Deterministic Generation of Arbitrary Photonic States Assisted by Dissipation

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(Received 4 May 2015; published 16 October 2015)

A scheme to utilize atomlike emitters coupled to nanophotonic waveguides is proposed for the generation of many-body entangled states and for the reversible mapping of these states of matter to photonic states of an optical pulse in the waveguide. Our protocol makes use of decoherence-free subspaces (DFSs) for the atomic emitters with coherent evolution within the DFSs enforced by strong dissipative coupling to the waveguide. By switching from subradiant to superradiant states, entangled atomic states are mapped to photonic states with high fidelity. An implementation using ultracold atoms coupled to a photonic crystal waveguide is discussed.

DOI: 10.1103/PhysRevLett.115.163603 PACS numbers: 42.50.-p, 03.67.Bg, 42.50.Ex

Recent work on optical emitters coupled to one-dimensional (1D) waveguides has opened new avenues to investigate light-matter interactions [1–20]. Particularly promising are the setups where atoms are strongly coupled to structured dielectrics [6–10], where large Purcell factors have been predicted [21,22]. Furthermore, collective effects can be enhanced by placing the atoms at particular positions [15,16,23–29]. The combination of atomlike emitters and nanophotonic waveguides may enable new regimes for the interaction of light and matter, leading to technologies that outperform current ones and qualitatively different physics. In this work we investigate the possibility of using atom nanophotonics interfaces to tailor arbitrary states for propagating photons on demand, which lies at the heart of many quantum information [30], metrology [31], and lithography [32] methods (see Ref. [33] for a review). We predict large fidelities even for relatively large numbers of photons, something which has been impossible to achieve with other platforms in the optical domain.

Our proposal uses $N+1$ three-level systems (with levels $\{|g\rangle, |s\rangle, |e\rangle\}$), where one of the optical transitions ($|g\rangle \leftrightarrow |e\rangle$) is strongly coupled to a 1D waveguide [see Figs. 1(a) and 1(b)]. We denote by P_{1D} the Purcell factor corresponding to that transition, i.e., the ratio of the emission rate into the waveguide mode, Γ_{1D}, and the one for all other modes, Γ^*. The atoms must be separated by distances proportional to $\lambda_a = 2\pi/q(\omega_a)$, where $q(\omega)$ is the wave number determined by the waveguide dispersion relation. Depending on their internal state, atoms may experience a collective decay into the waveguide, or become completely decoupled from it. The latter occurs if they are in a decoherence free subspace (DFS) [34–36]. Our protocol consists of two steps: in the first one, we generate certain states within the DFS, $|\Psi_D\rangle$, by driving the atoms with lasers and using the collective quantum Zeno effect [37–39] within the DFS with an infidelity $1-F_1 \propto m/\sqrt{P_{1D}}$, where m is the maximum number of photons we want to generate; in the second one, a laser pulse takes the atomic state out of the DFS so that atoms collectively emit into the waveguide, creating the desired state of a single propagating mode, $|\Psi_B\rangle$, with an infidelity $1-F_2 \propto m^2/(NP_{1D})$.

The atom-photon Hamiltonian of these systems is given by $H = H_{qb} + H_{field} + H_I$, with $H_{qb} = \sum_{n=1}^{N+1}(\omega_a \sigma_{ee}^n + \omega_s \sigma_{ss}^n)$ and $H_{field} = \sum_q \omega_q a_q^\dagger a_q$, (using $\hbar = 1$), where ω_a is the two-level system energy, ω_q is the field dispersion relation of the 1D photonic modes, and $\sigma_{ij}^n = |i\rangle_n\langle j|$. We consider only the coupling to a single polarization as justified for suitable dielectric waveguide modes [21], that is,

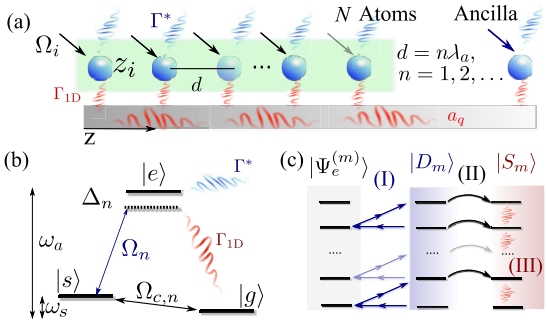

FIG. 1 (color online). (a) Setup: N atoms plus 1 ancilla atom coupled to a 1D photonic bath. The ancilla must be individually addressed. (b) Atomic Λ scheme: the transition $|g\rangle \leftrightarrow |e\rangle$ is coupled to the a_q modes. A laser controls the transition $|s\rangle \leftrightarrow |e\rangle$ with amplitude, Ω_n and detuning Δ_n. Another field, Ω_c, controls $|s\rangle \leftrightarrow |g\rangle$, with frequency ω_s. (c) Relevant states and steps for our protocol: (I) generation of superpositions of symmetric Dicke states, $|D_m\rangle$, by using the excited dark states $|\Psi_e^{(m)}\rangle$. (II) Flipping $|s\rangle \rightarrow |e\rangle$ to generate the superradiant state $|S_m\rangle$, which decays rapidly (III) to the desired photonic state.

$$H_I = \sum_{n,q} (g_q \sigma_{ge}^n a_q^\dagger e^{-iqz_n} + \text{H.c.}), \quad (1)$$

with g_q the single-photon coupling constant to the mode of interest and where we have used the rotating wave approximation. When the 1D baths have a much faster relaxation time scale than the atomic system, the atoms are described by a density matrix, ρ, which in the Born-Markov limit, is governed by a master equation [23,24,40] of the form: $d\rho/dt = \sum_{n,m}[(\Gamma_{1D}/2)e^{iq(\omega_a)|z_n - z_m|}(\sigma_{ge}^n \rho \sigma_{eg}^m - \rho \sigma_{eg}^m \sigma_{ge}^n) + \text{H.c.}]$ in the interaction picture rotating with H_{qb}. By appropriately choosing the atomic positions, e.g., $z_n = n\lambda_a = n2\pi/q(\omega_a)$, with $n \in \mathbb{N}$, the coherent atom-atom interactions are eliminated [41] and the effective interaction yields a pure Dicke model [42] described by

$$\mathcal{L}_D(\rho) = \frac{\Gamma_{1D}}{2}(S_{ge}\rho S_{eg} - S_{eg}S_{ge}\rho) + \text{H.c.}, \quad (2)$$

where we defined $S_{ij} = \sum_{n=1}^{N+1} \sigma_{ij}^n$. One of the assets of the Dicke model is the emergence of subradiant and superradiant states. The excited states with m atoms in $|e\rangle$ that are symmetric under permutations, denoted by $|S_m\rangle$, are superradiant with a decay rate proportional (at least) to the atom number N, and are unique for each m. On the other hand, the states satisfying $S_{ge}|\Psi\rangle = 0$ are dark states of the Liouvillian of Eq. (2), and therefore decoupled from collective dissipation. These dark states span the DFS that is highly degenerate for $m > 1$.

The atomic states that must be created in the first step of our protocol are very peculiar as (i) they must be prepared within the DFS to avoid dissipation, and (ii) they must be easily mapped to the appropriate superradiant states, to generate arbitrary superpositions of the photonic waveguide states. Our strategy consists of first, identifying states, denoted by $|D_m\rangle$, in the subspace spanned by the ground levels g, and s, which can be mapped one-to-one to a basis $|S_m\rangle$ of superradiant states using a simple laser pulse, and which in turn give rise to m photons in the waveguide via superradiance. Then, we use a more sophisticated scheme within the DFS to generate superpositions of $|D_m\rangle$, which requires only m steps.

By introducing another metastable state, $|s\rangle$, as depicted in Fig. 1(b), the candidates to map to superradiant states are the symmetric Dicke states $|D_m\rangle \propto \text{sym}\{|s\rangle^{\otimes m} \otimes |g\rangle^{\otimes N-m}\}$, as they can be turned superradiant by switching $|s\rangle \to |e\rangle$. Having identified the target, $|D_m\rangle$, we need to find ways to build efficiently arbitrary superpositions. Previous studies have proposed implementing one or two-qubit universal gates within DFS [34–38], but the number of steps increases rapidly with N, as well as using adiabatic passage methods [43–46], limited to small excitations number m. Here, we use the collective character of the interaction to deterministically generate arbitrary N-qubit states for which the number of steps is independent of the number of atoms.

The scheme that we use is depicted in Fig. 1(a): we consider a system of $N+1$ emitters, in which we aim to generate $|D_m\rangle$ in the first N emitters using the ancilla as an auxiliary state. As the $|D_m\rangle$'s are invariant under the permutation of the first N atoms, we choose the control fields with the same symmetry and the same detuning $\Delta_n = \Delta_e$ for all emitters:

$$H_c = \frac{\Omega_c}{2}\sigma_{sg}^{N+1} + \text{H.c.},$$
$$H_{\text{las}} = \left(\frac{\Omega_r}{2}\sum_{n=1}^{N}\sigma_{se}^n + \frac{\Omega_{\text{anc}}}{2}\sigma_{se}^{N+1} + \text{H.c.}\right) + \Delta_e \sum_{n=1}^{N+1}\sigma_{ee}^n, \quad (3)$$

written in the interaction picture with respect to H_{qb}. The Hamiltonian H_{las} allows us to control both the emitter state and the coupling to the 1D reservoir, while H_c allows us to control the atomic states of the ancilla independently of the coupling to the 1D reservoir. We are interested in working in the regime of strong collective dissipation, where $N\Gamma_{1D} \gg \Omega_r, \Omega_{\text{anc}}, \Omega_c, \Delta_e$. In this situation, the 1D bath is continuously monitoring the collective atomic state, as in the quantum Zeno regime [34–38], and projecting its state into the DFS of $\mathcal{L}_D$. Formally, we obtain the effective dynamics within the DFS by treating the control fields and dissipation into other decay channels as a perturbation to the collective dissipation $\mathcal{L}_D$ [47]. To explain the protocol, we first consider the action of the control fields projected into the DFS to first order, leaving the discussion of errors due to higher orders for later.

Because of the symmetry of the problem, it is convenient to introduce the following notation to describe any symmetric state over N atoms:

$$|F_{m,k}\rangle \propto \text{sym}\{|s\rangle^{\otimes m} \otimes |e\rangle^{\otimes k} \otimes |g\rangle^{\otimes N-m-k}\}, \quad (4)$$

that embeds both $|D_m\rangle \equiv |F_{m,0}\rangle$ and $|S_m\rangle \equiv |F_{0,m}\rangle$. In general, the DFS of the Liouvillian $\mathcal{L}_D$ grows exponentially with the number of atoms, but for each m, only three of these states fulfill the permutation symmetry of our system [47]. These states are

$$|\Psi_s^{(m)}\rangle = |F_{m-1,0}\rangle \otimes |s\rangle_A,$$
$$|\Psi_g^{(m)}\rangle = |F_{m,0}\rangle \otimes |g\rangle_A,$$
$$|\Psi_e^{(m)}\rangle = \sqrt{\frac{N_m}{N_m+1}}|F_{m-1,0}\rangle \otimes |e\rangle_A - \frac{1}{\sqrt{N_m+1}}|F_{m-1,1}\rangle \otimes |g\rangle_A, \quad (5)$$

where $|\psi\rangle_A$ denotes the state of the ancilla and $N_m = N - m + 1$. In Fig. 2(a) we sketch the protocol steps by a diagram of the projected Hamiltonian into the DFS. It consists of two parts, which are applied $m_{\max}$ times to reach any superposition of states over the first N atoms $|\Psi_D\rangle \otimes |g\rangle_A = (\sum_{m=0}^{m_{\max}} d_m |D_m\rangle) \otimes |g\rangle_A$ with maximum $m_{\max}$ excitations from the initial state $|\Psi_g^{(0)}\rangle$: (i) Use H_c to flip the ancilla state $|g\rangle_A \to |s\rangle_A$. This transition makes:

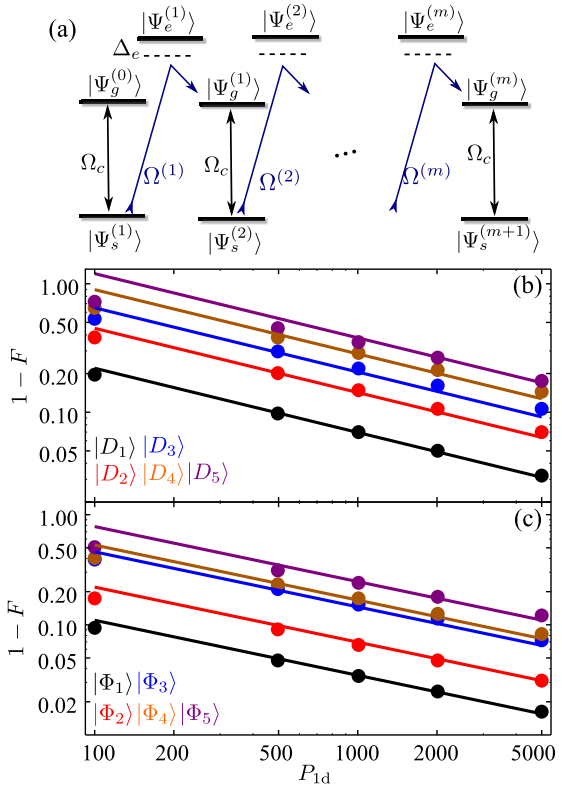

FIG. 2 (color online). (a) Preparation of arbitrary superpositions of $|D_m\rangle$. Alternating between σ_x gates on the ancilla and two-photon transitions via $|\Psi_e^{(m)}\rangle$ builds up excitations step by step. (b) [(c)] Error, $1-F$, as a function of P_{1D} for generating $|D_m\rangle$ $[|\Phi_m\rangle = (1/\sqrt{2})(|D_0\rangle + |D_m\rangle)]$ up to $m_{\max} = 5$ excitations. The dots correspond to the numerical fidelities, whereas the solid lines depict the $1/\sqrt{P_{1D}}$ scaling.

$|\Psi_g^{(m-1)}\rangle \to |\Psi_s^{(m)}\rangle$; (ii) two-photon transition $|\Psi_s^{(m)}\rangle \to |\Psi_g^{(m)}\rangle$. It can be shown that the dark states corresponding to a given excitation m form an effective Λ scheme within the DFS via the far-detuned state $|\Psi_e^{(m)}\rangle$ [47]. The two-photon transition can be made on-resonance if the intensities are chosen such that $|\Omega_{\text{anc}}| = |\Omega_r|\sqrt{m/N_m}$, which is possible because we demanded individual addressing of the ancilla. If this condition is not imposed, the $|\Psi_{s,g}^{(m)}\rangle$ experience different Stark shifts that spoil the two-photon process. The effective Hamiltonian for the mth excitation is

$$H_D = \Omega^{(m)}/2|\Psi_s^{(m)}\rangle\langle\Psi_g^{(m)}| + \text{H.c.}, \tag{6}$$

where $|\Omega^{(m)}| = (|\Omega_r|^2/(2|\Delta_e|))[m/(N_m+1)]$. The combination of both steps gives rise to a ladderlike structure [see Fig. 2(a)], which can be used to build any arbitrary superposition states of $|D_m\rangle$ from the ground state $|\Psi_g^{(0)}\rangle$ with $m_{\max}$ steps each. The necessary pulse sequence can be obtained by calculating the inverse evolution from the target to the initial state and successively removing excitations in the ladderlike structure [39].

Another advantage of our protocol is that it can be used to generate a superposition of photonic states in the 1D bath by dissipative means once we have $|\Psi_D\rangle$. To make sure the ancilla atom can be neglected, we flip the ancilla state $|g\rangle_A \to |s\rangle_A$ and apply no fields to it. In order to map to the superradiant state of N atoms, we apply a fast resonant π pulse ($\Delta_e = 0$ and $\Omega_r \gg N\Gamma_{1D}$) on the N emitters to switch all $|s\rangle_n \to |e\rangle_n$, thus generating the superposition of $|S_m\rangle \propto \text{sym}\{|e\rangle^{\otimes m} \otimes |g\rangle^{\otimes N-m}\}$. Because of their superradiant character, the $|S_m\rangle$ decay completely to 1D-reservoir modes. Because H_I conserves the number of excitations, the superradiant state of m excitations decays to the Fock-state of m photons [47]:

$$|S_m\rangle \to |m_{\{q\}}\rangle \equiv \sum_{\{q\}} \frac{A_{\{q\}}(t)}{m!} a_{q_1}^\dagger \ldots a_{q_m}^\dagger |\text{vac}\rangle, \tag{7}$$

where $\{q\} = \{q_1, \ldots, q_m\}$ is the set of relevant momenta which run over the whole Brillouin Zone $q_i \in$ B.Z. The scattering amplitude $A_{\{q\}}(t)$ is calculated using a generalized input-output formalism [47,52–54] and quantum regression theorem [55]:

$$A_{\{q\}}(t) = \prod_{r=1}^{m} \frac{\mathbf{i}g\sqrt{rN_r}e^{-\mathbf{i}\omega_{q_r}t}}{\mathbf{i}(\sum_{l=1}^{r}\omega_{q_l} - r\omega_a) + r\Gamma_{1D}N_r/2} + \text{P}[\{q\}] \tag{8}$$

for sufficiently large times $t \gg 1/N_m\Gamma_{1D}$ when the atomic state has decayed completely and defining P$[\{q\}]$ as all the permutations of $\{q\}$. The only dependence on t enters through $e^{-\mathbf{i}\sum_{r=1}^{m}\omega_{q_r}t}$, which describes the center-of-mass motion of the wave packet when going to the real space. In the low excitation regime, one can either use the Holstein-Primakoff approximation [56] or directly substitute $N_m \to N$ in the expression above, arriving at

$$A_{\{q\}}^{\text{HP}}(t) = \prod_{r=1}^{m} \frac{e^{-\mathbf{i}\omega_{q_r}t}\mathbf{i}g\sqrt{rN}}{\mathbf{i}(\omega_{q_r} - \omega_a) + \Gamma_{1D}N/2}, \tag{9}$$

that has a Lorentzian shape centered at ω_a with bandwidth $\Gamma_{1D}N/2$. Substituting $A_{\{q\}}(t) \to A_{\{q\}}^{\text{HP}}(t)$ into the definition of $|m_{\{q\}}\rangle$, yields a linear Fock state denoted by $|m_{\{q\}}^{\text{HP}}\rangle$. In principle, the emission into the waveguide is bidirectional ($\pm q$), but combining both fields in phase, e.g., by placing a mirror at an appropriate distance, or by engineering the atom-photon coupling appropriately [3,4,14], it is possible to achieve emission in one direction only. Furthermore, by shaping the pulse, $\Omega_r(t)$ (within a bandwidth $\lesssim N\Gamma_{1D}$) we generate any desired shape of the output photonic state, e.g., to create a time-symmetric photonic state [57] that ensures the reversibility of the process when mapping the photonic state to another sample. Moreover, because of the linearity of the calculation of $A_{\{q\}}$ with respect to the input state [47], superpositions of atomic states decay to superpositions

$$|\Psi_D\rangle \to \sum_{m=0}^{m_{\max}} d_m |S_m\rangle \to |\Psi_B\rangle = \sum_{m=0}^{m_{\max}} d_m |m_{\{q\}}\rangle, \quad (10)$$

that will be generated in a single-mode wave packet, as required for most applications [33], as long as $N \gg m$ because $\langle m_{\{q\}}|m^{\mathrm{HP}}_{\{q\}}\rangle \approx 1 - m^3/(20N^2)$ [47].

So far, we have only discussed the ideal protocol without considering, e.g., spontaneous emission into all other modes with rate Γ^*. For the error in the preparation of the many-body entangled state, we derive an error rate ϵ from perturbation theory, which, together with the time of the operation τ, gives an approximation of the fidelity, $F_1 = \sqrt{\langle\Psi_D|\rho(\tau)|\Psi_D\rangle} \approx 1 - \tau\epsilon$, with respect to the target state $|\Psi_D\rangle$. The dominant errors assuming $N \gg m$ and $\Gamma_{1D} \gg \Delta_e \gg \Gamma^*$ come from [47] (i) the spontaneously emitted photons from $|\Psi_e^{(m)}\rangle$ to decay channels other than the waveguide, which scales as $\epsilon_1 \approx \Gamma^* m|\Omega_r|^2/(4N\Delta_e^2)$, and (ii) from the photons emitted from the small populations of superradiant states. We estimate the rate of these errors by taking into account the second order corrections of the projected Hamiltonian which are finally given by $\epsilon_2 \approx N\Gamma_{1D}\{m|\Omega_r|^2/[4\Delta_e^2 + (N\Gamma_{1D})^2]\}$. Summing up, the error for the mth step of the process, which takes $\tau = (\pi/|\Omega^{(m)}|) \approx 2\pi N\Delta_e/(m|\Omega_r|^2)$ for full population transfer, is

$$1 - F_1^{(m)} \approx \frac{\pi}{2}\left(\frac{\Gamma^*}{\Delta_e} + \frac{\Delta_e}{\Gamma_{1D}}\right), \quad (11)$$

that is optimized for $\Delta_{e,\mathrm{opt}} = \sqrt{\Gamma^*\Gamma_{1D}}$, which yields a scaling: $1 - F^{(m)}_{1,\mathrm{opt}} \propto 1/\sqrt{P_{1D}}$. To create a superposition $|\Psi_D\rangle$, we require $m_{\max}$ steps. Thus, the total error of the first part of the protocol is $1 - F_1 \propto m_{\max}/\sqrt{P_{1D}}$, that can be improved via postselection conditioned on detecting no photons in the waveguide [47].

To validate the scaling analysis, we study numerically the preparation of two relevant sets of states: (i) the $|D_m\rangle$; (ii) the superpositions $|\Phi_m\rangle = (1/\sqrt{2})(|D_0\rangle + |D_m\rangle)$. Because of the imposed symmetry conditions the relevant Hilbert space depends only on the maximum number of excitations, $m_{\max}$, while the N only enters on the two-photon resonance condition, that fixes $\Omega^{(m)}$ [47]. With this restriction, we use a non-Hermitian evolution governed by $H_{\mathrm{eff}} = H_{\mathrm{las}} + H_c - \mathbf{i}\Gamma_{1D}S_{eg}S_{ge}/2 - \mathbf{i}\Gamma^*S_{ee}/2$. To generate the $|D_m\rangle$, the pulse sequence consists of a complete transfer of populations in each step of Fig. 2(a), which is ensured by fixing the time of interaction, t, to $t\Omega_c(\Omega^{(m)}) = \pi$ for the microwave (Raman) transitions, whereas for the $|\Phi_m\rangle$ the pulse sequence is calculated numerically. In Figs. 2(b) and 2(c), we show the numerical fidelities obtained when fixing the off-resonant transition to the optimal $\Delta_{e,\mathrm{opt}}$, confirming that our arguments give the correct scaling $\propto 1/\sqrt{P_{1D}}$.

Finally, we estimate the fidelity of the photonic state considering the effect of Γ^*. For a superradiant state with m excitations the error rate is $m\Gamma^*$, while the average time to decay is $1/(N\Gamma_{1D})$, which yields an error of $1 - F_2^{(m)} = 2m\Gamma^*/(N\Gamma_{1D})$. When there are $m_{\max}$ excitations in the system, the total fidelity of the process is

$$F_2 \approx 1 - \frac{m_{\max}^2\Gamma^*}{N\Gamma_{1D}}. \quad (12)$$

The dissipative character of this mapping allows for the efficient generation of (arbitrary superpositions of) photonic states, e.g., Fock states, that typically are generated probabilistically [58–60] or via nonlinear interactions [39].

An appealing platform to implement these ideas is cold atoms trapped near photonic crystal waveguides [6–10], where $\Gamma_{1D}/\Gamma_a = \xi n_g\sigma/(2A_m)$, where $n_g = c/v_g$ is the group index, $\sigma = 3\lambda_0^2/(2\pi)$ the radiative cross section, A_m the effective mode area, Γ_a the vacuum emission rate, and ξ a cavity enhancement factor. Current values for Cs atoms ($\lambda_0 = 894$ nm, $\Gamma_a/2\pi = 5.02$ MHz) and SiN alligator waveguides [6,10] have $A_m \approx 0.2\ \mu\mathrm{m}^2$, $n_g \approx 10$, $\xi \sim 5$, and Q factors of 10^6. Depending on the reduction of spontaneous emission, $\Gamma^* = \alpha\Gamma_a$, these numbers lead to $P_{1D} \approx 50/\alpha$. Intrinsic losses in the dielectric and reduced v_g set finite propagation lengths of waveguide modes, $L_{\mathrm{prop}}/\lambda_a \approx Q/(2\pi n_g)$, which is $> 10^4$ for state-of-art SiN values [6,10]. Retardation effects also set a maximum number of atoms for superradiant atom-photon mapping, e.g., assuming a separation $\lambda_a/2$, then $N\Gamma_{1D} < 2v_g/(N\lambda_a)$, which for current structures leads to $N \lesssim 500$ atoms. Possible ways of avoiding retardation are to increase Γ_{1D} by increasing ξ (not n_g) [10]; or by doing the atom-photon mapping off-resonantly, which decreases Γ_{1D} while keeping P_{1D} constant. Other potential problems that we neglected are (i) imperfect atomic separations limited ultimately by center of mass wave packets and atomic motion and (ii) group velocity dispersion that distorts the propagating wave packet. In Ref. [47], we estimate under which conditions they can be neglected; however, a thorough study should be done for each implementation to minimize the impact on the protocols described.

In conclusion, we have presented a protocol to generate deterministic superpositions of many-body entangled atomic states in the presence of strong dissipation. Remarkably, the errors in the preparation of complex superposition states increase only linearly with the number of excitations of the system and inversely with the square root of the Purcell factor. Furthermore, we have shown how to map these atomic states to photonic states with a very efficient scaling that depends linearly on the inverse collective Purcell Factor and how to engineer a time-symmetric wave packet that guarantees the reversibility of the mapping.

We gratefully acknowledge discussions with T. Shi. The work of A. G. T., V. P., and J. I. C. was funded by the European Union integrated project *Simulators and*

Interfaces with Quantum Systems (SIQS). A. G. T. also acknowledges support from Alexander Von Humboldt Foundation and Intra-European Fellowship NanoQuIS (625955). J. I. C. acknowledges support as a Moore Distinguished Scholar. D. E. C. acknowledges support from Fundacio Privada Cellex Barcelona, Marie Curie CIG ATOMNANO, and ERC Starting Grant FoQAL. H. J. K. acknowledges funding by the Institute of Quantum Information and Matter, a National Science Fundation (NSF) Physics Frontier Center with support of the Moore Foundation, by the Air Force Office of Scientific Research, Quantum Memories in Photon-Atomic-Solid State Systems (QuMPASS) Multidisciplinary University Research Initiative (MURI), by the Department of Defense National Security Science and Engineering Faculty Fellows (DoD NSSEFF) program, by NSF PHY1205729, and support as a Max Planck Institute for Quantum Optics Distinguished Scholar.

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