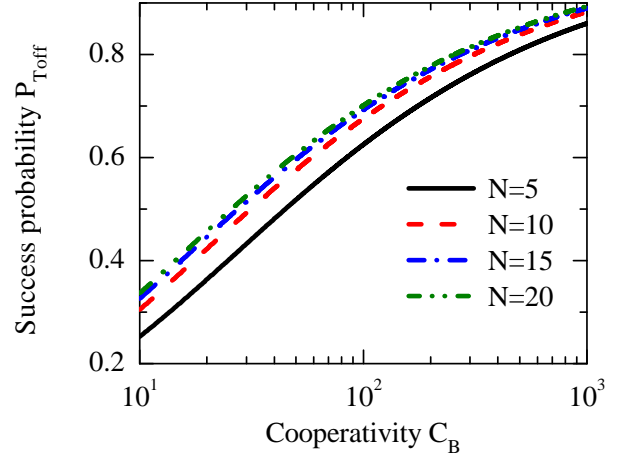


FIG. 3. (Color online) Success probability, P_{Toff} , of our multi-qubit Toffoli-like gate as a function of (a) the cooperativity C_B or (b) the inter-resonator-coupling strength λ , and for $N = 5$ (solid black curves), 10 (dashed red curves), 15 (dashed-dotted blue curves), and 20 (dashed-double dotted olive curves) three-level atoms. Here, we have assumed that $\lambda = 5$ in (a) and $C_B = 100$ in (b). In both panels, the damping rates are set as $\gamma_{g,1} = \gamma_{g,2} = \gamma$, $\gamma_0 = \gamma_1 = 0.5\gamma$, $g_A = g_B$, $C_B = g_B^2/(\kappa\gamma)$, $\lambda = J/(\kappa\sqrt{C_B})$, and $\alpha = \beta = 1$.

In order to quantify the quality of this gate, we need to define a conditional fidelity as

$$F_{\text{Toff}} = \langle \phi | \rho_{\text{qubit}}(t_{\text{Toff}}) | \phi \rangle, \quad (35)$$

where $|\phi\rangle$ is assumed to be the desired state after the gate operation. Correspondingly, the gate error is characterized by $1 - F_{\text{Toff}}$. Considering a generic initial state of the quartit and qutrit atoms,

$$|\Phi\rangle_{\text{ini}} = |g_1\rangle \left[\frac{1}{2^{N/2}} \prod_{k=1}^N (|0\rangle + |1\rangle)_k \right], \quad (36)$$

the success probability and the conditional fidelity is

given by

$$P_{\text{Toff}} = \frac{1}{2^N} \sum_{n=0}^N C_N^n \exp(-\Gamma_n t_{\text{Toff}}), \quad (37)$$

$$F_{\text{Toff}} = \frac{1}{2^{2N} P_{\text{Toff}}} \left[\sum_{n=0}^N C_N^n \exp(-\Gamma_n t_{\text{Toff}}/2) \right]^2, \quad (38)$$

respectively, with $C_N^n = N!/[n!(N-n)!]$ being the binomial coefficient. Again in the limit $\{G, C_B\} \gg 1$, we have $\Gamma_n t_{\text{T}} \ll 0$, which in turn results in

$$P_{\text{Toff}} = 1 - T_p \frac{\pi}{\sqrt{C_B}}, \quad (39)$$

$$F_{\text{Toff}} = 1 - T_f \frac{\pi^2}{C_B}, \quad (40)$$

where the scaling factors T_p and T_f are written as

$$T_p = 2r + \frac{1}{r} \left[\frac{1}{2^N} (1 + S_1) - 1 \right], \quad (41)$$

$$T_f = \frac{1}{2^{N+2} r^2} \left[(1 + S_2) - \frac{1}{2^N} (1 + S_1)^2 \right], \quad (42)$$

respectively. Here, $r = \sqrt{(1/\lambda^2 + \beta/\alpha + 1)/2}$ with $\lambda = G/\sqrt{C_B}$, and $S_x = \sum_{n=1}^N C_N^n/n^x$ for $x = 1, 2$. Together with a success probability

$$P_{\text{Toff}} \propto 1 - 1/\sqrt{C_B}. \quad (43)$$

Thus, the proposed protocol for our multi-qubit Toffoli-like gate has an error scaling as

$$1 - F_{\text{Toff}} \propto 1/C_B, \quad (44)$$

which is a quadratic improvement as compared to gate protocols making use of coherent unitary dynamics in cavities [36], as explained in Ref. [24]. The latter gate suffers errors from the resonator decay and atomic spontaneous emission. Instead our protocol exploits the combined effect of the unitary and dissipative processes, thus, resulting in the above improvement. In fact, the atom-resonator cooperativity C_B could experimentally reach $> 10^4$ in superconducting circuits [3, 54], thus, making the gate error very close to zero and the success probability close to unity.

The success probability is plotted as a function of either the cooperativity C_B or the inter-resonator coupling J , illustrated in Fig. 3. There we have assumed that $\gamma_{g,1} = \gamma_{g,2}$, $\gamma_0 = \gamma_1 = 0.5\gamma$, and the quartit atom is the same as the qutrit atoms, such that $\alpha = \beta = 1$. Similarly, we also plot the corresponding gate error in Fig. 4. As expected, we find that increasing the cooperativity not only makes the success probability very high [see Fig. 3(a)], but it also makes the gate error very low [see Fig. 4(a)]. For example, the success probability of up to ~ 0.9 and the gate error of up to $\sim 2.0 \times 10^{-6}$ can be

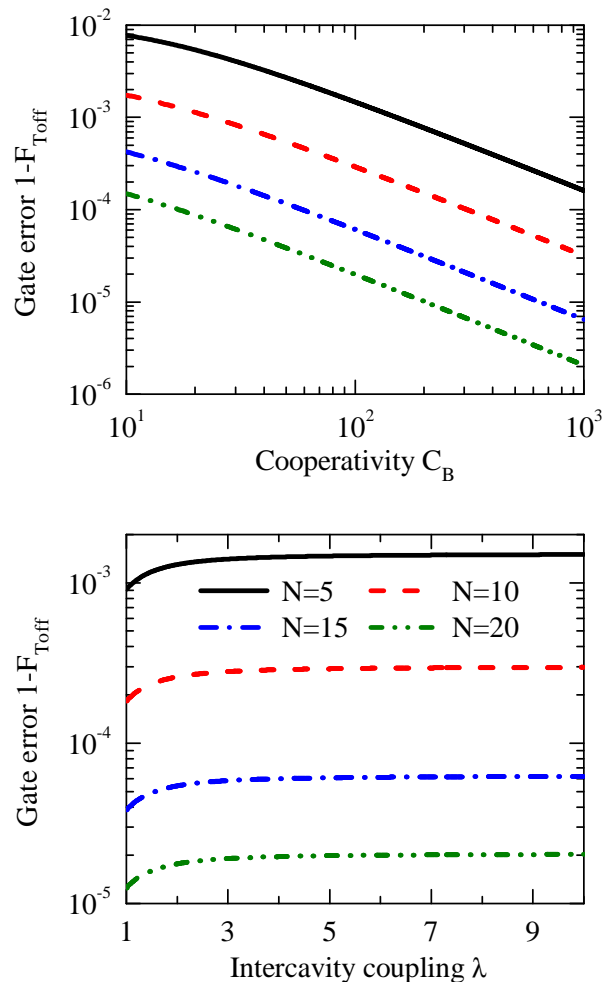


FIG. 4. (Color online) Analytical results for the gate error (gate infidelity), $1 - F_{\text{Toff}}$, of our multi-qubit Toffoli-like gate as a function of (a) the cooperativity C_B or (b) the inter-resonator coupling strength λ , and for $N = 5$ (solid black curves), 10 (dashed red curves), 15 (dashed-dotted blue curves), and 20 (dashed-double dotted olive curves) three-level atoms (qutrits). The other parameters are also set to be the same as in Fig. 3. There the success probability is in a one-to-one correspondence with the gate error here.

achieved when $N = 20$, $\lambda = 5$ and $C_B = 10^3$. Within the limit $\lambda \gg 1$, the $|g_1\rangle \leftrightarrow |e_1\rangle \leftrightarrow |e_2\rangle \leftrightarrow |g_2\rangle$ three-photon Raman transition off-resonantly mediated by means of the symmetric mode a_+ could be neglected, yielding $r \rightarrow \sqrt{(\beta/\alpha + 1)/2}$; hence, according to Eqs. (39) and (40), the gate error is limited by an upper bound [see Fig. 4(b)], along with the corresponding success probability also upper bounded [see Fig. 3(b)].

V. HERALDED CONTROLLED-Z GATE

Let us now consider the heralded near-deterministic realization of the two-qubit controlled-Z (CZ) gate. This gate is also known as the two-qubit controlled-sign gate, controlled-phase-flip gate, or controlled-phase gate. Specifically, the action of the CZ gate in our system can be explained as follows: Conditioned on the detection of the quartit in its ground state $|g_1\rangle$, the CZ gate flips the phase of the state $|11\rangle$ of an arbitrary two-qubit pure state $|\psi\rangle = c_0|00\rangle + c_1|01\rangle + c_1|10\rangle + c_1|11\rangle$ or any mixture of such states, where c_k are the complex normalized amplitudes and the qubit states correspond to the two lowest-energy levels of the two qutrits in the resonator B .

It follows from Eq. (28) that, in order to completely remove the gate error, the decay rate Γ_n needs to be independent of the qutrits. To this end, we retune the detunings δ and $\Delta_{e,2}$ to be

$$\frac{\delta}{\gamma} = \frac{1}{2(2D + G^{-1})}, \quad (45)$$

$$\frac{\Delta_{e,2}}{\gamma} = \alpha C_B (D + G^{-1}), \quad (46)$$

where $D = \sqrt{[1/G^2 + \beta/(\alpha C_B)]}/2$. Thus, in the limit $\{G, C_B\} \gg 1$, retuning δ and $\Delta_{e,2}$ as above yields an n -independent decay rate

$$\Gamma_n = \Gamma, \quad (47)$$

for all n , while Eq. (33) is valid only for $n = 0, 1$. The corresponding energy shifts are given by

$$\Delta_0 = -\frac{\Gamma D}{2}, \quad (48)$$

and

$$\Delta_{n>0} = -\frac{\tilde{\Omega}^2}{2\gamma} \frac{n(2D + 1/G)}{\alpha C_B (4nD^2 + 2nD/G + 1/C_B)}. \quad (49)$$

Subsequently, according to Eq. (28), the conditional density operator of the qutrits becomes

$$\rho_{\text{qutrit}}(t) = \sum_{n,n'=0}^N e^{-i(\Delta_n - \Delta_{n'})t} \mathcal{P}_n [\langle g_1 | \rho(0) | g_1 \rangle] \mathcal{P}_{n'}, \quad (50)$$

where the decay factor, $\exp(-\Gamma t)$, has been eliminated through a measurement conditional on the quartit atom being detected in the state $|g_1\rangle$. Together with single-qubit operations, we can utilize the dynamics described by the $\rho_{\text{qutrit}}(t)$ of Eq. (50) to implement a heralded CZ gate for $N = 2$.

For this purpose, the duration of the driving pulse is chosen to be

$$t_{\text{CZ}} = \pi |\Delta_2 - 2\Delta_1 + \Delta_0|^{-1}, \quad (51)$$

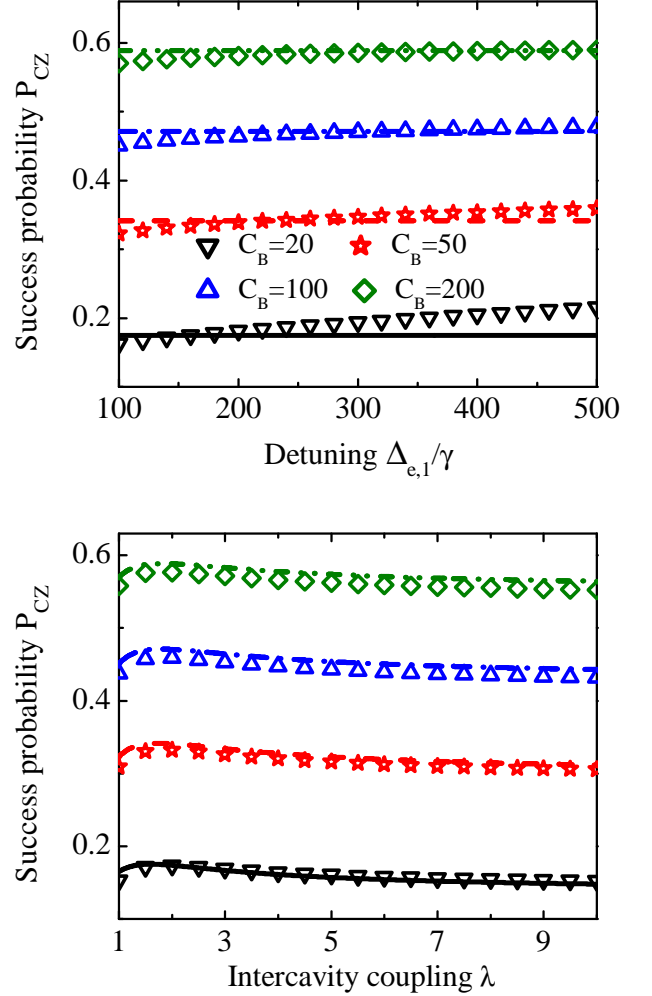


FIG. 5. (Color online) Numerical simulations (marked by symbols) for the success probability, P_{CZ} , of the heralded CZ gate, and for the cooperativity $C_B = 20$ (black down-triangles), 50 (red stars), 100 (blue up-triangles), and 200 (olive diamonds). The success probability is displayed versus the detuning $\Delta_{e,1}$ given in terms of the overall decay rate γ in (a), as well as versus the inter-resonator coupling strength λ in (b). For comparison, we also plot the analytical success probability (curves), and find that the analytical results are in good agreement with the exact numerics. Here, we have assumed $\lambda = 1.84$ in (a) and $\Delta_{e,1} = 150\gamma$ in (b). In both panels we have set that $\gamma_{g,1} = \gamma_{g,2} = \gamma$, $\gamma_0 = \gamma_1 = 0.5\gamma$, $\kappa = 10\gamma$, $g_A = g_B$, $C_B = g_B^2/(\kappa\gamma)$, $\lambda = J/(\kappa\sqrt{C_B})$, $\alpha = \beta = 1$, $\Omega_1 = \Delta_{e,1}/(10C_B^{1/4})$, and $\Omega_2 = 4\gamma C_B^{1/4}$.

and the unitary operation on each qubit, applied after the pulse, has the following action

$$\mathcal{U}|0\rangle = e^{i\Delta_0 t_{\text{CZ}}/2}|0\rangle, \quad \mathcal{U}|1\rangle = e^{i(2\Delta_1 - \Delta_0)t_{\text{CZ}}/2}|1\rangle, \quad (52)$$

so as to ensure the right phase evolution. The resulting gate is either to flip the phase of the qubit state $|11\rangle$, or to leave the otherwise qubit states unchanged. The

associated success probability for any initial pure state is

$$P_{CZ} = \exp(-\Gamma t_{CZ}), \quad (53)$$

which can, as long as $\{G, C_B\} \gg 1$, be approximated as

$$P_{CZ} = 1 - Z_p \frac{\pi}{\sqrt{C_B}}, \quad (54)$$

with a scaling factor

$$Z_p = 2d + \frac{3}{2(2d+1/\lambda)} + \frac{1}{4d(2d+1/\lambda)^2}, \quad (55)$$

where $d = \sqrt{(1/\lambda^2 + \beta/\alpha)}/2$. If assuming that the desired state is $|\psi\rangle$, we can calculate the conditional fidelity for this gate via

$$F_{CZ} = \langle \psi | (\mathcal{U} \otimes \mathcal{U}) \rho_{\text{qudit}}(t_{CZ}) (\mathcal{U} \otimes \mathcal{U})^\dagger | \psi \rangle, \quad (56)$$

and then can directly find $F_{CZ} = 1$. This implies that with the single-qubit operation, we achieve a more significant improvement than that shown in Eq. (44), and, thus, by decreasing Ω_1 and increasing $\Delta_{e,1}$, an arbitrarily small gate error can even be achievable.

In order to confirm the heralded CZ gate, we now perform numerical simulations exactly using, instead of the effective master equation in Eq. (18), the full zero-temperature master equation, given by Eq. (13), for the density operator $\rho_T(t)$ of the total system initially in

$$|\Psi\rangle_{\text{ini}} = |\Phi\rangle_{\text{ini}} \otimes |\text{vac}\rangle, \quad (57)$$

with $|\text{vac}\rangle$ being the vacuum state of the coupled resonators and $|\Phi\rangle_{\text{ini}}$ is given in Eq. (36). The numerical simulations calculate the success probability P_{CZ} , the conditional density operator $\rho_{\text{qudit}}(t_{CZ})$ and fidelity F_{CZ} using the expressions below:

$$P_{CZ} = \sum_{n=0,1,2} \text{Tr}[(|g_1\rangle\langle g_1| \otimes \mathcal{P}_n \otimes \mathcal{I}) \rho_T(t_{CZ})], \quad (58)$$

$$\rho_{\text{qudit}}(t_{CZ}) = \frac{1}{P_{CZ}} \text{Tr}_{\text{cav}}[|g_1\rangle\langle g_1| \rho_T(t_{CZ})], \quad (59)$$

$$F_{CZ} = \langle \psi | (\mathcal{U} \otimes \mathcal{U}) \rho_{\text{qudit}}(t_{CZ}) (\mathcal{U} \otimes \mathcal{U})^\dagger | \psi \rangle, \quad (60)$$

where Tr and Tr_{cav} are trace operations over the total system and the resonators, respectively, and \mathcal{I} is an identity operator related to the two resonators. For the simulations, we assume that $\gamma_{g,1} = \gamma_{g,2}$, $\gamma_0 = \gamma_1 = 0.5\gamma$, and $\alpha = \beta = 1$; moreover, we take $\kappa = 10\gamma$, $\Omega_1 = \Delta_{e,1}/(10C_B^{1/4})$, $\Omega_2 = 4\gamma C_B^{1/4}$, and determine the detunings δ and $\Delta_{e,2}$ according to Eqs. (45) and (46), respectively. Then, we calculate the success probability P_{CZ} and the gate error $1 - F_{CZ}$, as a function of either the detuning $\Delta_{e,1}$ or the inter-resonator coupling J , for different cooperativities [55, 56]. The numerical results

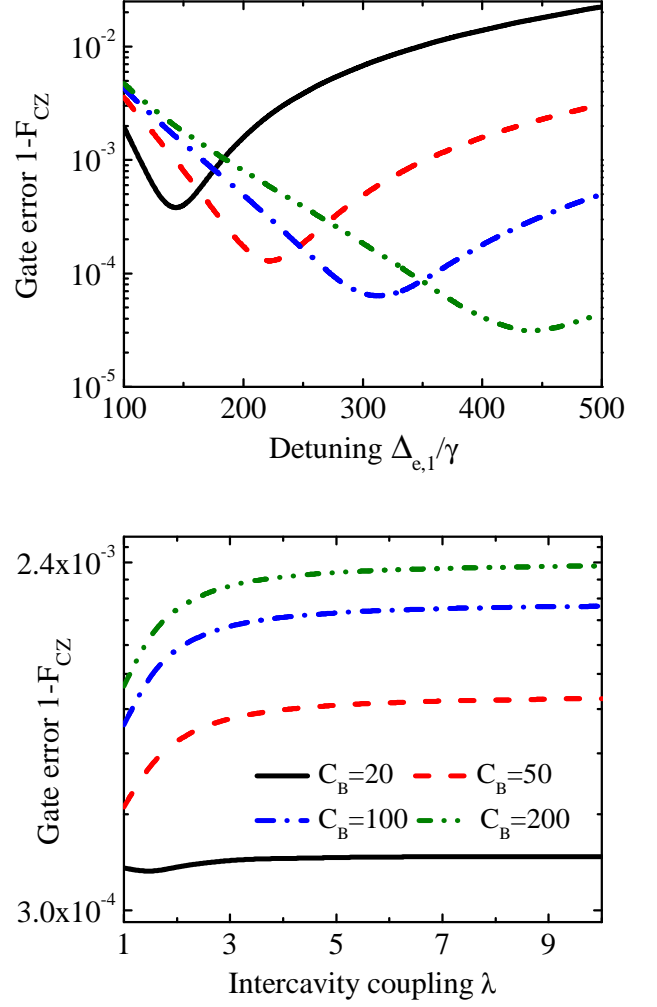


FIG. 6. (Color online) Numerical simulations for the gate error, $1 - F_{CZ}$, of the heralded CZ gate, and for the cooperativity $C_B = 20, 50, 100$, and 200 . Upon setting the same parameters as used in Fig. 5, the gate error is displayed versus the detuning $\Delta_{e,1}/\gamma$ in (a), as well as versus the inter-resonator coupling strength λ in (b). There is a one-to-one correspondence between the gate error here and the numerically calculated success probability P_{CZ} in Fig. 5.

(marked by symbols) are plotted in Fig. 5 for the success probability and in Fig. 6 for the gate error. The analytical success probability, which we use curves to represent, is also plotted in Fig. 5, and shows good agreement with the exact results. In Fig. 6(a) we find that as the detuning $\Delta_{e,1}$ increases, the gate error first decreases and then increases. This is because, in addition to suppressing the error from the $|e_1\rangle \rightarrow |g_1\rangle$ decay (see the Appendix), such an increase in $\Delta_{e,1}$, however, increases the driving strength $\Omega_1 = \Delta_{e,1}/(10C_B^{1/4})$, and hence the non-adiabatic error. There is a trade-off between the two errors. Again within the limit $\lambda \gg 1$, the off-resonant Raman transition could be removed as before, yielding

$d \rightarrow \sqrt{\beta/(2\alpha)}$. As a result, one finds that the scaling factor, given by Eq. (55), is equal to

$$Z_p \rightarrow 2d + \frac{3}{4d} + \frac{1}{16d^3}. \quad (61)$$

Hence, when λ is sufficiently large, both the success probability and the gate error will be independent of λ , illustrated in Figs. 5(b) and 6(b). It should be noted that to calculate the success probability in Fig. 5(a), and the gate error in Fig. 6(a), we have chosen $\lambda = 1.84$, because this value can lead to the shortest driving pulse for the two-qubit gate if $\alpha = \beta$.

Finally, we note that the energy shifts, Δ_n , for both Toffoli-like and CZ gates involve only the ground states, because the effective master equation (18) is obtained by adiabatically eliminating the excited states. According to the effective Hamiltonian of Eq. (19), a quantum state, with n , of the qutrit atoms has an energy shift, Δ_n . Consequently, a dynamical phase, $\exp(-i\Delta_n t)$, can be imposed by choosing the appropriate evolution times for a given quantum state. Thus, the quantum gates, including the multi-qubit Toffoli-like (see Table II) and two-qubit CZ gates (Table III), are realized by choosing appropriate evolution times.

VI. CONCLUSIONS

We have described a method for performing a heralded near-deterministic controlled phase gates in two distant resonators in the presence of decoherence, including the two-qubit controlled-Z (CZ) gate and its multi-qubit-controlled generalization known, which can be referred to as a Toffoli-like gate. Our proposal is a generalization of the single-resonator method introduced by Borregaard *et al.* [24]. The method in our paper uses an auxiliary microwave-driven four-level atom (qutrit) inside one resonator to serve as both an intra-resonator photon source and a detector, and, thus, to control and also herald the gates on atomic qutrits inside the other resonator. In addition to the quantum gate fidelity scalings, which are superior to traditional cavity-assisted deterministic gates, this method demonstrates a macroscopically-distant herald for controlled quantum gates, and at the same time avoids the difficulty in individually addressing a microwave-driven atom. We note that the original method of Ref. [24] is based on the heralding and operational atoms coupled to the same resonator mode acting as a cavity bus. Here, we are not applying this cavity bus.

The operator formalism used in our paper to calculate the effective Hamiltonian [i.e., Eq. (15)] and the effective Lindblad operators [i.e., Eq. (16)] follows the approach in Ref. [12]. The dynamical behavior of the system can be fairly-well approximated by the effective master equation. For example, the results obtained respectively by using the full and effective master equations in Figs. 4 and 6 in Ref. [12], or in Figs. 5 and 6 in our paper, are in

good agreement. The physics of this approximation can be understood by the simple qutrit atom in Fig. 2. A direct picture is that the three-level atom in state $|0\rangle$ is first excited to the state $|2\rangle$, then can go to the state $|1\rangle$ by an atomic decay, or can go back to the state $|0\rangle$ also by another atomic decay. If starting in the state $|1\rangle$, the atom has a similar behavior. Roughly speaking, there exists an indirect interaction between the states $|0\rangle$ and $|1\rangle$, as well as the direct interaction between the states $|0\rangle$, $|1\rangle$, and the state $|2\rangle$, mediated by the atomic decays. Thus, after adiabatically eliminating the state $|2\rangle$, the energy shifts of the states $|0\rangle$, $|1\rangle$, and even a direct interaction between them, would be induced by the atomic decays. Upon combining the above processes, mediated by the atom decays with the coherent drives and other interactions, the effective master equation was obtained here in analogy to that in Ref. [12].

We have also proposed a circuit-QED system with superconducting qutrits and quartits implementing the proposed protocol as shown in Fig. 1. Our circuit includes two coupled transmission-line resonators, which are linearly coupled via a SQUID. These resonators can be coupled to qutrits via a capacitance. For example, we assumed a phase quartit coupled to one of the resonators serving a herald, and N identical flux qutrits coupled to the other resonator for performing the controlled phase gates. Typically, the length of a transmission-line resonator can be of up to the order of cm [57], but the size of a superconducting atoms is of $\sim \mu\text{m}$ [3].

We assume realistic parameters from experiments with superconducting quantum circuits [54, 57]. Specifically: $\gamma/2\pi = 10$ MHz, $\kappa/2\pi = 6$ MHz, $C_A = C_B = 170$ ($\alpha = \beta = 1$), $\Delta_{e,1}/2\pi = 420$ MHz, $\Omega_1/2\pi = 70$ MHz, $\lambda = 1.84$, $J/2\pi = 144$ MHz, and $\Omega_2/2\pi = 87$ MHz. The implementation of the CZ gate with these parameters can result in a success probability of ~ 0.55 , a gate error of ~ 0.006 , and a gate time of $\sim 6 \mu\text{s}$. These gate performance parameters are quite good for a coupled-resonator system. Indeed by increasing the capacitance of the capacitor between the TLRs, it is possible to achieve a stronger J [39] and, thus, a smaller gate error and a shorter gate time. Furthermore, the decoherence time T_2 of a flux qubit can be improved to $\sim 85 \mu\text{s}$ [58], which is much larger than $6 \mu\text{s}$. This justifies neglecting such decoherence effect on the flux qutrits. But for a phase qubit, T_2 reaches $\sim 1 \mu\text{s}$, which is smaller than $6 \mu\text{s}$ [41]. Nevertheless, the phase quartit in our protocol works only as an auxiliary atom to herald the quantum gate by measuring the $|g_1\rangle$ population on which the decoherence, quantified T_2 , has no effect. Hence, such decoherence of both phase quartit and flux qutrits can be effectively neglected in our protocol.

Another possible implementation can be based on ultracold atoms coupled to nanoscale optical cavities. In this situation, the atoms ^{87}Rb can be used for both quartit and qutrit atoms [24, 59, 60]. Furthermore, nanoscale optical cavities have been realized by the use of defects in a two dimensional photonic crystal, and one can place

such nanocavities very close to each other to directly couple them by evanescent fields [61], or one can use a common waveguide (a quantum bus) to indirectly couple such cavities [62].

Although we have chosen to discuss the specific case of two coupled resonators, this description may be extended to a coupled-cavity array [63–67]. Hence, it would enable applications such as scalable quantum information processing and long-distance quantum communication.

VII. ACKNOWLEDGMENTS

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Appendix: Derivation of effective master equation

In this Appendix, we will derive the effective master equation by using the method in Ref. [12]. We start with the total Hamiltonian H_T in the main text. Upon introducing the symmetric and antisymmetric modes, $a_{\pm} = (a_A \pm a_B)/\sqrt{2}$, of the coupled resonators, and then switching into a rotating frame with respect to

$$H_{\text{rot}} = \sum_{k=1}^N \left[(\omega_c + \omega_1 - J) |2\rangle_k \langle 2| + \sum_{z=0,1} \omega_z |z\rangle_k \langle z| \right] + \bar{\omega} |e_1\rangle \langle e_1| + (\bar{\omega} + \omega_{m,2}) |e_2\rangle \langle e_2| + \sum_{x=1,2} \omega_{g,x} |g_x\rangle \langle g_x| + (\omega_c - J) (a_+^\dagger a_+ + a_-^\dagger a_-), \quad (\text{A.1})$$

where $\bar{\omega} = \omega_{m,1} + \omega_{g,1}$, the total Hamiltonian is transformed to

$$H_T = H_e + V + V^\dagger, \quad (\text{A.2})$$

as given in the main text. With the Lindblad operators in Eq. (14), we obtain the no-jump Hamiltonian of the

form

$$H_{\text{NH}} = \sum_{k=1}^N \left\{ \bar{\delta} |2\rangle_k \langle 2| + \frac{g_B}{\sqrt{2}} [(a_+ - a_-) |2\rangle_k \langle 1| + \text{H.c.}] \right\} + \bar{\Delta}_{e,1} |e_1\rangle \langle e_1| + \bar{\Delta}_{e,2} |e_2\rangle \langle e_2| + \bar{J} a_+^\dagger a_+ - \frac{i\kappa}{2} a_-^\dagger a_- + \frac{g_A}{\sqrt{2}} [(a_+ + a_-) |e_2\rangle \langle g_2| + \text{H.c.}] + \frac{\Omega_2}{2} (|e_2\rangle \langle e_1| + \text{H.c.}), \quad (\text{A.3})$$

where $\bar{\Delta}_{e,1} = \Delta_{e,1} - i\gamma_{g,1}/2$, $\bar{\Delta}_{e,2} = \Delta_{e,2} - i\gamma_{g,2}/2$, $\bar{\delta} = \delta - i\gamma/2$, and $\bar{J} = 2J - i\kappa/2$. Following the formalism in Ref. [12], the effective Hamiltonian and Lindblad operators are given, respectively, by

$$H_{\text{eff}} = \frac{1}{2} V^\dagger \left[H_{\text{NH}}^{-1} + (H_{\text{NH}}^{-1})^\dagger \right] V, \quad (\text{A.4})$$

$$L_{\text{eff}}^j = L_j H_{\text{NH}}^{-1} V. \quad (\text{A.5})$$

To proceed, we work within the single-excitation subspace and, after a straightforward calculation, obtain

$$\Delta_n = -\frac{\Omega_1}{2\sqrt{\gamma_{g,1}}} \text{Re}(r_{g_1,n}), \quad (\text{A.6})$$

$$r_{+,n} = \frac{\Omega_1 \tilde{\Omega}_m \sqrt{C_A}}{4\tilde{J} \mathcal{X}_n \sqrt{2\gamma}} (i\tilde{\delta} + 2nC_B), \quad (\text{A.7})$$

$$r_{-,n} = -\frac{\Omega_1 \tilde{\Omega}_m \sqrt{C_A}}{2\mathcal{X}_n \sqrt{2\gamma}} (\tilde{\delta} - nC_B/\tilde{J}), \quad (\text{A.8})$$

$$r_{g_1,n} = \frac{\Omega_1 \sqrt{\gamma_{g,1}}}{2\gamma \mathcal{X}_n} \left[i\tilde{\delta} \tilde{\Delta}_{e,2} + (C_A \tilde{\delta} + nC_B \tilde{\Delta}_{e,2}) \times \left(1 - i/(2\tilde{J}) \right) - 2nC_A C_B / \tilde{J} \right], \quad (\text{A.9})$$

$$r_{g_2,n} = -\frac{\Omega_1 \tilde{\Omega}_m \sqrt{\gamma_{g,2}}}{4\gamma \mathcal{X}_n} \left\{ i\tilde{\delta} + nC_B \left[1 - i/(2\tilde{J}) \right] \right\}, \quad (\text{A.10})$$

$$r_{k,n} = -\frac{\Omega_1 \tilde{\Omega}_m \sqrt{\gamma_k C_A C_B}}{4\gamma \mathcal{X}_n} \left[1 + i/(2\tilde{J}) \right], \quad (\text{A.11})$$

where $C_A = g_A^2/(\kappa\gamma)$, $\tilde{\Delta}_{e,1} = \Delta_{e,1}/\gamma - i\gamma_{g,1}/(2\gamma)$, $\tilde{J} = 2J/\kappa - i/2$, $\tilde{\Omega}_2 = \Omega_2/\gamma$, $Z = \tilde{\Delta}_{e,1} \tilde{\Delta}_{e,2} - (\tilde{\Omega}_m/2)^2$, $\mathcal{X}_n = iZ\tilde{\delta} + (C_A \tilde{\delta} \tilde{\Delta}_{e,2} + nC_B Z) \left[1 - i/(2\tilde{J}) \right] - 2nC_A C_B \tilde{\Delta}_{e,1}/\tilde{J}$, and $k = 0, 1$. In the limit $\Omega_2 \ll \Delta_{e,1}$, we can make a Taylor expansion around $\Omega_2/\Delta_{e,1} = 0$, yielding

$$\Delta_n = -\frac{\Omega_1^2}{4\Delta_{e,1}} - \frac{\tilde{\Omega}_2^2}{4\gamma} \text{Re} \left\{ \frac{1}{\mathcal{Y}_n} \left\{ i\tilde{\delta} + nC_B \left[1 - \frac{i}{2\tilde{J}} \right] \right\} \right\}, \quad (\text{A.12})$$

$$r_{+,n} = \frac{\tilde{\Omega}\sqrt{C_A}}{2\tilde{\mathcal{J}}\mathcal{Y}_n\sqrt{2\gamma}} \left(i\tilde{\delta} + 2nC_B \right), \quad (\text{A.13})$$

$$r_{g_2,n} = -\frac{\tilde{\Omega}\sqrt{\gamma_{g,2}}}{2\gamma\mathcal{Y}_n} \left\{ i\tilde{\delta} + nC_B \left[1 - i/(2\tilde{\mathcal{J}}) \right] \right\}, \quad (\text{A.16})$$

$$r_{-,n} = -\frac{\tilde{\Omega}\sqrt{C_A}}{\mathcal{Y}_n\sqrt{2\gamma}} \left(i\tilde{\delta} - nC_B/\tilde{\mathcal{J}} \right), \quad (\text{A.14})$$

$$r_{k,n} = -\frac{\tilde{\Omega}\sqrt{\gamma_k C_A C_B}}{2\gamma\mathcal{Y}_n} \left[1 + i/(2\tilde{\mathcal{J}}) \right], \quad (\text{A.17})$$

$$r_{g_1,n} = \frac{\Omega_1\sqrt{\gamma_{g,1}}}{2\Delta_{e,1}} + \frac{\tilde{\Omega}\sqrt{\tilde{\gamma}_{g,1}}}{2\gamma\mathcal{Y}_n} \left\{ i\tilde{\delta} + nC_B \left[1 - \frac{i}{2\tilde{\mathcal{J}}} \right] \right\}, \quad (\text{A.15})$$

where $\tilde{\Omega} = \Omega_1\Omega_2/(2\Delta_{e,1})$, $\tilde{\gamma}_{g,1} = \gamma_{g,1}\Omega_2^2/(4\Delta_{e,1}^2)$, and $\mathcal{Y}_n = i\tilde{\delta}\tilde{\Delta}_{e,2} + (C_A\tilde{\delta} + nC_B\tilde{\Delta}_{e,2}) \left[1 - i/(2\tilde{\mathcal{J}}) \right] - 2nC_A C_B/\tilde{\mathcal{J}}$. In Eq. (A.12), the term, $-\Omega_1^2/(4\Delta_{e,1})$, of Δ_n can be removed because it has no effect on the phase gates. Meanwhile, in Eq. (A.15), we find that the $|e_1\rangle \rightarrow |g_1\rangle$ decay is suppressed by increasing $\Delta_{e,1}$, such that the second term of $r_{g_1,n}$ can be removed as long as $\tilde{\gamma}_{g,1} \ll 1$. Thus, Eqs. (A.12)–(A.17), under the assumption that $\kappa \ll J$, reduce to Eqs. (20) and (26) in the main text.

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