Stimulated adiabatic passage in a dissipative ensemble of atoms with strong Rydberg-state interactions

David Petrosyan
Institute of Electronic Structure and Laser, FORTH, GR-71110 Heraklion, Crete, Greece

Klaus Mølmer
Department of Physics and Astronomy, University of Aarhus, DK-8000 Aarhus C, Denmark
(Received 8 February 2013; published 21 March 2013)

We study two-photon excitation of Rydberg states of atoms under stimulated adiabatic passage with delayed laser pulses. We find that the combination of strong interaction between the atoms in Rydberg state and the spontaneous decay of the intermediate exited atomic state leads to the Rydberg excitation of precisely one atom within the atomic ensemble. The quantum Zeno effect offers a lucid interpretation of this result: the Rydberg blocked atoms repetitively scattering photons effectively monitor a randomly excited atom, which therefore remains in the Rydberg state. This system can be used for deterministic creation and, possibly, extraction of Rydberg atoms or ions one at a time. The sympathetic monitoring via decay of ancilla particles may find wider applications for state preparation and probing of interactions in dissipative many-body systems.

DOI: 10.1103/PhysRevA.87.033416 PACS number(s): 32.80.Ee, 32.80.Rm, 32.80.Qk

I. INTRODUCTION

Rydberg atoms strongly interact with each other via long-range dipole-dipole (DD) or van der Waals (vdW) potentials [1]. Within a certain interatomic distance, the interaction-induced level shifts can suppress resonant optical excitation of multiple Rydberg atoms [2–9]. A collection of atoms in the corresponding blockade volume then forms a "superatom," which can accommodate at most one shared Rydberg excitation [10–13].

The Rydberg blockade mechanism constitutes the basis for a number of promising quantum information schemes [2,3,14] and interesting multiatom effects [15–34]. Resonant two-photon excitation of Rydberg states is employed in several schemes [34–46] utilizing the effects of atomic coherence, such as electromagnetically induced transparency [47] and coherent population trapping and transfer [48]. Stimulated adiabatic passage with delayed pulses in an ensemble of three-level atoms was previously considered in Ref. [46], where all the atomic states were assumed to be stable, while the lower atomic transition was driven either by a microwave field or by a pair of optical fields in the Raman configuration. It was shown that, under the Rydberg blockade, the application of the "counterintuitive" pulse sequence results in a multiatom entangled state with strongly correlated population of the two lower states.

The purpose of the present work is to investigate the more typical experimental situation [36–40] in which both transitions of three-level atoms are driven by optical fields in a ladder (Ξ) configuration, while the intermediate excited state of the atoms undergoes rapid spontaneous decay. For noninteracting (distant) atoms, the situation is analogous to what is usually referred to as stimulated Raman adiabatic passage (STIRAP) in a Λ -configuration [48]. Adding interatomic interactions leads to highly nontrivial behavior of the Rydberg superatom. In the earlier part of the process, the system is in a completely symmetric superposition of N atoms, each undergoing adiabatic passage towards the Rydberg state without populating the intermediate excited state. But once any one atom is excited to the Rydberg state, it blocks further Rydberg excitations and triggers the cycling excitation and decay of the intermediate excited state of all the other $N-1$ atoms. This destroys the interatomic coherences and dephases the single Rydberg excitation, which therefore decouples from the field. Through the exact solution of the N-atom master equation, we obtain at the end of the process a mixed state of the system with a single Rydberg excitation incoherently shared among all N atoms. We can also understand the underlying physical mechanism in terms of the quantum Zeno effect [49], in which atoms emitting spontaneous photons through the decay of the intermediate excited state reveal that interactions block their adiabatic passage towards the Rydberg state and thereby perform frequent projective measurements of the presence of a Rydberg excitation in the ensemble. Remarkably, the larger is the number of atoms within the blockade volume the more robust is the transfer process resulting in a single Rydberg excitation of the superatom.

II. ADIABATIC PASSAGE IN A MULTIATOM SYSTEM

A. Single-atom STIRAP

Let us first recall the essence of adiabatic transfer of population in an isolated three-level atom using a pair of delayed laser pulses (STIRAP) [48]. A coherent optical field with Rabi frequency Ω_{ge} resonantly couples the stable ground state $|g\rangle$ to an unstable (decaying) excited state $|e\rangle$, which in turn is resonantly coupled to another stable state $|r\rangle$ by the second coherent field of Rabi frequency Ω_{er} [Fig. 1(a)]. The eigenstates of the corresponding Hamiltonian $\mathcal{V}_{\text{af}} = \hbar(\Omega_{ge}|e\rangle\langle g| + \Omega_{er}|r\rangle\langle e| + \text{H.c.})$ are given by $|\psi_0\rangle = (\cos\theta\,|g\rangle - \sin\theta\,|r\rangle)$ and $|\psi_\pm\rangle = \frac{1}{\sqrt{2}}(\sin\theta\,|g\rangle \pm |e\rangle + \cos\theta\,|r\rangle)$, where the mixing angle θ is defined via $\tan\theta = \Omega_{ge}/\Omega_{er}$. The "dark" state $|\psi_0\rangle$ with energy $\lambda_0 = 0$ does not have any contribution from the fast decaying state $|e\rangle$, while the "bright" states $|\psi_\pm\rangle$ having energies $\lambda_\pm = \pm\hbar\sqrt{\Omega_{ge}^2 + \Omega_{er}^2}$ do contain $|e\rangle$ and thus are unstable against spontaneous decay. The aim of the STIRAP process is to completely transfer the population between the

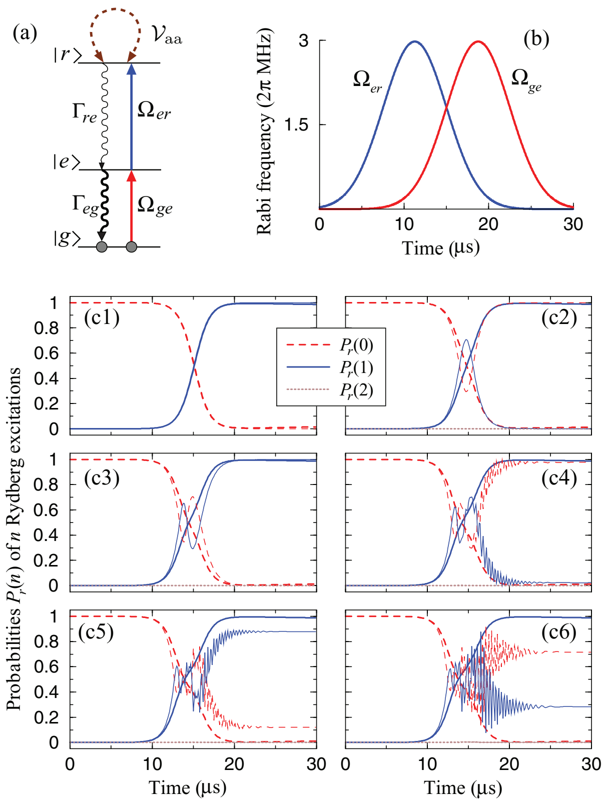

FIG. 1. (Color online) (a) Level scheme of atoms interacting with the fields Ω_{ge} and Ω_{er} on the transitions $|g\rangle \to |e\rangle$ and $|e\rangle \to |r\rangle$, while Γ_{eg} and Γ_{re} are (population) decay rates of states $|e\rangle$ and $|r\rangle$. $\mathcal{V}_{\rm aa}$ denotes the interaction between atoms in Rydberg state $|r\rangle$. (b) Time dependence of the Ω_{ge} and Ω_{er} fields. (c) The corresponding time-dependent probabilities $P_r(n)$ of $n = 0,1,2$ Rydberg excitations in a compact ensemble of $N = 1 - 6$ atoms, (c1)–(c6), obtained from the exact solutions of Eq. (1); thicker lines correspond to the dissipative system (Γ_{eg},Γ_{re} given in the text), while thinner lines to the fully coherent dynamics of nondecaying atoms ($\Gamma_{eg},\Gamma_{re} \simeq 0$).

two stable states $|g\rangle$ and $|r\rangle$ without populating the unstable state $|e\rangle$, which is achieved by adiabatically changing the dark state superposition. With the system initially in state $|g\rangle$, one first applies the Ω_{er} field, resulting in $\langle g|\psi_0\rangle = 1$ ($\Omega_{ge} \ll \Omega_{er}$ and therefore $\theta = 0$). This is then followed by switching on Ω_{ge} and switching off Ω_{er} [Fig. 1(b)], resulting in $|\langle r|\psi_0\rangle| = 1$ ($\Omega_{ge} \gg \Omega_{er}$ and therefore $\theta = \pi/2$). If the mixing angle is rotated slowly enough, $\dot{\theta} \ll \frac{1}{\hbar}|\lambda_\pm - \lambda_0|$, the system adiabatically follows the dark state $|\psi_0\rangle$, and the bright states $|\psi_\pm\rangle$, and thereby $|e\rangle$, are never populated. Hence, the decay of $|e\rangle$ is neutralized and the population of the system is completely transferred from $|g\rangle$ to $|r\rangle$. For what follows, it is useful to remember that the dark state $|\psi_0\rangle$ does not contain $|e\rangle$ because the resonant coupling of $|g\rangle$ to $|e\rangle$ by Ω_{ge} interferes destructively with the resonant coupling of $|r\rangle$ to $|e\rangle$ by Ω_{er}.

B. The *N*-atom master equation

Consider now an ensemble of N three-level atoms confined in a small volume with linear dimension L of several μm. All the atoms uniformly interact with two optical fields of Rabi frequencies Ω_{ge} and Ω_{er} as shown in Fig. 1(a). The atom-field interaction Hamiltonian reads $\mathcal{V}_{\rm af}^j = \hbar(\Omega_{ge}\hat{\sigma}_{eg}^j + \Omega_{er}\hat{\sigma}_{re}^j + \text{H.c.})$, where $\hat{\sigma}_{\mu\nu}^j \equiv |\mu\rangle_{jj}\langle\nu|$ are the transition operators for atom j. The intermediate excited state $|e\rangle$ decays to the ground state $|g\rangle$ with the rate Γ_{eg}; the corresponding Liouvillian acting on the density matrix $\hat{\rho}$ of the system is given by $\mathcal{L}_{eg}^j\hat{\rho} = \frac{1}{2}\Gamma_{eg}[2\hat{\sigma}_{ge}^j\hat{\rho}\hat{\sigma}_{eg}^j - \hat{\sigma}_{ee}^j\hat{\rho} - \hat{\rho}\hat{\sigma}_{ee}^j]$. The decay rate Γ_{re} of the highly excited Rydberg state $|r\rangle$ is typically much smaller (and can be neglected when $\Gamma_{re} \ll \Omega_{er}^2/\Gamma_{eg}$), but for completeness we include it via $\mathcal{L}_{re}^j\hat{\rho} = \frac{1}{2}\Gamma_{re}[2\hat{\sigma}_{er}^j\hat{\rho}\hat{\sigma}_{re}^j - \hat{\sigma}_{rr}^j\hat{\rho} - \hat{\rho}\hat{\sigma}_{rr}^j]$. Note that both transitions of the three-level atoms are assumed closed.

We next include the interatomic interactions. The long-range potential between pairs of atoms i,j in the Rydberg state $|r\rangle$ induces level shifts $\Delta_{ij} = C_p/d_{ij}^p$ of states $|r_ir_j\rangle$, where d_{ij} is the interatomic distance and C_p is the DD ($p = 3$) or vdW ($p = 6$) coefficient. The atom-atom interaction Hamiltonian reads $\mathcal{V}_{\rm aa}^{ij} = \hbar\hat{\sigma}_{rr}^i\Delta_{ij}\hat{\sigma}_{rr}^j$. We assume that all the atoms are within a blockade distance from each other, $\Delta_{ij} \gg \max[w]\ \forall\, i,j \in [1,N]$, where $w = \frac{\Omega_{ge}^2+\Omega_{er}^2}{\sqrt{2\Omega_{ge}^2+\Gamma_{eg}^2/4}}$ is the Rydberg-state excitation linewidth of a single three-level atom. We define the probabilities $P_r(n) = \langle\hat{\Sigma}_r^{(n)}\rangle$ of n Rydberg excitations of superatom through the corresponding projectors $\hat{\Sigma}_r^{(0)} \equiv \prod_{i=1}^N(\hat{\sigma}_{gg}^i + \hat{\sigma}_{ee}^i) = \prod_{i=1}^N(\mathbb{1} - \hat{\sigma}_{rr}^i)$, $\hat{\Sigma}_r^{(1)} \equiv \sum_{j=1}^N \hat{\sigma}_{rr}^j \prod_{i\neq j}^N(\mathbb{1} - \hat{\sigma}_{rr}^i)$, etc. Note that $\hat{\sigma}_{gg}^i + \hat{\sigma}_{ee}^i + \hat{\sigma}_{rr}^i = \mathbb{1}\ \forall\, i \in [1,N]$.

The density operator $\hat{\rho}$ of the N-atom system obeys the master equation [50]

$$\partial_t\hat{\rho} = -\frac{i}{\hbar}[\mathcal{H},\hat{\rho}] + \mathcal{L}\hat{\rho}, \tag{1}$$

with the Hamiltonian $\mathcal{H} = \sum_j \mathcal{V}_{\rm af}^j + \sum_{i<j}\mathcal{V}_{\rm aa}^{ij}$ and the Liouvillian $\mathcal{L}\hat{\rho} = \sum_j(\mathcal{L}_{eg}^j\hat{\rho} + \mathcal{L}_{re}^j\hat{\rho})$.

III. NUMERICAL SIMULATIONS

We solve the master equation (1) numerically assuming the atoms are irradiated by the two pulsed fields having Gaussian temporal shapes

$$\Omega_{ge,er}(t) = \Omega_0 \exp\left[-\frac{\left(t - \frac{1}{2}t_{\rm end} \mp \sigma_t\right)^2}{2\sigma_t^2}\right],$$

where $\Omega_0 = 2\pi \times 3$ MHz is the peak amplitude, $2\sigma_t = \frac{1}{4}t_{\rm end}$ is the temporal width and relative delay of the pulses, and $t_{\rm end} = 30\ \mu s$ is the process duration [see Fig. 1(b)]. We take cold ^{87}Rb atoms [36–40], with the ground state $|g\rangle \equiv 5S_{1/2}\,|F = 2,m_F = 2\rangle$, the intermediate excited state $|e\rangle \equiv 5P_{3/2}\,|F = 3,m_F = 3\rangle$ with $\Gamma_{eg} = 38$ MHz, and the highly excited Rydberg state $|r\rangle \equiv nS_{1/2}$ with the principal quantum number $n \sim 80$ and $\Gamma_{re} = 1$ kHz. Within the trapping volume of linear dimension $L \sim 5\ \mu$m we then have large interatomic (vdW) interactions [51] $\Delta_{ij} \geqslant 10w_0\ \forall\, d_{ij} \leqslant L$, where $w_0 = \frac{2\Omega_0^2}{\sqrt{2\Omega_0^2+\Gamma_{eg}^2/4}} \simeq 2\pi \times 3.5$ MHz.

The results of simulations for $N = 1,\ldots,6$ atoms are summarized in Fig. 1(c). For any N, even or odd, the "counterintuitive" sequence of pulses $\Omega_{ge,er}(t)$ leads, with large probability $P_r(1) \geqslant 0.98$, to a single Rydberg excitation of the superatom, while the probabilities of multiple excitations $P_r(n>1)$ are negligible, due to the strong blockade. Once a Rydberg excitation is produced, the small decay Γ_{re} of state $|r\rangle$ leads to a slow decrease of $P_r(1)$.

The response of the Rydberg superatom to the "counterintuitive" sequence of pulses may look analogous to the coherent adiabatic passage of a single three-level atom, but this similarity is superficial and the physics behind it is more involved. This is perhaps best illustrated in Fig. 1(c) by the strikingly different behavior of superatom in the absence of dissipation, $\Gamma_{eg},\Gamma_{re}=0$, which was studied in Ref. [46]. Without dissipation, in the transition region $\Omega_{ge}(t)\sim\Omega_{er}(t)$ the probabilities of zero $P_r(0)$ and one $P_r(1)$ Rydberg excitation do not change monotonically but alternate $N-1$ times, with the result that for an even number of atoms N the final state of the system does not contain a Rydberg excitation. (For $N\geqslant 4$, the fast oscillations of probabilities $P_r(0,1)$ and their final values noticeably different from 0 and 1 are due to the violation of adiabaticity with the increased system size and the corresponding decrease in the separation between its eigenstates [46]).

A. Analysis

For a dissipationless system, it is convenient to use the fully symmetrized states $|n_g,n_e,n_r\rangle$ denoting n_g atoms in state $|g\rangle$, n_e atoms in $|e\rangle$, and n_r atoms in $|r\rangle$. Due to the Rydberg blockade, only $n_r=0,1$ values are allowed, while $n_g+n_e+n_r=N$. The field Ω_{ge} couples the ground state of the superatom $|N_g,0_e,0_r\rangle$ successively to the collective single $|(N-1)_g,1_e,0_r\rangle$, double $|(N-2)_g,2_e,0_r\rangle$, etc. excitation states, which are in turn coupled to the single Rydberg excitation states $|(N-1)_g,0_e,1_r\rangle$, $|(N-2)_g,1_e,1_r\rangle$, etc. by the field Ω_{er} [43]. The corresponding Hamiltonian can be expressed as $\mathcal{H}=\hbar(\Omega_{ge}\hat{e}^\dagger\hat{g}+\Omega_{er}\hat{r}^\dagger\hat{e}+\text{H.c.})$, where operators $\hat{g}$ ($\hat{g}^\dagger$), $\hat{e}$ ($\hat{e}^\dagger$), and $\hat{r}$ ($\hat{r}^\dagger$) annihilate (create) an atom in the corresponding state $|g\rangle$, $|e\rangle$, and $|r\rangle$; $\hat{g}$ and $\hat{e}$ are standard bosonic operators, while $\hat{r}$ describes a hard-core boson $(\hat{r}^\dagger)^2=0$. As was shown in Ref. [46] for nondecaying atoms, an ideal adiabatic passage leads to the final state of the system $|J_x=0\rangle$ for N even, and $|J_x=0\rangle\otimes|1_r\rangle$ for N odd, where $|J_x=0\rangle$ is the eigenstate of operator $\hat{J}_x\equiv\frac{1}{2}(\hat{e}^\dagger\hat{g}+\hat{g}^\dagger\hat{e})$ with zero eigenvalue. The state $|J_x=0\rangle$ involves an equal number of atoms [$N/2$ or $(N-1)/2$] in $(|g\rangle\pm|e\rangle)/\sqrt{2}$.

In the presence of strong decay $\Gamma_{eg}\gtrsim\Omega_{ge}$ of state $|e\rangle$, the dynamic of the system is completely different. If we had noninteracting atoms, the adiabatic passage would yield a product state $(\cos\theta\,|g\rangle-\sin\theta\,|r\rangle)^{\otimes N}$ containing multiple Rydberg excitations but no atoms in state $|e\rangle$. The strong interatomic interactions, however, shift the energies of multiply excited Rydberg states out of resonance with the Ω_{er} field. Starting from the ground state $|N_g,0_e,0_r\rangle=\prod_{i=1}^N|g\rangle_i$ the population transfer beyond the symmetric single Rydberg excitation state $|(N-1)_g,0_e,1_r\rangle=\frac{1}{\sqrt{N}}\sum_{j=1}^N|r\rangle_j\prod_{i\neq j}^N|g\rangle_i$ is then blocked. In this superposition, the atoms in state $|g\rangle$ can be now excited to state $|e\rangle$ by the strong resonant field Ω_{ge} since the coupling of $|e\rangle$ to $|r\rangle$ by Ω_{er} and the resulting destructive interference are suppressed. Atoms excited to $|e\rangle$ rapidly decay back to $|g\rangle$ with random phases. This leads to continuous dephasing of the superposition $|(N-1)_g,0_e,1_r\rangle$, turning it into an incoherent mixture of single Rydberg excitation, which is decoupled from states $|(N-n)_g,n_e,0_r\rangle$ and $|N_g,0_e,0_r\rangle$. A related effect is described in Ref. [12], where the dephasing of collective Rydberg excitation was brought about by an inhomogeneous light field with short-range space and time correlations. Here, instead, the superposition containing a Rydberg atom is dephased by the field Ω_{ge} and decay Γ_{eg} through the cycling transition $|g\rangle\leftrightarrow|e\rangle$ of all the atoms that do not populate the Rydberg state.

The field $\Omega_{ge}(t)$ driving the $N-1$ blocked (two-level) atoms is spatially uniform but varies slowly in time. At any time t, the average populations of state $|g\rangle$ for all the atoms is therefore approximately given by the steady-state expression for an independent two-level atom [50], $\langle\hat{\sigma}_{gg}\rangle\approx\frac{\Omega_{ge}^2+\Gamma_{eg}^2/4}{2\Omega_{ge}^2+\Gamma_{eg}^2/4}\equiv\kappa$. The probability that all but the Rydberg excited atom are in the ground state $|g\rangle$ is then $\langle\hat{\sigma}_{rr}^j\prod_{i\neq j}^N\hat{\sigma}_{gg}^i\rangle\approx[\kappa(t)]^{N-1}$. At large times $\kappa(t_{\text{end}})\simeq 1$ and the above probability approaches unity, with the system in the mixed state $\hat{\rho}=\frac{1}{N}\sum_{j=1}^N\hat{\sigma}_{rr}^j\prod_{i\neq j}^N\hat{\sigma}_{gg}^i$. Yet, the total probability of finding one Rydberg excitation is $P_r(1)=\langle\hat{\Sigma}_r^{(1)}\rangle\simeq 1$ already when $\Omega_{ge}>\Omega_{er}$ [Fig. 1]. These results are fully reproduced by the exact numerical solution of the density matrix equations (1).

B. Quantum Zeno effect

An alternative and perhaps more elegant explanation as to why the superatom attains near unity Rydberg excitation $\langle\hat{\Sigma}_r^{(1)}\rangle\simeq 1$ is based on the quantum Zeno effect [49]: Upon repetitive excitation to state $|e\rangle$ and spontaneous decay back to the ground state $|g\rangle$, the Rydberg blocked atoms perform continuous projective measurements of the Rydberg excitation in the ensemble.

To verify this physical picture, we have performed quantum Monte Carlo simulations [52] of the dissipative dynamics of a few-atom system. In such simulation, the state of the system $|\Psi\rangle$ evolves according to the Schrödinger equation $\partial_t|\Psi\rangle=-\frac{i}{\hbar}\tilde{\mathcal{H}}|\Psi\rangle$ with an effective non-Hermitian Hamiltonian $\tilde{\mathcal{H}}=\mathcal{H}-i\hbar\sum_j\frac{1}{2}(\Gamma_{eg}\hat{\sigma}_{ee}^j+\Gamma_{re}\hat{\sigma}_{rr}^j)$, which does not preserve the norm of $|\Psi\rangle$. The evolution is interrupted by random quantum jumps $|\Psi\rangle\to\hat{\sigma}_{ge(er)}^j|\Psi\rangle$ with probabilities determined by the decay rates $\Gamma_{eg(re)}$.

In the early part of evolution, the states of all the atoms share the same overlap with the current dark and bright states, and the state of the system $|\Psi\rangle$ is symmetric under permutation of the atoms. But already the first quantum jump breaks this symmetry by transferring one randomly selected atom from $|e\rangle$ to $|g\rangle$ (or, with a much smaller probability $\sim\Gamma_{re}/\Gamma_{eg}$, from $|r\rangle$ to $|e\rangle$), while the dark state overlap increases for the atoms that did not jump. During the subsequent evolution under the Rydberg blockade, the bright state overlap grows for all the atoms, while the following jump again suddenly increases the dark state contribution to the states of the atoms which did not jump. This proceeds until eventually only one atom has experienced no quantum jump. This atom closely follows the dark state superposition, and while small excursions away from the dark state occur, they are reduced by each jump of the other atoms. Hence, atoms undergoing quantum jumps stabilize the almost deterministic evolution of a non-decaying atom towards the Rydberg state.

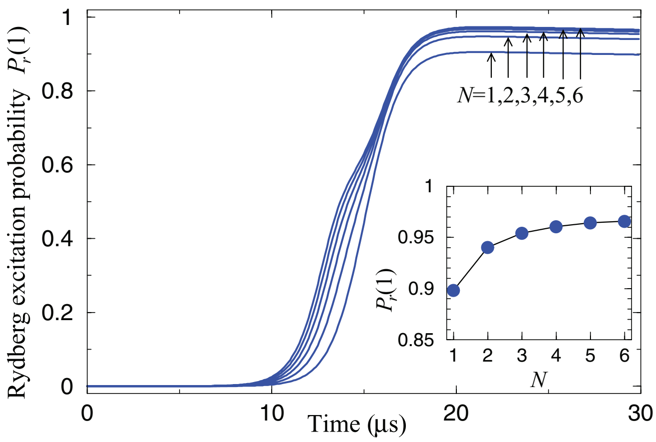

FIG. 2. (Color online) Probabilities $P_r(1)$ of single Rydberg excitation of superatom composed of $N = 1 - 6$ atoms, obtained from the exact solutions of Eq. (1) including dephasing γ_r and decays Γ_{eg}, Γ_{re}. Main panel shows the time dependence of $P_r(1)$, and inset shows $P_r(1)$ for different N at time $t_{\rm end} = 30\ \mu$s.

C. Coherence relaxation

We finally demonstrate the robustness of adiabatic passage in a Rydberg superatom. In a single three-level atom, the STIRAP—while immune to the decay and moderate detuning of the intermediate state $|e\rangle$—is very sensitive to coherence relaxation between the long-lived states $|g\rangle$ and $|r\rangle$ [48]. The physical origins of relaxation of the atomic coherence include nonradiative collisions, Doppler shifts, laser phase fluctuations and electromagnetic field noise. In addition to the small decay Γ_{re} of the Rydberg state $|r\rangle$, we now include its dephasing γ_r via the Liouvillian $\mathcal{L}_r^j \hat{\rho} = \frac{1}{2}\gamma_r[(\hat{\sigma}_{rr}^j - \hat{\sigma}_{ee}^j - \hat{\sigma}_{gg}^j)\hat{\rho}(\hat{\sigma}_{rr}^j - \hat{\sigma}_{ee}^j - \hat{\sigma}_{gg}^j) - \hat{\rho}]$ [50]. In Fig. 2 we show the probability of Rydberg excitation of the superatom obtained from the solution of the master equation (1) with the dephasing rate $\gamma_r = 2\pi \times 0.1$ MHz. For a single atom, the excitation probability is now reduced to $\langle\hat{\sigma}_{rr}\rangle \simeq 0.9$, since the decoherence of the dark state superposition leads to the population of the bright states [47]. With increasing the number of atoms N, however, the Rydberg excitation probability of the superatom grows according to $\langle\hat{\Sigma}_r^{(1)}\rangle \approx \frac{N\langle\hat{\sigma}_{rr}\rangle}{(N-1)\langle\hat{\sigma}_{rr}\rangle+1}$ [33], approaching $P_r(1) \gtrsim 0.97$ for $N = 6$. The spontaneous decay, perhaps surprisingly, counteracts the detrimental effect of decoherence of the dark state superposition and facilitates the efficient production of a single Rydberg excitation. A similar result has been obtained for Rydberg superatoms composed of incoherently driven two-level atoms [33].

IV. CONCLUSION

To conclude, we have examined the excitation of a Rydberg superatom using adiabatic passage with delayed laser pulses. We have found that spontaneous decay of atoms from the intermediate excited state facilitates a single Rydberg excitation of the superatom, with nearly unit probability. An ensemble of $N > 1$ atoms in a tight trap of linear dimension smaller than the Rydberg blockade distance can repeatedly and reliably produce single Rydberg atoms [53] or ions [40], which can then be extracted and possibly deposited elsewhere. This process can proceed down to $N = 1$ atoms, and the final count of the extracted Rydberg atoms would correspond to the initial number of ground-state atoms in the trap.

Experiments with the atomic ensembles much larger than the blockade length have revealed significant suppression of the number of Rydberg excitations [4–7], which is consistent with a regular spatial arrangement of Rydberg atoms [25,32]. We envisage that employing stimulated adiabatic passage to produce with unit probability singe Rydberg atoms per blockade volume can result in long-range order and tighter crystallization of Rydberg excitations in extended systems [33,34] and may shine more light onto the spatial correlation patterns of Rydberg excitations.

Finally, our studies provide an interesting element to the very active field of dissipative generation of states and processes in open many-body systems [54,55]. In our many-particle system, dissipation assigns different roles to different particles, such that some particles become dissipative probes for the coherent dynamics of the others. The observation that the decay of the probe particles may counteract decoherence in the target particle may find application in, e.g., ancilla-assisted protocols for quantum computing [56].

ACKNOWLEDGMENTS

K.M. gratefully acknowledges support from the FET-Open grant MALICIA (265522).

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