Directional quantum state transfer in a dissipative Rydberg-atom-cavity system

D. X. Li and X. Q. Shao*

Center for Quantum Sciences and School of Physics, Northeast Normal University, Changchun 130024, China and Center for Advanced Optoelectronic Functional Materials Research, and Key Laboratory for UV Light-Emitting Materials and Technology of Ministry of Education, Northeast Normal University, Changchun 130024, China

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Quantum state transfer is a basic task for quantum information processing, which can be easily executed using the coherent dynamics for a closed system instead of the dissipative dynamics for an open system. Here we propose a dissipation-assisted scheme to directionally transfer an arbitrary quantum state from the sender A to the receiver B by virtue of the Rydberg antiblockade mechanism, the laser-induced Raman transition, and the photon loss of an optical cavity. The prominent advantage of the current proposal is that it does not require accurate control over the relevant parameters of the system, as the target state is the steady state of the whole process. The effect of atomic spontaneous emission for the excited states is dramatically restricted by the adiabatic elimination, and a high population of the transferred state around 99% is achievable with the current experimental technology.

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I. INTRODUCTION

Quantum state transfer (QST) is one of the pivotal building blocks for quantum computation and quantum communication [1], and it entails the coherent transmission of an arbitrary quantum state from the sender A to the receiver B with a high fidelity. Up to the present, much literature has been proposed to implement undamaged and robust QST in many physical systems [2–10]. However, the traditional state transfer schemes governed by the unitary dynamics are sensitive to the small variations of the Hamiltonian. In order to stabilize a quantum state, which periodically oscillates between different systems, into the target site, one has to impose severe operations on the transport time such as turning on and turning off the interaction. Besides these defects, the known unitary-dynamics methods have to resist against the quantum dissipation, which will account for the undesired decoherent effect and destroy the transferred information.

It is well recognized that quantum dissipation was considered a primary obstacle for quantum information processing [11] until notable techniques regarding dissipation as an important resource [12] were proposed. Now considerable dissipative applications have appeared in succession [13–29], among which the most general strategy is to make the quantum system evolve into a unique steady state via dissipation dynamics. Nevertheless, it is notoriously difficult to design a proposal for dissipatively transferring an arbitrary quantum state with the standard form of QST, such as $|\psi\rangle_A|0\rangle_B \rightarrow |0\rangle_A|\psi\rangle_B$ ($|\psi\rangle = \alpha|0\rangle + \beta|1\rangle$ is the transferred logical qubit with $|\alpha|^2 + |\beta|^2 = 1$). Recently Wang and Gertler [30] have proven the minimum system construction for dissipative QST and put forward a new type of cascaded system, where the dimension is 3×2 (between one physical qutrit and one physical qubit), plus one auxiliary reservoir, to autonomously transfer a quantum state between qubits without time-dependent controls as

$$|\psi\rangle_A|\mathrm{vac}\rangle_B \rightarrow |\mathrm{vac}\rangle_A|\psi\rangle_B, \tag{1}$$

where $|\mathrm{vac}\rangle$ designates a predefined state void of information, and $|\mathrm{vac}\rangle_A|\psi\rangle_B$ is the steady state of the whole transfer process. They first performed their design in superconducting circuit quantum electrodynamics (QED) using nonlinear couplings between transmon and cavity modes and transferred a logical equator state with a fidelity of up to 93%. However, for the circuit QED scheme, the interaction term engineered by the four-wave mixing Hamiltonian is too weak to be acquired in many physical systems. Therefore, they harnessed another proof-of-principle protocol with only bilinear interaction to prove their theory. However, the logical states have to be encoded by three-atom superposition states, causing a waste of quantum resource. Meanwhile, Matsuzaki *et al.* [31] presented a similar one-way transfer of the quantum states via decoherence of a tailored environment, where they considered two qubits interacting via flip-flop interaction and collectively coupling to a cavity. They also required that the resonance frequencies of the two qubits be different. Obviously, the above scenarios demand rigorous conditions or special systems, which motivates us to find an alternative proposal to dissipatively achieve a QST with a higher fidelity.

The Rydberg atoms have been the subject of intensive studies in the context of quantum information processing [32], since its many exaggerated properties are tempting for controlling matter and electromagnetic fields at the quantum level, such as large geometrical size, long lifetime, large transition dipole moments between neighboring levels, and large polarizability. Particularly, the Rydberg-mediated interaction is a preferred choice for various applications [33–51]. Furthermore, great attention is also paid to combining the fields of Rydberg atoms and cavity QED [52–58]; e.g., Parigi *et al.* [53] discussed and measured dispersive optical nonlinearities in an ensemble of cold Rydberg atoms placed inside a low-finesse

*Corresponding author: shaoxq644@nenu.edu.cn

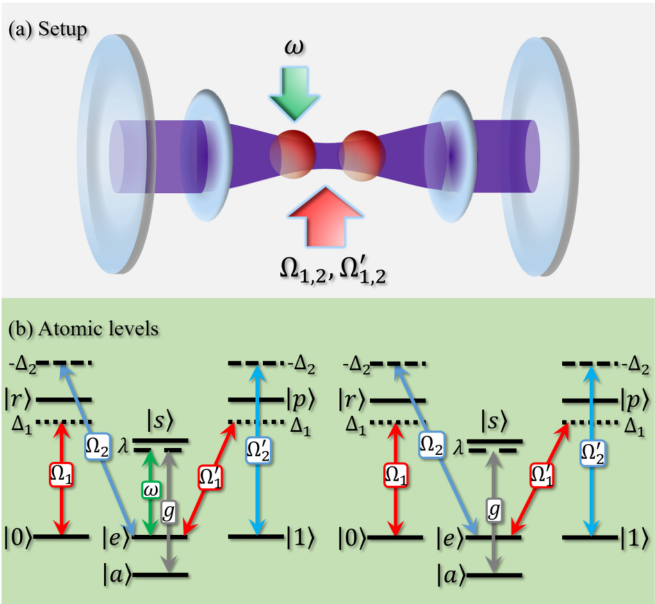

FIG. 1. (a) Setup for the dissipative scheme. Atoms A and B collectively interact with four classical lasers and one optical cavity. (b) Diagram of corresponding atomic energy levels.

optical cavity. Maghrebi *et al.* [57] investigated a proposal to realize fractional quantum Hall states of light, where the quasi-two-dimensional cloud of Rydberg atoms overlapped with an array of cavity modes created by arrays of microlenses or spherical micromirrors.

Motivated by these existing works, we present a reliable dissipation-assisted scheme to accomplish directional QST by the combination of a Rydberg-antiblockade mechanism, a laser-induced Raman transition, and the photon loss of an optical cavity. The corresponding setup and the atomic levels are shown in Figs. 1(a) and 1(b), respectively. We assume that two identical Rydberg atoms, consisting of four ground states ($|0\rangle$, $|1\rangle$, $|e\rangle$, $|a\rangle$), one excited state ($|s\rangle$), and two Rydberg states ($|r\rangle$, $|p\rangle$), are trapped to a common cavity, where the states $|0\rangle$ and $|1\rangle$ are utilized as encoded quantum bits.

The scheme can be divided into two compositions: (i) Rydberg-interaction-induced swap interaction, making use of the ingenious integration of the Rydberg antiblockade mechanism and the laser-induced Raman transition, swaps an arbitrary quantum state between atom A (left one) and atom B (right one) as $|\psi\rangle_A|e\rangle_B \leftrightarrow |e\rangle_A|\psi\rangle_B$, and (ii) photon-loss-induced dissipative dynamics stabilizes the state $|e\rangle_A|\psi\rangle_B$ at the state $|a\rangle_A|\psi\rangle_B$ by virtue of the laser-induced Raman transition and cavity loss. The cooperation of the two compositions turns the state $|a\rangle_A|\psi\rangle_B$ into the unique steady state for the subsystem beginning with $|\psi\rangle_A|e\rangle_B$, which means that an arbitrary quantum state initialized in the sender A will accumulate in the receiver B via the dissipation.

On the whole, our scheme has five features.

(i) The photon loss of an optical cavity is a powerful resource to stabilize an arbitrary quantum state.

(ii) Owing to the adiabatic elimination of the excited states and the Rydberg states, the adverse effect of the atomic spontaneous emission is remarkably depressed.

(iii) The target state is the steady state of the whole process; thus no strictly determinate timescale is needed.

(iv) The population of the target state can be about 99% with the state-of-the-art technology.

(v) The scheme can be generalized to defeat the collective dephasing from the environment.

The structure of the paper is as follows. In Sec. II, we introduce the principle of our directional QST scheme in detail. In Sec. III, we systematically investigate the influences of relevant parameters on our scheme and substantially discuss the feasibility of the scheme with the current experimental parameters. In Sec. IV, we successfully propose a generalized scheme to resist the collective dephasing from the environment. Finally, we sum up the work in Sec. V.

II. PRINCIPLE OF THE DIRECTIONAL QST SCHEME

The setup and the atomic energy levels of the total system are illustrated in Figs. 1(a) and 1(b), respectively. Atoms A and B collectively interact with four classical lasers and one optical cavity. Their transitions, $|0\rangle \leftrightarrow |r\rangle$ and $|e\rangle \leftrightarrow |r\rangle$, are driven by two lasers in possession of Rabi frequencies Ω_1 and Ω_2 and detunings Δ_1 and $-\Delta_2$, while other two lasers (Rabi frequencies Ω'_1 and Ω'_2 and detunings Δ_1 and $-\Delta_2$) complete the transitions $|e\rangle \leftrightarrow |p\rangle$ and $|1\rangle \leftrightarrow |p\rangle$, respectively. In the meantime, the excited states $|s\rangle$ are both coupled with the quantized cavity mode (strength g, detuning λ) to fulfill the transition $|s\rangle \leftrightarrow |a\rangle$. Besides the above operations, atom A is also manipulated with an extra laser (Rabi frequency ω, detuning λ) to realize the transition $|s\rangle \leftrightarrow |e\rangle$. Moreover, since the initial state of the second atom is set in $|e\rangle$, the interaction between atom B and the optical cavity becomes useless during the entire process, which is ignored in the following discussion.

In the interaction picture, the total Hamiltonian reads ($\hbar = 1$)

$$H = H_X + H_D, \tag{2}$$

where

$$\begin{aligned} H_X = & \sum_{j=A,B} (\Omega_1 |r\rangle_j\langle 0| + \Omega'_1 |p\rangle_j\langle e|)e^{i\Delta_1 t} + (\Omega_2 |r\rangle_j\langle e| \\ & + \Omega'_2 |p\rangle_j\langle 1|)e^{-i\Delta_2 t} + \text{H.c.} + U_{rr}|rr\rangle\langle rr| \\ & + U_{pp}|pp\rangle\langle pp| + U_{rp}|rp\rangle\langle rp| + U_{pr}|pr\rangle\langle pr|, \\ H_D = & (\omega|s\rangle_A\langle e| + g|s\rangle_A\langle a|c)e^{i\lambda t} + \text{H.c.} \end{aligned}$$

Here c stands for the annihilation operator of the optical cavity. $U_{\alpha\beta}$ bridges the Rydberg-mediated interaction between atom A in $|\alpha\rangle$ and atom B in $|\beta\rangle$, which can be induced by the dipole-dipole potential of the scale C_3/r^3 or the long-range van der Waals interaction proportional to C_6/r^6 (r is the distance of two atoms and C_3 and C_6 depend on the quantum numbers of the Rydberg state) [32,59]. The relation of $U = U_{rr} = U_{pp} \gg U_{rp} = U_{pr}$ can be obeyed by regulating the interatomic distance and the atomic principal quantum numbers [59,60], whereby the terms $U_{rp}|rp\rangle\langle rp| + U_{pr}|pr\rangle\langle pr|$ can be neglected. Then taking into account $\Omega_1 = \Omega'_1$ and $\Omega_2 = \Omega'_2$, the full Hamiltonian can be rewritten as

$$H = H_X + H_D, \tag{3}$$

with

$$H_X = \sum_{j=A,B} \Omega_1(|r\rangle_j\langle 0| + |p\rangle_j\langle e|)e^{i\Delta_1 t} + \Omega_2(|r\rangle_j\langle e| + |p\rangle_j\langle 1|)e^{-i\Delta_2 t} + \text{H.c.} + U(|rr\rangle\langle rr| + |pp\rangle\langle pp|),$$

$$H_D = (\omega|s\rangle_A\langle e| + g|s\rangle_A\langle a|c)e^{i\lambda t} + \text{H.c.}$$

A. Rydberg-interaction-induced swap interaction

In this subsection, we elaborate on the operational principle of Rydberg-interaction-induced swap interaction only characterized by the H_X of Eq. (3), where the purpose is to change the initial state $|\psi\rangle_A|e\rangle_B|0\rangle_c$ into $|e\rangle_A|\psi\rangle_B|0\rangle_c$, i.e., actualizing the communication of $|\psi\rangle$ between atoms A and B. The cavity state of this subsection can be neglected because it is kept at the vacuum state $|0\rangle_c$.

Here, we consider $U = (\Delta_2 - \Delta_1) + \delta$ (δ is a difference between U and $\Delta_2 - \Delta_1$). In the limiting of $\{\Delta_1, \Delta_2, |\Delta_1 - \Delta_2|\} \gg \Omega_{1,2}$, the Hamiltonian H_X is further simplified as

$$H_X = \Omega'(|rr\rangle\langle 0e| + |rr\rangle\langle e0| + |pp\rangle\langle 1e| + |pp\rangle\langle e1|) + \text{H.c.} + (\delta + \delta')(|rr\rangle\langle rr| + |pp\rangle\langle pp|), \quad (4)$$

where we have ignored the high-frequency oscillating terms and abbreviated $\Omega' = \Omega_1\Omega_2/\Delta_2 - \Omega_1\Omega_2/\Delta_1$ and $\delta' = 2(\Omega_1^2/\Delta_2 - \Omega_2^2/\Delta_1)$. Besides, the Stark-shift terms $(\Omega_2^2/\Delta_2 - \Omega_1^2/\Delta_1)(|0e\rangle\langle 0e| + |e0\rangle\langle e0| + |1e\rangle\langle 1e| + |e1\rangle\langle e1|)$ are also omitted, since they can be compensated by the auxiliary levels. On the basis of Eq. (4), we further consider the case of $\delta + \delta' \gg \Omega'$, and the final form of the effective Hamiltonian for the Rydberg-interaction-induced swap interaction can be written as

$$H_{\text{eff}}^X = \Omega_{\text{eff}}(|e0\rangle\langle 0e| + |e1\rangle\langle 1e|) + \text{H.c.}, \quad (5)$$

where $\Omega'^2/(\delta + \delta')$ has been parametrized as Ω_{eff} and the corresponding Stark-shift terms have been disregarded. According to the H_{eff}^X, the swap of information between two systems, $|\psi\rangle_A|e\rangle_B \leftrightarrow |e\rangle_A|\psi\rangle_B$, will occur.

In Fig. 2, we show the evolution of populations for ρ_0 (empty squares) and ρ_R (empty circles) governed by the H_X of Eq. (3), where $\rho_0 = |\psi\rangle_A\langle\psi| \otimes |e\rangle_B\langle e| \otimes |0\rangle_c\langle 0|$ is the initial state and ρ_R is $|e\rangle_A\langle e| \otimes |\psi\rangle_B\langle\psi| \otimes |0\rangle_c\langle 0|$. We find ρ_0 and ρ_R interconvert into each other as predicted by Eq. (5), which reflects the transferred state $|\psi\rangle$ flows between atoms A and B successfully. The population of ρ_R arriving at 99.66% with $\Omega t = \pi\Omega/(2\Omega_{\text{eff}}) \approx 3040$ also confirms the validity of the Rydberg-interaction-induced swap interaction and the accuracy of the corresponding effective Hamiltonian.

B. Photon-loss-induced dissipative dynamics

Referring to the Rydberg-interaction-induced swap interaction, the Hamiltonian of the photon-loss-induced dissipative dynamics, H_D of Eq. (3), will only affect the states $|e0\rangle|0\rangle_c$ and $|e1\rangle|0\rangle_c$. On account of the large detuning condition $\lambda \gg \{\omega, g\}$, it can be simplified as (the Stark-shift terms have been ignored)

$$H_{\text{eff}}^D = \frac{\omega g}{\lambda}(|e0\rangle\langle a0| + |e1\rangle\langle a1| \otimes |0\rangle_c\langle 1| + \text{H.c.}, \quad (6)$$

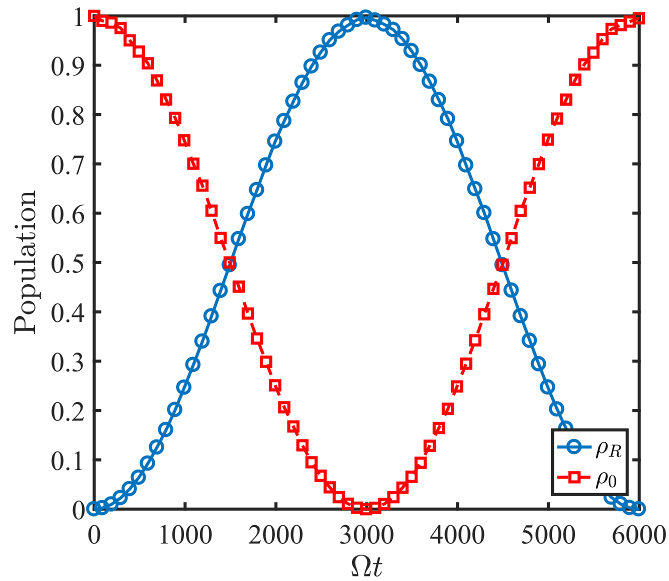

FIG. 2. The evolution of populations for ρ_0 and ρ_R governed by the H_X of Eq. (3). The relevant parameters are $\Omega_1 = \Omega_2 = \delta = \Omega$, $\Delta_1 = 30\Omega$, $\Delta_2 = 90\Omega$, $\alpha = \sin\theta$, $\beta = \cos\theta$, and $\theta = \pi/3$, where θ can be varied at will.

and this Hamiltonian will transmit the state $|e\rangle_A|\psi\rangle_B|0\rangle_c$ into $|a\rangle_A|\psi\rangle_B|1\rangle_c$. Subsequently, the photon loss of the cavity, characterized by $L_c = \sqrt{\kappa}c$ with the decay rate κ, will further decay the state $|a\rangle_A|\psi\rangle_B|1\rangle_c$ into the steady state $|a\rangle_A|\psi\rangle_B|0\rangle_c$. Ultimately, the system will be stable at $|a\rangle_A|\psi\rangle_B|0\rangle_c$ and the quantum state transfer will be finished.

Uniting the two compositions, we can obtain a full master equation of the total system as

$$\dot{\rho} = -i[H, \rho] + L_c\rho L_c^\dagger - \tfrac{1}{2}(L_c^\dagger L_c\rho + \rho L_c^\dagger L_c), \quad (7)$$

and the associated effective master equation reads

$$\dot{\rho} = -i[H_{\text{eff}}^X + H_{\text{eff}}^D, \rho] + L_c\rho L_c^\dagger - \tfrac{1}{2}(L_c^\dagger L_c\rho + \rho L_c^\dagger L_c). \quad (8)$$

It is evident that the state $|a\rangle_A|\psi\rangle_B|0\rangle_c$ is the steady state as the system initialized at $|e\rangle_A|\psi\rangle_B|0\rangle_c$. In Fig. 3, we inspect the time evolution of the populations for states ρ_0 (dotted line), ρ_R (dashed line), ρ_1 (dash-dotted line), and ρ_S (solid line) governed by the full master equation, Eq. (7), where $\rho_S = |a\rangle_A\langle a| \otimes |\psi\rangle_B\langle\psi| \otimes |0\rangle_c\langle 0|$ is the target state of the entire process and $\rho_1 = |a\rangle_A\langle a| \otimes |\psi\rangle_B\langle\psi| \otimes |1\rangle_c\langle 1|$. In view of the four curves, we find that with a decrease of the population of the initial state ρ_0, the system evolves into the state ρ_R by the Rydberg-interaction-induced swap interaction. Whereafter, taking advantage of the photon-loss-induced dissipative dynamics, ρ_R is driven into state ρ_1 which rapidly decays to the target state ρ_S. Ultimately, the system is stabilized into ρ_S with a high population of 99.99%; i.e., the quantum state $|\psi\rangle$ is directionally transferred from atom A to atom B without time-dependent external controls. Additionally, we also plot the populations of ρ_S (empty circles) as functions of Ωt governed by the effect master equation, Eq. (8). The brilliant agreement between the solid line and the empty circles significantly demonstrates the correctness of the reduced system. So we can accurately explicate the behaviors of the

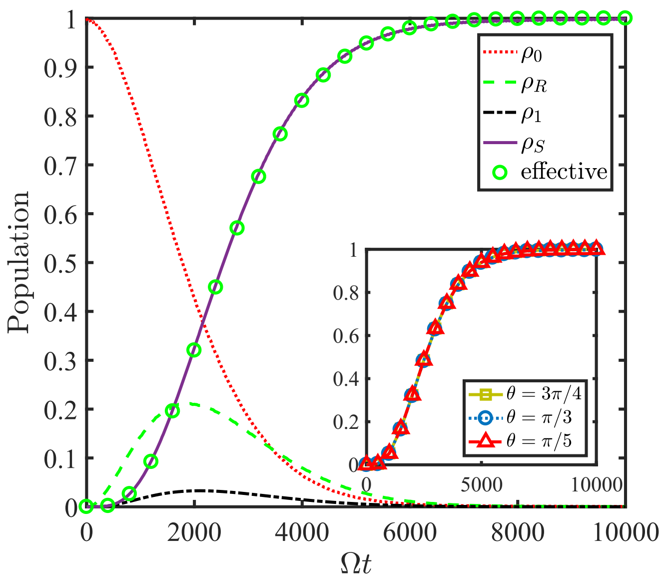

FIG. 3. The populations as functions of Ωt governed by the full master equation, Eq. (7), and the effective master equation, Eq. (8). The initial state is ρ_0 and the target state is ρ_S. The relevant parameters are set as $\Omega_1 = \Omega_2 = \omega = g = \delta = \Omega$, $\Delta_1 = 30\Omega$, $\Delta_2 = 90\Omega$, $\lambda = 500\Omega$, $\kappa = 0.01g$, $\alpha = \sin\theta$, $\beta = \cos\theta$, and $\theta = \pi/3$. The inset is the populations of ρ_S as functions of Ωt with different θ.

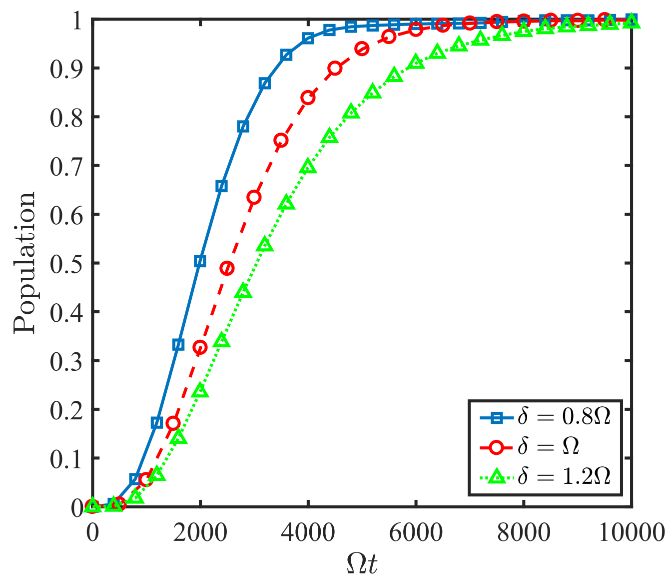

FIG. 4. The evolutions of the populations for ρ_S governed by the full master equation with different δ, where the initial state is ρ_0 and the other parameters are the same as those in Fig. 3.

original system by the effective master equation. The inset of Fig. 3 is the populations of ρ_S as functions of Ωt with different θ. It reflects that our scheme is independent of the values for α and β.

III. INFLUENCES OF RELEVANT PARAMETERS

As is well known, the Rydberg-mediated interaction will cause a large enough level shift to suppress the multiple atoms excited simultaneously, which is the marvellous Rydberg blockade effect [33]. On the contrary, one can realize the Rydberg antiblockade by the detuning of driving fields compensating the energy shift of Rydberg states [34] [equivalent to setting $U = (\Delta_2 - \Delta_1) + \delta$ in our scheme], which results in the simultaneous transitions of two atoms into Rydberg states and the inhibition of the Rydberg blockade. Distinctly, a perfect realization of the Rydberg antiblockade relies on the rigorous relations between the strength of Rydberg-mediated interaction and the detuning of the driving fields, which increases the experimental difficulty. However, there is no precisely tailored condition of the strength and detuning in our scheme, where we just need $U = (\Delta_2 - \Delta_1) + \delta$ and $\delta + \delta' \gg \Omega'$ (δ is not an exactly fixed value). In Fig. 4, we present the evolutions of the populations for ρ_S with different δ. And the qualities of the directional QST are all notably excellent, where the populations equal to 99.93% (empty squares), 99.92% (empty circles), and 99.23% (empty triangles) at $\Omega t = 10\,000$ illustrate the flexibility and feasibility of our scheme. In addition, under the limiting condition $\delta + \delta' \gg \Omega'$, the smaller the δ is, the shorter the convergence time is. This tendency can be interpreted by the $\Omega_{\rm eff}$ of Eq. (5). With the increase of δ, the decrease of $\Omega_{\rm eff}$ will retard the transport efficiency of the transferred quantum state.

In Fig. 5, we investigate the influence of the decay rate of the cavity by the population $\langle\psi|{\rm Tr}_{A\&c}[\rho]|\psi\rangle$ for the transferred state $|\psi\rangle$ in atom B with different κ, where ${\rm Tr}_{A\&c}$ denotes the partial trace over atom A and the cavity. Corresponding to $\kappa = 0.02g$, $0.015g$, and $0.01g$, the populations of $|\psi\rangle_B$ can be up to 98.97%, 98.38%, and 97.67% at $\Omega t = 6000$, respectively. The appearances also reliably affirm the feasibility of directional QST. Moreover, at the appropriate range of κ/g, the convergence time will be prolonged by weakening κ/g, since the photon loss of the cavity promotes

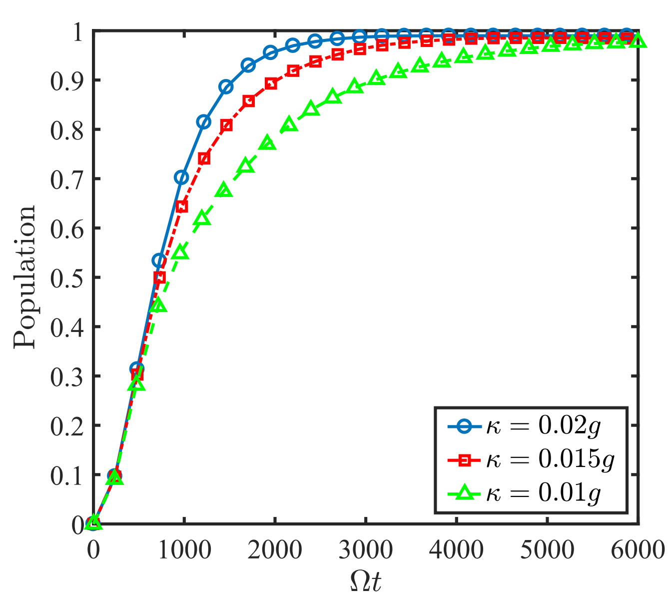

FIG. 5. The population of the transferred state in atom B, $\langle\psi|{\rm Tr}_{A\&c}[\rho]|\psi\rangle$, with different κ, where the initial state is ρ_0. The parameters are $\Omega_1 = \Omega_2 = \omega = g = \delta = \Omega$, $\Delta_1 = 20\Omega$, $\Delta_2 = 70\Omega$, $\lambda = 200\Omega$, $\alpha = \sin\theta$, $\beta = \cos\theta$, and $\theta = \pi/3$.

the quantum state $|\psi\rangle$ evolving into atom B. It is worth mentioning that a large κ/g will also delay the convergence time. Although there may be a group of optimal relevant parameters leading to the fastest convergence of steady state, the final quality of the QST is not affected.

In an experiment, one can always find a scenario to excite the atoms from ground states to Rydberg states by two-step transitions [36,61–63]. The first transition is dispersively pumping a ground state to an intermediate state by a laser with Rabi frequency Ω_a and detuning $-\Delta_a$. Then the intermediate state will be pumped to the desired Rydberg state via another laser with Rabi frequency Ω_b and detuning Δ_a. In the limiting of $|\Delta_a| \gg \Omega_{a,b}$, the intermediate state can be adiabatically eliminated. An equivalent direct transition from the ground state to the Rydberg state can be obtained with an effective Rabi frequency $\Omega_a\Omega_b/\Delta_a$, which is similar to the lasers with Rabi frequencies $\Omega_{1,2}$ for our scheme. So we can adjust the relation between $\Omega_{a,b}$ and Δ_a to obtain the desired values of $\Omega_{1,2}$ at will in an experiment. In addition, making use of two ^{87}Rb atoms individually trapped in optical tweezers at a distance of 4 μm, Gaëtan *et al.* [36] found that the strength of the Rydberg-mediated interaction (dipole-dipole type) between the two Rydberg states $|58d_{3/2}, F = 3, M_F = 3\rangle$ is equal to $U = 2\pi \times 50$ MHz. It can be calculated by C_3/r^3, where $C_3 \approx 2\pi \times 3200$ MHz μm^3. In Ref. [64], Müller *et al.* investigated the implementation of a controlled-Z gate on a pair of Rydberg atoms in spatially separated dipole traps at a distance of 0.3 μm, where the Rydberg states were the same as those of Ref. [36] and the corresponding interaction strength was 118 GHz. Therefore, we can consider that the Rydberg-mediated interaction for our scheme is dipole-dipole type, and the distance of two atoms can be changed into [0.3,4] μm to derive a suitable strength, $U = (\Delta_2 - \Delta_1) + \delta$. However, it should be noted that the conclusion of the work is independent of the interaction type (C_3/r^3 or C_6/r^6) and only the strength of the interaction at a given fixed distance of the two atoms is relevant.

On the other hand, selecting different atomic energy levels will correspond to different channels of atomic spontaneous emission during the course of an experiment. Nevertheless, the difference is negligible for the final results of our scheme, because the atomic spontaneous emission is significantly inhibited by the adiabatical elimination of the excited states and the Rydberg states. Thus we just consider the atomic spontaneous emission of our scheme consisting of $L_{1(2)} = \sqrt{\gamma/2}|0\rangle_{A(B)}\langle r|, L_{3(4)} = \sqrt{\gamma/2}|e\rangle_{A(B)}\langle r|$, $L_{5(6)} = \sqrt{\gamma/2}|e\rangle_{A(B)}\langle p|$, $L_{7(8)} = \sqrt{\gamma/2}|1\rangle_{A(B)}\langle p|$, $L_{9(10)} = \sqrt{\gamma_e/2}|e\rangle_{A(B)}\langle s|$, and $L_{11(12)} = \sqrt{\gamma_e/2}|a\rangle_{A(B)}\langle s|$, where the decay rates of the Rydberg states and the excited states $|s\rangle$ denoted by γ and γ_e can be chosen as $2\pi \times 1$ kHz and $2\pi \times 3$ MHz [54,65,66], respectively. Moreover the experiments of cavity QED with a Fabry-Perot cavity [67] and a microscopic optical resonator [68] provide us with two groups of parameters: $(g, \kappa) = 2\pi \times (770, 21.7)$ MHz and $(g, \kappa) = 2\pi \times (70, 5)$ MHz. Furthermore, another parameter can be supplied by the experimental scheme [69] with $(g, \kappa) = 2\pi \times (14.4, 0.66)$ MHz. In terms of the above state-of-the-art technologies and a full master equation including the atomic spontaneous emission, we depict the time evolutions of the populations for ρ_S with different experimental parameters in Fig. 6. At $gt = 3 \times 10^4$, the populations arrive at 99.98% (empty circles), 99.47% (empty squares), and 98.44% (empty triangles), respectively. The above upshot adequately indicates the experimental feasibility of our scheme.

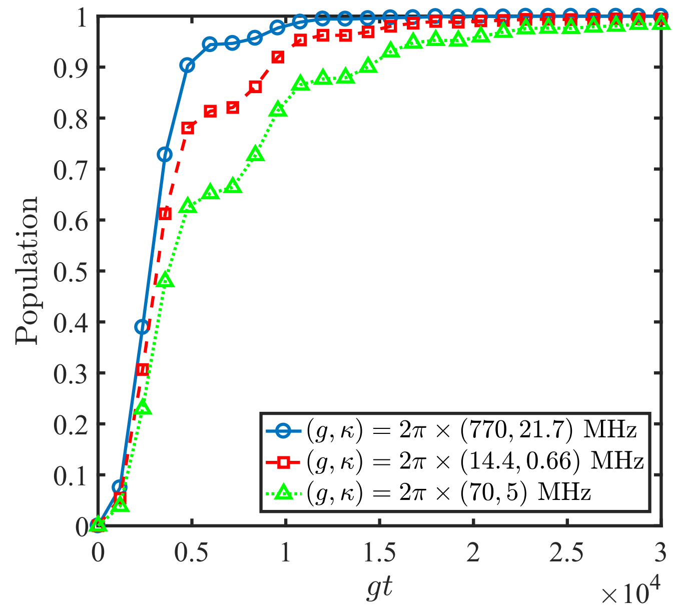

FIG. 6. The time evolutions of the populations for ρ_S governed by a full master equation including the atomic spontaneous emission with different experimental parameters. The initial state is ρ_0 and the other parameters are chosen as $\Omega_1 = \Omega_2 = \omega = \delta = g$, $\Delta_1 = 30g$, $\Delta_2 = 90g$, $\lambda = 400g$, $\gamma_e = 2\pi \times 3$ MHz, $\gamma = 2\pi \times 1$ kHz, $\alpha = \sin\theta$, $\beta = \cos\theta$, and $\theta = \pi/3$.

IV. GENERALIZED SCHEME TO RESIST THE COLLECTIVE DEPHASING FROM THE ENVIRONMENT

In this section, we investigate the influence of relative phase noise between the $|0\rangle$ and $|1\rangle$ states, i.e., the collective dephasing resulting from the collective interaction of the total system with a common environment [70–72]. In our system, it can be described by a Lindblad operator as

$$L_d = \sqrt{\kappa_d}\sum_{j=1}^{N}(|0\rangle_j\langle 0| - |1\rangle_j\langle 1|), \tag{9}$$

where κ_d is the collective dephasing rate and N is the number of atoms. Quite evidently, our method with two atoms is fragile for the effect of collective dephasing. Therefore, we generalize our scheme from two identical atoms to four identical atoms. The atomic energy levels are the same as those of Fig. 1(b). The generalized scheme aims to transfer an arbitrary state, $|\psi\rangle = \alpha|0\rangle_L + \beta|1\rangle_L = \alpha|01\rangle_P + \beta|10\rangle_P$, from system A consisting of atoms 1 and 2 to system B composed of atoms 3 and 4, where the subscripts L and P stand for logical and physical. It is significant that the negative effects of collective dephasing can be efficaciously overcome since there is only a common phase between states $|01\rangle_P$ and $|10\rangle_P$.

In particular, the first process to realize the generalized scheme is plotted in Fig. 7(a). We apply all the operations

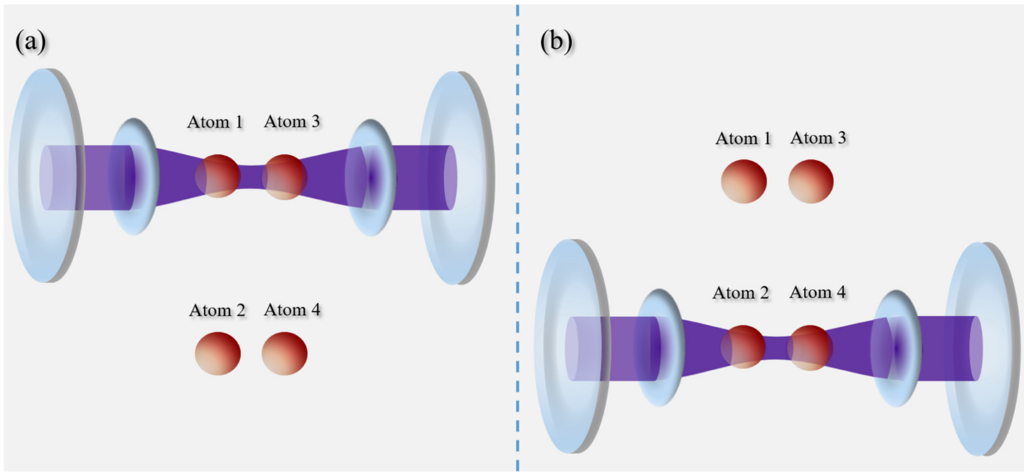

FIG. 7. The diagram of the generalized scheme to resist the collective dephasing from the environment, where the purpose is to transfer an arbitrary state, $|\psi\rangle = \alpha|0\rangle_L + \beta|1\rangle_L = \alpha|01\rangle_P + \beta|10\rangle_P$, from system A consisting of atoms 1 and 2 to system B composed of atoms 3 and 4. (a) First, only apply all the operations identical to those of Fig. 1(a) to atoms 1 and 3. (b) Second, remove all the operations from atoms 1 and 3 to atoms 2 and 4.

identical to those of Fig. 1(a) to atoms 1 and 3, where the initial state of the total system is $|\psi_0\rangle = (\alpha|01\rangle_{12} + \beta|10\rangle_{12}) \otimes |ee\rangle_{34} \otimes |0\rangle_c$. Building on the Rydberg-interaction-induced swap interaction and the photon-loss-induced dissipative dynamics, the information of atoms 1 and 3 can be swapped and the system will be stabilized at the state $|\psi_1\rangle = (\alpha|a10e\rangle_{1234} + \beta|a01e\rangle_{1234}) \otimes |0\rangle_c$. Second, as exhibited in Fig. 7(b), we remove all the operations from atoms 1 and 3 to atoms 2 and 4. The state $|\psi_1\rangle$ will evolve into the target state $|\psi_S\rangle = |aa\rangle_{12} \otimes (\alpha|01\rangle_{34} + \beta|10\rangle_{34}) \otimes |0\rangle_c$, and the QST of an arbitrary state $|\psi\rangle$ between two systems is achieved. Naturally, the two steps can be also performed at the same time with two cavities respectively trapping atoms 1 and 3 and atoms 2 and 4.

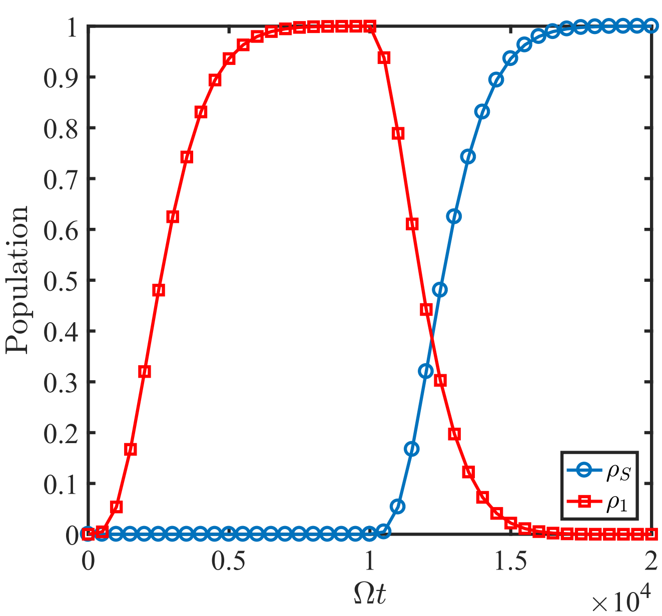

FIG. 8. The dynamics evolutions of populations for ρ_S and ρ_1. The relevant parameters are equal to those of Fig. 3. In addition, the collective dephasing rate is chosen as $\kappa_d = 0.1\Omega$.

To testify the feasibility of the generalized scheme, we utilize the master equation including the terms of collective dephasing, $L_d\rho L_d^\dagger - (L_d^\dagger L_d\rho + \rho L_d^\dagger L_d)/2$, to inspect the behaviors of populations for $\rho_S = |\psi_S\rangle\langle\psi_S|$ (empty circles) and $\rho_1 = |\psi_1\rangle\langle\psi_1|$ (empty squares) in Fig. 8. At the timescale $\Omega t \in [0, 10\,000)$, we implement the first step of Fig. 7(a), and the system will be steady at $|\psi_1\rangle$ with a high population. Furthermore, we carry out the operations as plotted in Fig. 7(b) while $\Omega t \in [10\,000, 20\,000]$. The target state ρ_S is prepared on schedule. It affirms that the generalized scheme not only actualizes a dissipation-assisted QST but also possesses outstanding robustness against the collective dephasing.

V. SUMMARY

In summary, we successfully show that it is possible to dissipatively transfer an arbitrary quantum state from one place to another without precise time-dependent controls, where we organically integrate the Rydberg-antiblockade effect, the laser-induced Raman transition, and the photon loss of an optical cavity to make the target state become the unique steady state of the whole process. Owing to the adiabatical elimination of the excited states and the Rydberg states, the adverse effect of the atomic spontaneous emission is also depressed pronouncedly. We also discuss the influence of the relevant parameters and demonstrate the feasibility of the scheme by using state-of-the-art technologies and showing that a high population of the transferred state around 99% is achievable. Finally, we design a generalized scheme to resist the collective dephasing from the environment. We believe our scheme supplies a viable prospect with regard to QST via quantum dissipation.

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[1] J. I. Cirac, P. Zoller, H. J. Kimble, and H. Mabuchi, Quantum State Transfer and Entanglement Distribution among Distant Nodes in a Quantum Network, Phys. Rev. Lett. **78**, 3221 (1997).

[2] S. Bose, Quantum Communication through an Unmodulated Spin Chain, Phys. Rev. Lett. **91**, 207901 (2003).

[3] D. N. Matsukevich and A. Kuzmich, Quantum state transfer between matter and light, Science **306**, 663 (2004).

[4] M. Christandl, N. Datta, A. Ekert, and A. J. Landahl, Perfect State Transfer in Quantum Spin Networks, Phys. Rev. Lett. **92**, 187902 (2004).

[5] M. A. Sillanpää, J. I. Park, and R. W. Simmonds, Coherent quantum state storage and transfer between two phase qubits via a resonant cavity, Nature (London) **449**, 438 (2007).

[6] J. Majer, J. M. Chow, J. M. Gambetta, J. Koch, B. R. Johnson, J. A. Schreier, L. Frunzio, D. I. Schuster, A. A. Houck, A. Wallraff, A. Blais, M. H. Devoret, S. M. Girvin, and R. J. Schoelkopf, Coupling superconducting qubits via a cavity bus, Nature (London) **449**, 443 (2007).

[7] N. Y. Yao, L. Jiang, A. V. Gorshkov, Z.-X. Gong, A. Zhai, L.-M. Duan, and M. D. Lukin, Robust Quantum State Transfer in Random Unpolarized Spin Chains, Phys. Rev. Lett. **106**, 040505 (2011).

[8] B. Vermersch, P.-O. Guimond, H. Pichler, and P. Zoller, Quantum State Transfer via Noisy Photonic and Phononic Waveguides, Phys. Rev. Lett. **118**, 133601 (2017).

[9] Z.-L. Xiang, M. Zhang, L. Jiang, and P. Rabl, Intracity Quantum Communication via Thermal Microwave Networks, Phys. Rev. X **7**, 011035 (2017).

[10] F. Mei, G. Chen, L. Tian, S.-L. Zhu, and S. Jia, Robust quantum state transfer via topological edge states in superconducting qubit chains, Phys. Rev. A **98**, 012331 (2018).

[11] H. P. Breuer and F. Petruccione, *The Theory of Open Quantum Systems* (Oxford University, Oxford, England, 2002).

[12] J. F. Poyatos, J. I. Cirac, and P. Zoller, Quantum Reservoir Engineering with Laser Cooled Trapped Ions, Phys. Rev. Lett. **77**, 4728 (1996).

[13] J. P. Paz and W. H. Zurek, Continuous error correction, Proc. R. Soc. London, Ser. A **454**, 355 (1998).

[14] M. B. Plenio, S. F. Huelga, A. Beige, and P. L. Knight, Cavity-loss-induced generation of entangled atoms, Phys. Rev. A **59**, 2468 (1999).

[15] M. B. Plenio and S. F. Huelga, Entangled Light from White Noise, Phys. Rev. Lett. **88**, 197901 (2002).

[16] C. Ahn, A. C. Doherty, and A. J. Landahl, Continuous quantum error correction via quantum feedback control, Phys. Rev. A **65**, 042301 (2002).

[17] X. X. Yi, C. S. Yu, L. Zhou, and H. S. Song, Noise-assisted preparation of entangled atoms, Phys. Rev. A **68**, 052304 (2003).

[18] M. J. Kastoryano, F. Reiter, and A. S. Sørensen, Dissipative Preparation of Entanglement in Optical Cavities, Phys. Rev. Lett. **106**, 090502 (2011).

[19] Y. Lin, J. P. Gaebler, F. Reiter, T. R. Tan, R. Bowler, A. S. Sørensen, D. Leibfried, and D. J. Wineland, Dissipative production of a maximally entangled steady state of two quantum bits, Nature (London) **504**, 415 (2013).

[20] A. W. Carr and M. Saffman, Preparation of Entangled and Antiferromagnetic States by Dissipative Rydberg Pumping, Phys. Rev. Lett. **111**, 033607 (2013).

[21] K. Fujii, M. Negoro, N. Imoto, and M. Kitagawa, Measurement-Free Topological Protection Using Dissipative Feedback, Phys. Rev. X **4**, 041039 (2014).

[22] Z. Leghtas, S. Touzard, I. M. Pop, A. Kou, B. Vlastakis, A. Petrenko, K. M. Sliwa, A. Narla, S. Shankar, M. J. Hatridge, M. Reagor, L. Frunzio, R. J. Schoelkopf, M. Mirrahimi, and M. H. Devoret, Confining the state of light to a quantum manifold by engineered two-photon loss, Science **347**, 853 (2015).

[23] F. Reiter, D. Reeb, and A. S. Sørensen, Scalable Dissipative Preparation of Many-Body Entanglement, Phys. Rev. Lett. **117**, 040501 (2016).

[24] F. Reiter, A. S. Sørensen, P. Zoller, and C. A. Muschik, Dissipative quantum error correction and application to quantum sensing with trapped ions, Nat. Commun. **8**, 1822 (2017).

[25] X. Q. Shao, J. H. Wu, X. X. Yi, and G.-L. Long, Dissipative preparation of steady Greenberger-Horne-Zeilinger states for Rydberg atoms with quantum zeno dynamics, Phys. Rev. A **96**, 062315 (2017).

[26] D. X. Li, X. Q. Shao, J. H. Wu, and X. X. Yi, Engineering steady-state entanglement via dissipation and quantum Zeno dynamics in an optical cavity, Opt. Lett. **42**, 3904 (2017).

[27] D.-X. Li, X.-Q. Shao, J.-H. Wu, and X. X. Yi, Dissipation-induced *W* state in a Rydberg-atom-cavity system, Opt. Lett. **43**, 1639 (2018).

[28] X.-Q. Shao, Engineering steady entanglement for trapped ions at finite temperature by dissipation, Phys. Rev. A **98**, 042310 (2018).

[29] D.-X. Li, X.-Q. Shao, J.-H. Wu, X. X. Yi, and T.-Y. Zheng, Engineering steady Knill-Laflamme-Milburn state of Rydberg atoms by dissipation, Opt. Express **26**, 2292 (2018).

[30] C. Wang and J. M. Gertler, Autonomous quantum state transfer by dissipation, arXiv:1809.03571.

[31] Y. Matsuzaki, V. M. Bastidas, Y. Takeuchi, W. J. Munro, and S. Saito, One-way transfer of quantum states via decoherence, arXiv:1810.02995.

[32] M. Saffman, T. G. Walker, and K. Mølmer, Quantum information with Rydberg atoms, Rev. Mod. Phys. **82**, 2313 (2010).

[33] D. Jaksch, J. I. Cirac, P. Zoller, S. L. Rolston, R. Côté, and M. D. Lukin, Fast Quantum Gates for Neutral Atoms, Phys. Rev. Lett. **85**, 2208 (2000).

[34] C. Ates, T. Pohl, T. Pattard, and J. M. Rost, Antiblockade in Rydberg Excitation of an Ultracold Lattice Gas, Phys. Rev. Lett. **98**, 023002 (2007).

[35] E. Urban, T. A. Johnson, T. Henage, L. Isenhower, D. D. Yavuz, T. G. Walker, and M. Saffman, Observation of Rydberg blockade between two atoms, Nat. Phys. **5**, 110 (2009).

[36] A. Gaëtan, Y. Miroshnychenko, T. Wilk, A. Chotia, M. Viteau, D. Comparat, P. Pillet, A. Browaeys, and P. Grangier, Observation of collective excitation of two individual atoms in the Rydberg blockade regime, Nat. Phys. **5**, 115 (2009).

[37] M. Saffman and K. Mølmer, Efficient Multiparticle Entanglement via Asymmetric Rydberg Blockade, Phys. Rev. Lett. **102**, 240502 (2009).

[38] H. Weimer, M. Muller, I. Lesanovsky, P. Zoller, and H. P. Buchler, A Rydberg quantum simulator, Nat. Phys. **6**, 382 (2010).

[39] T. Amthor, C. Giese, C. S. Hofmann, and M. Weidemüller, Evidence of Antiblockade in an Ultracold Rydberg Gas, Phys. Rev. Lett. **104**, 013001 (2010).

[40] M. Mayle, W. Zeller, N. Tezak, and P. Schmelcher, Rydberg-Rydberg interaction profile from the excitation dynamics of ultracold atoms in lattices, Phys. Rev. A **84**, 010701 (2011).

[41] L. Béguin, A. Vernier, R. Chicireanu, T. Lahaye, and A. Browaeys, Direct Measurement of the van der Waals Interaction between Two Rydberg Atoms, Phys. Rev. Lett. **110**, 263201 (2013).

[42] D. D. B. Rao and K. Mølmer, Dark Entangled Steady States of Interacting Rydberg Atoms, Phys. Rev. Lett. **111**, 033606 (2013).

[43] Y.-Y. Jau, A. M. Hankin, T. Keating, I. H. Deutsch, and G. W. Biedermann, Entangling atomic spins with a Rydberg-dressed spin-flip blockade, Nat. Phys. **12**, 71 (2016).

[44] S.-L. Su, E. Liang, S. Zhang, J.-J. Wen, L.-L. Sun, Z. Jin, and A.-D. Zhu, One-step implementation of the Rydberg-Rydberg-interaction gate, Phys. Rev. A **93**, 012306 (2016).

[45] S.-L. Su, Y. Gao, E. Liang, and S. Zhang, Fast Rydberg antiblockade regime and its applications in quantum logic gates, Phys. Rev. A **95**, 022319 (2017).

[46] J. Song, C. Li, Z.-J. Zhang, Y.-Y. Jiang, and Y. Xia, Implementing stabilizer codes in noisy environments, Phys. Rev. A **96**, 032336 (2017).

[47] S.-L. Su, Y. Tian, H. Z. Shen, H. Zang, E. Liang, and S. Zhang, Applications of the modified Rydberg antiblockade regime with simultaneous driving, Phys. Rev. A **96**, 042335 (2017).

[48] I. I. Beterov, G. N. Hamzina, E. A. Yakshina, D. B. Tretyakov, V. M. Entin, and I. I. Ryabtsev, Adiabatic passage of radio-frequency-assisted Förster resonances in Rydberg atoms for two-qubit gates and the generation of Bell states, Phys. Rev. A **97**, 032701 (2018).

[49] J. T. Young, T. Boulier, E. Magnan, E. A. Goldschmidt, R. M. Wilson, S. L. Rolston, J. V. Porto, and A. V. Gorshkov, Dissipation-induced dipole blockade and antiblockade in driven Rydberg systems, Phys. Rev. A **97**, 023424 (2018).

[50] S. L. Su, H. Z. Shen, E. Liang, and S. Zhang, One-step construction of the multiple-qubit Rydberg controlled-phase gate, Phys. Rev. A **98**, 032306 (2018).

[51] D. X. Li and X. Q. Shao, Unconventional Rydberg pumping and applications in quantum information processing, Phys. Rev. A **98**, 062338 (2018).

[52] C. Guerlin, E. Brion, T. Esslinger, and K. Mølmer, Cavity quantum electrodynamics with a Rydberg-blocked atomic ensemble, Phys. Rev. A **82**, 053832 (2010).

[53] V. Parigi, E. Bimbard, J. Stanojevic, A. J. Hilliard, F. Nogrette, R. Tualle-Brouri, A. Ourjoumtsev, and P. Grangier, Observation and Measurement of Interaction-Induced Dispersive Optical Nonlinearities in an Ensemble of Cold Rydberg Atoms, Phys. Rev. Lett. **109**, 233602 (2012).

[54] X.-F. Zhang, Q. Sun, Y.-C. Wen, W.-M. Liu, S. Eggert, and A.-C. Ji, Rydberg Polaritons in a Cavity: A Superradiant Solid, Phys. Rev. Lett. **110**, 090402 (2013).

[55] J.-F. Huang, J.-Q. Liao, and C. P. Sun, Photon blockade induced by atoms with Rydberg coupling, Phys. Rev. A **87**, 023822 (2013).

[56] H. Wu, Z.-B. Yang, and S.-B. Zheng, Two-photon absorption and emission by Rydberg atoms in coupled cavities, Phys. Rev. A **88**, 043816 (2013).

[57] M. F. Maghrebi, N. Y. Yao, M. Hafezi, T. Pohl, O. Firstenberg, and A. V. Gorshkov, Fractional quantum Hall states of Rydberg polaritons, Phys. Rev. A **91**, 033838 (2015).

[58] J. Sheng, Y. Chao, S. Kumar, H. Fan, J. Sedlacek, and J. P. Shaffer, Intracavity Rydberg-atom electromagnetically induced transparency using a high-finesse optical cavity, Phys. Rev. A **96**, 033813 (2017).

[59] T. G. Walker and M. Saffman, Consequences of Zeeman degeneracy for the van der Waals blockade between Rydberg atoms, Phys. Rev. A **77**, 032723 (2008).

[60] Y.-M. Liu, X.-D. Tian, D. Yan, Y. Zhang, C.-L. Cui, and J.-H. Wu, Nonlinear modifications of photon correlations via controlled single and double Rydberg blockade, Phys. Rev. A **91**, 043802 (2015).

[61] T. Wilk, A. Gaëtan, C. Evellin, J. Wolters, Y. Miroshnychenko, P. Grangier, and A. Browaeys, Entanglement of Two Individual Neutral Atoms Using Rydberg Blockade, Phys. Rev. Lett. **104**, 010502 (2010).

[62] G. Higgins, W. Li, F. Pokorny, C. Zhang, F. Kress, C. Maier, J. Haag, Q. Bodart, I. Lesanovsky, and M. Hennrich, Single Strontium Rydberg Ion Confined in a Paul Trap, Phys. Rev. X **7**, 021038 (2017).

[63] X. Q. Shao, J. H. Wu, and X. X. Yi, Dissipative stabilization of quantum-feedback-based multipartite entanglement with Rydberg atoms, Phys. Rev. A **95**, 022317 (2017).

[64] M. M. Müller, M. Murphy, S. Montangero, T. Calarco, P. Grangier, and A. Browaeys, Implementation of an experimentally feasible controlled-phase gate on two blockaded Rydberg atoms, Phys. Rev. A **89**, 032334 (2014).

[65] F. Brennecke, T. Donner, S. Ritter, T. Bourdel, M. Köhl, and T. Esslinger, Cavity QED with a Bose-Einstein condensate, Nature (London) **450**, 268 (2007).

[66] A. Grankin, E. Brion, E. Bimbard, R. Boddeda, I. Usmani, A. Ourjoumtsev, and P. Grangier, Quantum statistics of light transmitted through an intracavity Rydberg medium, New J. Phys. **16**, 043020 (2014).

[67] S. M. Spillane, T. J. Kippenberg, K. J. Vahala, K. W. Goh, E. Wilcut, and H. J. Kimble, Ultrahigh-q toroidal microresonators for cavity quantum electrodynamics, Phys. Rev. A **71**, 013817 (2005).

[68] B. Dayan, A. S. Parkins, T. Aoki, E. P. Ostby, K. J. Vahala, and H. J. Kimble, A photon turnstile dynamically regulated by one atom, Science **319**, 1062 (2008).

[69] K. W. Murch, K. L. Moore, S. Gupta, and D. M. Stamper-Kurn, Observation of quantum-measurement backaction with an ultracold atomic gas, Nat. Phys. **4**, 561 (2008).

[70] T. Yu and J. H. Eberly, Phonon decoherence of quantum entanglement: Robust and fragile states, Phys. Rev. B **66**, 193306 (2002).

[71] D. Braun, Creation of Entanglement by Interaction with a Common Heat Bath, Phys. Rev. Lett. **89**, 277901 (2002).

[72] T. Yu and J. H. Eberly, Qubit disentanglement and decoherence via dephasing, Phys. Rev. B **68**, 165322 (2003).