Selective Rydberg pumping via strong dipole blockade

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The resonant dipole-dipole interaction between highly excited Rydberg levels dominates the interaction of neutral atoms at short distances scaling as \(1/r^3\). Here we take advantage of the combined effects of strong dipole-dipole interaction and multifrequency driving fields to propose one type of selective Rydberg pumping mechanism. In the ground state space of two atoms \(\{|00\rangle, |01\rangle, |10\rangle, |11\rangle\}\), this mechanism allows \(|11\rangle\) to be resonantly pumped upwards to the single-excited Rydberg states while the transitions of the other three states are suppressed. From the perspective of mathematical form, we achieve an analogous Förster resonance for ground states of neutral atoms. Thence we apply this mechanism to one-step implementation of controlled-Z gate, realization of ground-state blockade, and speedup preparation of maximally entangled state with dissipation, and further discuss respectively each feasibility with the state-of-the-art technology.

I. INTRODUCTION

Rydberg blockade and Rydberg antiblockade are two representative phenomena observed in neutral atom systems [1–6]. The former prevents two or more atoms from being resonantly excited to the Rydberg state due to the strong Rydberg-Rydberg interaction (RRI), while the later permits a resonant two-photon transition as the shifting energy of Rydberg states is compensated by the two-photon detuning. In addition to the above two characteristics, the highly excited Rydberg states possess a long lifetime proportional to the third power of the principle quantum number [7–9], which makes them very suitable for encoding qubits or as a medium in the field of quantum information science [10–21].

Nevertheless, the spontaneous emission of Rydberg state cannot be ignored once the long-time dynamic evolution is considered, such as a process of a many-body system reaching equilibrium. Using a continuous laser field coupling the ground state to the Rydberg state can form two dressed states, these Rydberg-dressed states have relative long lifetimes as compared to the bare Rydberg state, while maintains a certain strength of the RRI to induce the blockade effect [22, 23]. Therefore the Rydberg dressing technique provides a tunable interaction for quantum control and quantum computing [24–33]. It should be emphasized that at the end of each Rydberg-dressed scheme, an additional operation that adiabatically transferring the Rydberg-dressed state back to the ground state is necessary to avoid further decoherence.

Different from the Rydberg dressing techniques, the spontaneous emission of the Rydberg state can be neglectable during the whole process of the ground-state blockade mechanism [34]. The transitions between ground states governed by the weak Raman couplings are “frequently measured” by a strong Rydberg antiblockade interaction. Thence the Rydberg states are virtually excited and the population of a two-atom ground state that directly coupled to the two-atom Rydberg state is blockaded. The prominent advantage of the ground-state blockade is that although the strength of the Rydberg antiblockade is usually much smaller than the Rabi coupling of Rydberg atoms, it can be made far greater than the Raman coupling of ground state, and simultaneously reduces the atomic spontaneous emission for the Raman coupling dynamics [35, 36].

In this paper, we aim to engineer an alternative interaction mechanism for neutral atom systems, which is able to replace the traditional second-order dynamics in the Rydberg antiblockade with first-order interaction strength. This idea mainly comes from our precious work about the unconventional Rydberg pumping (URP) [37], where the evolution of two atoms initialized in the same ground state is frozen even under driven by laser fields. Instead of using the van der Waals interaction of Rydberg atoms for URP, here we take advantage of the combined effects of strong dipole-dipole interaction (which is not replaceable) and multifrequency driving fields to propose a selective Rydberg pumping (SRP) mechanism, i.e., in the ground state space of two atoms \(\{|00\rangle, |01\rangle, |10\rangle, |11\rangle\}\), this mechanism allows \(|11\rangle\) to be resonantly pumped upwards to the single-excited Rydberg states while the transitions of the other three states are suppressed. Since the dynamics is dominated by the first-order Rabi coupling of Rydberg atoms, our work may have a great application value on speeding up the Rydberg-antiblockade-based schemes [12, 13, 20, 34].

The structure of the paper is organized as follows. In Sec. II, we first illustrate the mechanism of the SRP in detail and explain why the dipole-dipole interaction of Rydberg atoms cannot be replaced by the van der Waals interaction. In Sec. III, we then apply the technology of SRP to one-step implementation of controlled-Z gate, realization of ground-state blockade, and dissipative preparation of maximally entangled state with available experimental parameters and compare them with the relevant Rydberg-antiblockade-based schemes. In Sec. IV, we further discuss the effects of imperfection such as Förster defect and decoherence on the performance of the SRP, and finally give a summary of our manuscript in Sec. V.
In general the strong dipole-dipole interaction between Rydberg atoms can result in a splitting of the excited Rydberg states. Thus, this part can be reformulated in the rotating frame $U = \exp[-\sqrt{2}Jt(|E_+\rangle\langle E_+| - |E_-\rangle\langle E_-|)]$, the above Hamiltonian is written as

$$H_1 = H_1 + H_2,$$

$$H_1 = \sqrt{2}\Omega_s \langle 00 \rangle \left\langle |E_+\rangle e^{-i\Delta t} + |E_-\rangle e^{i\Delta t} \right\rangle$$

$$+ |01 \rangle \left[ \frac{\Omega}{\sqrt{2}}(|T_0 \rangle - |S_0 \rangle) + 2\Omega_s \cos (\Delta t)|1r \rangle \right]$$

$$+ |10 \rangle \left[ \frac{\Omega}{\sqrt{2}}(|T_0 \rangle + |S_0 \rangle) - 2i\Omega_s \sin (\Delta t)|1r \rangle \right]$$

$$+ \Omega|11 \rangle (|r \rangle + |1r \rangle) + \text{H.c.},$$

where $H_1$ describes the interaction between ground states and single-excited Rydberg states, and $H_2$ bridges the transitions of single-excited and two-excited Rydberg states. In the regime of the large detuning limits $\Delta = \sqrt{2}J$ and $\Delta \gg \Omega_s, \Omega_s$, the terms oscillating with high frequencies $\{\Omega_s, \Omega_s, \Omega_s + \Omega_s + \Delta + \sqrt{2}J\}$ in Eq. (2) can be safely disregarded, then we have a concise form as

$$H_1 = \Omega|11 \rangle (|r \rangle + |1r \rangle) + \frac{\Omega}{\sqrt{2}}|01 \rangle (|T_0 \rangle - |S_0 \rangle)$$

$$\Omega_s |10 \rangle (|T_0 \rangle + |S_0 \rangle) + \Omega_s |T_0 \rangle |E_+ \rangle$$

$$\Omega_s |S_0 \rangle |E_- \rangle + \text{H.c.}. \quad (3)$$

This Hamiltonian can be further simplified by considering the limitation of $\Omega_s \gg \Omega_s$. Analogous to the process of dipole blockade, the strong Rabi coupling $\Omega_s$ leads to large shifting energies of states $|T_0 \rangle$ and $|S_0 \rangle$, which is big enough to block the jumps from $|01 \rangle (|10 \rangle)$ towards to the excited states driven by Rabi frequency $\Omega_s$. Most importantly, there is no extra stark shifts or detuning-induced Raman resonance, because these transition paths mediated by the independent channels $|\alpha_\pm \rangle$ and $|\beta_\pm \rangle$ interfere destructively, where

$$|\alpha_\pm \rangle = \frac{1}{2}(|T_0 \rangle - |S_0 \rangle) \pm (|E_+ \rangle + |E_- \rangle),$$

$$|\beta_\pm \rangle = \frac{1}{2}(|T_0 \rangle + |S_0 \rangle) \mp (|E_+ \rangle - |E_- \rangle),$$

are the corresponding eigenstates of the part governed by $\Omega_s$. Therefore the Hamiltonian of our current model reduces to an effective form

$$H_{\text{eff}} = \Omega|11 \rangle (|r \rangle + |1r \rangle) + \text{H.c.}. \quad (4)$$

Now we finish the mechanism of SRP. The effective transitions for all ground states are also displayed in Fig. 1(c) in order to offer a better physical picture. Note that there are many
we also analyze the dynamical form, it describes an analogous Föster resonance, where the ground states \( |11 \rangle \) are coupled to other two states resonantly. So one potential application of our SRP is to simulate the Föster resonance-related phenomena with ground states of neutral atoms. In addition, the first-order Rabi coupling is usually much larger than the second-order interaction as required by the Rydberg antiblockade, the SRP mechanism may greatly speed up the Rydberg-antiblockade-based schemes, while reducing the decay of other excited Rydberg states.

Before we go on to the next step, let us first discuss the feasibility of the SRP with typical parameters used in experiments. According to the works of Browaeys et al. [38, 39], the Föster resonance of excited Rydberg states is reached using \( |r' \rangle = |61P_{5/2}, m_J = 1/2, \rangle, |r \rangle = |59D_{3/2}, m_J = 3/2, \rangle, \) and \( |r'' \rangle = |57F_{5/2}, m_J = 5/2, \rangle \) of two \(^{87}\)Rb atoms in the presence of an electric field and the measured \( C_3/2\pi = 2.39 \pm 0.03 \text{ GHz } \mu\text{m}^3 \) which is close to the theoretical value \( C_3/2\pi \approx 2.54 \text{ GHz } \mu\text{m}^3 \). This allows the dipole interaction strength \( J/2\pi \) to be continuously varied between 2.39 MHz and 152.96 MHz corresponding to the distance between atoms are adjusted from 10 \( \mu\text{m} \) to 2.5 \( \mu\text{m} \). The Rabi coupling of the ground states and the Rydberg states is accomplished by a tunable two-photon process that can be up to \( 2\pi \times 5 \text{ MHz} \) in Ref. [39]. So the condition for realization of the SRP \( (\Delta = \sqrt{2J} \gg \Omega_\chi \gg \Omega) \) is easy to access by choosing \( \Omega/2\pi = 0.02 \text{ MHz}, \Omega_\chi/2\pi = 1 \text{ MHz}, \) and \( J/2\pi = 50 \text{ MHz} \). In the left panel of Fig. 2, we depict the temporal evolution of all ground states obtained from the full Hamiltonian of Eq. (1). It can be seen that this result has an excellent agreement with our prediction. The states \(|00\rangle, |01\rangle, \) and \(|10\rangle \) always keep in their initial states with fidelities higher than 99.5% during the coherent-oscillation process of states \(|11\rangle \) and \((|r'1\rangle + |r1\rangle)/\sqrt{2} \). As we mentioned earlier, the dipole-dipole interaction is the key factor to implement the SRP, which cannot be substituted by the van der Waals-type interaction. In the right panel of Fig. 2 we also analyze the dynamical behavior of each ground state based on the van der Waals interaction \((U = J = 2\pi \times 50 \text{ MHz})\) as a comparison. In this case, the energy of the excited Rydberg states will only shift without splitting, which results in an undesired resonant transition between states \((|01\rangle - |10\rangle)/\sqrt{2} \) and \((|r'\rangle - |r1\rangle)/\sqrt{2} \).

### III. Applications

#### A. Controlled-Z gate

The mechanism of SRP itself defines a kind of quantum logic operation, i.e. \( |00\rangle \rightarrow |00\rangle, |01\rangle \rightarrow |01\rangle, |10\rangle \rightarrow |10\rangle, \) and \( |11\rangle \rightarrow \cos(\sqrt{2}\Omega t)|11\rangle + i\sin(\sqrt{2}\Omega t)(|r'1\rangle + |r1\rangle)/\sqrt{2} \) Thus the two-qubit controlled-Z gate corresponding \( \frac{\sqrt{2}}{2}\pi t = \pi \) is readily implemented only in one step. Without loss of generality, we use the fidelity between quantum states \( \langle \psi_{\text{ideal}} | \rho(t) | \psi_{\text{ideal}} \rangle \) as a measure to check the performance of the gate, where the initial (ideal final state) is \((|00\rangle + |01\rangle + |10\rangle + |11\rangle)/2 \). In Fig. 3(a) we employ the same group of parameters used in Fig. 2 to obtain a nearly perfect two-qubit controlled-Z gate with fidelity of 99.9% at \( t = 17.7 \mu\text{s} \). In Fig. 3(b) we further enhance the Rabi frequency of the resonant driving field to \( \Omega/2\pi = 0.06 \text{ MHz}, \) which thence shortens the gating time at the expense of just decreasing the performance of the quantum gate by 0.11%. And the parameters used in Fig. 3(c) guarantees a more faster implementation of logic gate at \( t = 2.36 \mu\text{s} \) with a fidelity of 99.31%. Compared with the schemes of Ref. [34, 40], our work does provide a way to perform quantum information processing quickly with...
available experimental parameters.

B. Ground-state blockade

Due to the resemblance between the Förster resonance interaction and our SRP, it is natural to extend the effect of Rydberg blockade to the ground states. Now we introduce a weak interaction \( \Omega_w \) driving the transition between ground states, which can be realized directly by a Raman laser field or microwave field [28]. In this case the effective Hamiltonian of system is

\[
H_{\text{eff}} = 2 \sum_{n=1}^{2} \Omega_w |n\rangle \langle n| + \Omega |1\rangle (|1\rangle + |r\rangle) + H.c.,
\]

where the energy of state \( |1\rangle \) in the ground state space is split by coupling to the single-excited Rydberg state, as shown in Fig. 4. If the Rabi coupling strength \( \Omega \) is much larger than the weak driving strength \( \Omega_w \), the population of state \( |1\rangle \) will be blockaded, and the evolution of atoms is confined to the subspace \( \{|0\rangle, |01\rangle, |10\rangle\} \) governed by \( \dot{\rho}_{gb} = \dot{H}_{gb} = \Omega_w |0\rangle \langle 0| + \langle 0| + |10\rangle + H.c. \). Fig. 5 characterizes the evolutions of all ground states beginning with state \( |0\rangle \). The weak Raman coupling is selected as \( \Omega_w/\pi = 0.001 \text{ MHz} \), and other parameters are the same as in Fig. 2. It clearly shows the phenomenon of ground-state blockade because the population of state \( |1\rangle \) is suppressed well below the order of \( 10^{-4} \) during the whole process.

C. Steady entanglement

A maximally entangled state of two Rydberg atoms in the form of \( \frac{|01\rangle + |10\rangle}{\sqrt{2}} \) can be produced through the ground-state blockade effect, as illustrated by the dashed line of Fig. 5. This unitary-dynamics-based protocol requires to exactly tailor the initial state and precisely control the interaction time. In contrast, the reservoir-engineering approaches to entanglement generation is able to release the above two restrictions by converting the decoherence factor into a resource [12–14, 19, 41–44]. On the basis of the previous model of Fig. 4, we consider the dissipative dynamics of system described by the effective master equation

\[
\dot{\rho} = -i[H_{\text{eff}}, \rho] + \sum_{n=1}^{2} \sum_{k=0}^{L_n} L_n^k \rho L_n^{k\dagger} - \frac{1}{2} \left( L_n^k L_n^{k\dagger} + L_n^{k\dagger} L_n^k \right) + \left( L_n^{k\dagger} \right) \rho L_n^k + \rho L_n^{k\dagger} L_n^k \]

where \( H_{\text{eff}} \) is the Hamiltonian of Eq. (5) and \( \mathcal{L}_n^{(1)} = \sqrt{1/2} |0\rangle \langle 0| \) is the Lindblad operator indicating the atom decay from the excited Rydberg state \( |r\rangle \) into ground states \( |0\rangle \).
we numerically simulate the temporal evolution of atoms using to dissipatively produce the maximally entangled state of two Rydberg states is about 0.2 ms ($\gamma/2\pi \approx 0.01$ MHz) from a fully mixed state, whose convergence time can be inferred to 20 ms according to the above discussion. Fortunately the lifetime of Rydberg state is adjustable utilizing the method of SRP. In this part we take into account both of them in order to give a more convincing result. The lifetimes of Rb $59D_{3/2}$ Rydberg states is about 0.2 ms ($\gamma = 5$ kHz) at $T = 0$ K [45], whose convergence time can be inferred to be 20 ms according to the above discussion. Fortunately the lifetime of Rydberg state is adjustable utilizing the method of engineered spontaneous emission [19, 46]. By coupling the Rydberg state to a short-lived state (decay rate $\Gamma$) with a weak driving field $\Omega_d$, we may obtain an extra effective decay rate for the Rydberg state as $\Delta + \delta/2$. Therefore we find a fast way to dissipatively produce the maximally entangled state of two atoms using $\gamma/2\pi = 0.03$ MHz in Fig. 6, and the dynamics described by the full master equation is consistent with the effective prediction.

and $|1\rangle$ with the same branching ratio $\gamma/2$. A simple inspection shows that the singlet state $|01\rangle = \frac{1}{\sqrt{2}} (|0\rangle - |1\rangle)$ is the unique stationary state solution of Eq. (6). In the inset of Fig. 6, we investigate the influence of different spontaneous emission rates on the convergence time of the maximally entangled state by fixing other parameters as $\Omega_\text{w}/2\pi = 0.005$ MHz, $\Omega_\text{i}/2\pi = 0.01$ MHz from an fully mixed state $\rho_0 = \sum_{i,j=0,1} |ij\rangle\langle ij|/4$. It reveals that a large decay rate can accelerate the convergence time. For example, the choice of $\gamma = 0.1$ MHz ensures a fidelity of 99.39% is achievable at $t = 1$ ms (blue circle), while a selection of $\gamma = 0.001$ MHz delay the time to 100 ms (yellow rhombic). We know the effective lifetime of Rb $59D_{3/2}$ Rydberg states is about 0.2 ms ($\gamma = 5$ kHz) at $T = 0$ K [45], whose convergence time can be inferred to be 20 ms according to the above discussion. Fortunately the lifetime of Rydberg state is adjustable utilizing the method of engineered spontaneous emission [19, 46]. By coupling the Rydberg state to a short-lived state (decay rate $\Gamma$) with a weak driving field $\Omega_d$, we may obtain an extra effective decay rate for the Rydberg state as $\Delta + \delta/2$. Therefore we find a fast way to dissipatively produce the maximally entangled state of two atoms using $\gamma/2\pi = 0.03$ MHz in Fig. 6, and the dynamics described by the full master equation is consistent with the effective prediction.

FIG. 7. The temporal evolution of populations for states $|00\rangle$, $|01\rangle$, and $|10\rangle$ in the presence of Föster defect $\delta/2\pi = 8.5$ MHz, and the other parameters are the same as in Fig. 2.

IV. INFLUENCE OF IMPERFECTION

In the presence of Föster defect, the dipole-dipole coupling between two Rydberg states $|rr\rangle$ and $\langle pp'\rangle$ is modified by $H_{dd} = \frac{1}{\sqrt{2}J} \frac{\delta}{\delta}$, where the $\delta$ is the Föster defect measuring the detuning of the above two states, and its value is only $2\pi \times 8.5$ MHz in the absence of an electric field [38]. This Föster defect alters the eigenvalues of the dipole-dipole interaction and then bring about a deviation $\epsilon = |\Delta| = \sqrt{2}J - (\Delta + \delta + \sqrt{\delta^2 + \delta^2})/2$ of the condition for realization of SRP, which is approximately expanded to $\epsilon \approx \delta/2 + \sqrt{\delta^2}/16J$ for a small ratio $\delta/J$. In Fig. 7 we numerically simulate the temporal evolution of populations for states $|00\rangle$, $|01\rangle$, and $|10\rangle$ in the presence of Föster defect $\delta/2\pi = 8.5$ MHz, which shows that our SRP still works because the quantum states are well suppressed to their initial states, although the corresponding values are slightly lower than that in Fig. 2. In this sense we can claim that our scheme is robust against the fluctuation of Föster defect.

Until now we have not considered the spontaneous emissions of the Rydberg states $|p\rangle$ and $|p'\rangle$, since they are virtually excited in the mechanism of SRP. In this part we take into account both of them in order to give a more convincing result. The lifetimes of Rb $61F_{1/2}$ and $57F_{3/2}$ Rydberg states can be estimated as 0.48 ms and 0.13 ms respectively, from Refs. [45, 47]. By substituting all these factors into the master equation, we plot the corresponding results for controlled-Z gate, ground-state blockade and dissipative entanglement generation in Fig. 8, which are basically consistent with the ideal
conditions.

V. SUMMARY

In summary, we have exhibited how to selectively pump the quantum state of neutral atoms using the current technical means and experimental parameters. This SRP mechanism has an analogous form to the Förster resonance interaction, which is very useful for implementation of controlled-\(Z\) gate, realization of ground-state blockade, and speedup dissipative entanglement generation. Note that all the parameters in our paper are selected as the simplest time-independent form in order to explain our scheme more clearly. The generalization of our scheme to the cases of soft temporal modulation and multipartite interaction will be investigated in our future study. We hope that our work may provides a new prospect with regard to quantum information processing of neutral atoms.

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