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Efficient multiphoton generation in waveguide QED

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Engineering quantum states of light is at the basis of many quantum technologies such as quantum cryptography, teleportation or metrology among others. Though, single-photons can be generated in many scenarios, the efficient and reliable generation of complex *single-mode* multiphoton states is still a long-standing goal in the field, as current methods either suffer from low fidelities or small probabilities. Here we discuss several protocols which harness the strong and long-range atomic interactions induced by waveguide QED to efficiently load excitations in a collection of atoms, which can then be triggered to produce the desired multi-photon state. In order to boost the success probability and fidelity of each excitation process, atoms are used to both generate the excitations in the rest, as well as to herald the successful generation. Furthermore, to overcome the exponential scaling of the probability of success with the number of excitations, we design a protocol to merge excitations that are present in different internal atomic levels with a polynomial scaling.

On-demand generation of optical propagating photons is at the basis of many applications in quantum information science, including multipartite teleportation [1], quantum repeaters [2], cryptography [3, 4], and metrology [5]. While single photons are routinely produced in different experimental setups [6], e.g., by using natural/artificial atoms coupled to cavities or waveguides [7–12], *single–mode multiphoton states* are much harder to generate [13]. Current methods are limited by either exponentially small success probabilities or low fidelities. The enhancement of light-matter interactions provided by quantum nanophotonics opens up new avenues to create high-fidelity multiphoton states. For example, m quantum emitters can be strongly coupled to structured waveguides, which show large Purcell factors, $P_{1\mathrm{d}}$, so that m atomic excitations can be mapped to a waveguide mode with an error (or infidelity, I_m) scaling as $m/P_{1\mathrm{d}}$. However, the resulting state is not a single mode, but a complex entangled state of several modes [14], so that it cannot be directly used for quantum information purposes. Single-mode multiphoton states can be created by storing m collective excitations in $N \gg m$ atoms, which are then mapped to a photonic state of the waveguide. While the latter process can be achieved with very low infidelity, scaling as $m^2/(NP_{1\mathrm{d}})$ [14, 15], present schemes for the first part scale like $I_m \propto m/\sqrt{P_{1\mathrm{d}}}$ [14], as they still do not fully exploit the strong coupling to the waveguide nor collective effects. This ultimately limits the fidelity of the whole procedure.

In this work we show how to overcome this limitation with new schemes for the heralded generation of m collective excitations in $N \gg m$ atoms coupled to a waveguide. The idea is to use the atoms to both create the excitations one-by-one, and to herald the success of the process. In this way, arbitrarily small infidelities, I_m, can be obtained at the expense of making the process non-deterministic. Depending on the scheme, we find that the global probability of success (or, inversely, the average number of operations, R_m) decreases (increases) exponentially with m, and thus it cannot be scaled to arbitrarily large photon production. Finally, we also show how to overcome this exponential law by using additional atomic states, atom number resolved detection, and a specific protocol to merge excitations, while keeping a low global infidelity, $I_m \sim \mathrm{poly}(m)/(NP_{1\mathrm{d}})$.

Structured waveguide setups with trapped atoms offer interesting characteristics that we exploit to design our protocols, namely, i) regions of large Purcell Factor $P_{1\mathrm{d}} \gg 1$, e.g., due to the reduced group velocity (v_g) in engineered dielectrics [9, 16–19], with simultaneously ii) long-propagation lengths of the guided modes compared to the characteristic wavelength (λ_{a}) that give rise to long-range dissipative couplings [20]. Moreover, as shown in, e.g. Ref. [20], in order to avoid dipole-dipole interactions and fully exploit superradiance effects we assume iii) the atoms to be placed at distances $z_n = n\lambda_{\mathrm{a}}$, with $n \in \mathbb{N}$ [21]. Choosing those positions, and assuming the timescales of the waveguide modes are much faster than the atomic ones (Born-Markov approximation), the waveguide induces an effective dynamics of the atoms described by [22]:

$$\dot{\rho} = \frac{\Gamma_{1\mathrm{d}}}{2}\left(2S_{ge}\rho S_{eg} - S_{eg}S_{ge}\rho - \rho S_{eg}S_{ge}\right) \tag{1}$$

where ρ is the density matrix describing the atomic state, $S_{\alpha\beta} = (N)^{-1/2}\sum_{n=1}^{N}\sigma^n_{\alpha\beta}$ are the collective spin operators and denoting $\sigma^n_{\alpha\beta} = |\alpha\rangle_n\langle\beta|$ the spin operator corresponding to the n-th atom. Finally, iv) it is possible to read the collective atomic state efficiently through the waveguide due to the short timescales and large collection efficiencies. We use atomic detection for heralding, which has reached accuracies of 10^{-4} [23] in trapped ion setups, where the collection efficiency is not enhanced by the presence of a waveguide. Thus, we assume perfect atomic detection, and consider $P_{1\mathrm{d}}$ and N as the resources to analyze the figures of merit of the proposal.

For the analysis of three different schemes, we adopt the following strategy: first, we study the process of heralding the generation of a single collective excitation in an ensemble already storing m excitations, i.e., $|\Psi_m\rangle$, denoted as *target* ensemble, by measuring a given state of an auxiliary system,

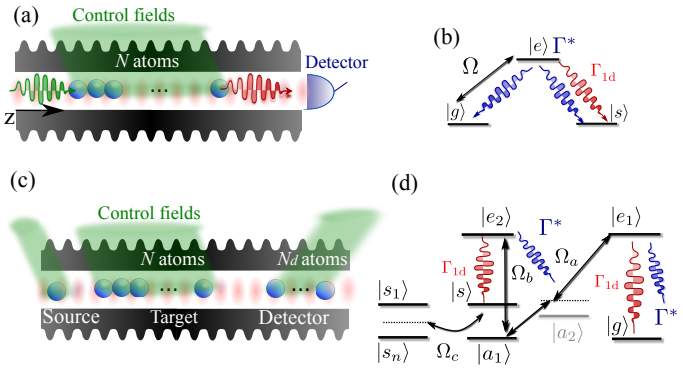

Figure 1. (a) [(c)] Setup for the first [second/third] protocol: N *target* atoms [one *source* and N_d *detector* atoms] are coupled to the waveguide. Photon detection [atomic detection in the detector ensemble] heralds the addition of collective excitations. (b) Atomic level structure for the first protocol: waveguide modes are coupled to $e \leftrightarrow s$ transition with an emission rate Γ_{1d}. Transition $g \leftrightarrow e$ is driven by a Raman laser Ω and the spontaneous emission rate to other modes is denoted by Γ^*. (d) Atomic level structure for second/third protocols: waveguide modes are coupled to the $e_1 \leftrightarrow g$ and $e_2 \leftrightarrow s$ transition with an emission rate Γ_{1d}. A two-photon transition $a_1 \leftrightarrow e_1$ is driven by laser light via level a_2 with effective Rabi frequency Ω_a. Transition $a_1 \leftrightarrow e_2$ is driven by other laser with Rabi frequency Ω_b. Coupling between levels $s \leftrightarrow s_n$ is given by a Rabi frequency Ω_c.

e.g., $|\Phi_{\rm her}\rangle$. The analysis begins by considering the initial state of the combined system $\rho_T(0) = |\Psi_m\rangle\langle\Psi_m| \otimes \rho_{\rm aux}(0)$, where the auxiliary state is initialized in a state different from $|\Phi_{\rm her}\rangle$. Then, we apply our protocols consisting of a combination of laser pulses plus the collective dissipation of Eq. 1 induced by the waveguide, which evolves $\rho_T(0)$ after a time T_f to $\rho_T(T_f)$. We characterize the heralding by calculating both the probability of success (p) and the corresponding error or infidelity ($I_{m\to m+1}$) as follows:

$$p = {\rm Tr}\left[\langle\Phi_{\rm her}|\rho_T(T_f)|\Phi_{\rm her}\rangle\right]. \tag{2}$$

$$I_{m\to m+1} = 1 - \frac{\langle\Psi_{m+1}|\langle\Phi_{\rm her}|\rho_T(T_f)|\Phi_{\rm her}\rangle|\Psi_{m+1}\rangle}{p}. \tag{3}$$

Finally, we analyze how to combine the heralded single excitations to accumulate m excitations in the target ensemble, characterizing the process by the average number of operations R_m one has to perform, together with the final infidelities I_m. To simplify the expressions along the main text, we assume $N \gg m$ and $P_{1d} \gg 1$, though complete expressions can be found in Supplementary Material [24].

Let us first analyze a protocol inspired in a method originally devised to create long-distant entangled states in atomic ensembles [25], that requires N atoms placed close to a 1d waveguide [see Fig. 1(a)]. The atoms must have the level structure depicted in Fig. 1(b), where the transition $e \leftrightarrow s$ is coupled to the guided modes at a rate Γ_{1d} and the transition $g \leftrightarrow e$ is driven equally by a laser with Rabi frequency Ω. The excited states e also radiate into leaky modes (of the waveguide or outside) other than the relevant one, that give rise to a decay rate Γ^*, leading to the Purcell factor $P_{1d} = \Gamma_{1d}/\Gamma^*$. The excitations are stored in the states g and s, which are assumed to be decoherence-free like any other hyperfine ground state. We assume to start with m collective excitations in level s, i.e., $|\Phi_m^s\rangle \propto {\rm sym}\{|s\rangle^{\otimes m} \otimes |g\rangle^{N-m}\}$, where sym denotes the symmetrizing operator. The idea is to weakly excite the atoms in g to level e with a short laser pulse of duration $T \ll 1/(N\Gamma_{1d})$ and if a photon in the waveguide is detected, it heralds the addition of an excitation in state s, i.e., $|\Phi_{m+1}^s\rangle$ with respect to which the infidelity $I_{m\to m+1}$ is defined. As all the atoms are equally coupled to the waveguide, the excitation will be collective.

Let us denote by $x = \Omega\sqrt{N}T/2 \ll 1$, so that right after the pulse the state (up to a normalization constant) is $\left[\mathbf{1} + xS_{eg} + x^2 S_{eg}^2 + O(x^3)\right]|\Phi_m^s\rangle$. After the pulse, the system evolves under Eq. 1 and the interaction with the leaky photonic bath (for a time $t \gg 1/\Gamma^*$), in which case the wavefunction terms with excitations in e decay either to waveguide/leaky photons. If a waveguide photon is detected, either it comes from the lowest order term, $O(x)$, in which case the atomic state is the desired one, i.e., $S_{sg}|\Phi_m^s\rangle$, or from the double excited state, in which case we will prepare the wrong state introducing an error. Emission of leaky photons also produces errors, but of smaller order [24]. Denoting by η the detection efficiency, the success probability and infidelity in heralding the desired state (to lowest order in x) reads:

$$p \approx \eta x^2\left(1 - \frac{1}{P_{1d}}\right), I_{m\to m+1} \propto (1-\eta)x^2, \tag{4}$$

To create m excitations, one has to detect m consecutive photons, leading to $R_m = p^{-m}$ and $I_m \propto m(1-\eta)x^2$. The error can be made arbitrarily small at the expense of decreasing the success probability. Thus, if a high fidelity is required the method is practicable only for few excitations. To reduce R_m one can use an additional metastable state s_1 in which the heralded excitation is stored after each successful addition. This can be done by combining, e.g., a two-photon Raman transition, to make a beam splitter transformation between levels s and s_1, and post-selection conditioned on no atomic detection in s. Then, assuming that m excitations are stored in s_1, we generate $m+1$ within the same level. This strategy leads to $R_m \propto e^m/p$ [24, 26].

To overcome the trade-off between probabilities and fidelities coming from zero and double excitations, we propose a protocol relying on a configuration depicted in Fig. 1(c): the write field of the previous scheme is replaced by a single *source* atom that guarantees the transfer of at most a single excitation to the *target* ensemble. Furthermore, in a second step, N_d *detector* atoms herald the transfer of excitations, replacing the photon detector. Both the source and detector atoms should be separated from the target ensemble for independent addressing with external fields. The protocol requires a level structure as shown in Fig. 1(d) where two dipolar transitions $e_1 \leftrightarrow g$ and $e_2 \leftrightarrow s$ are coupled to the same waveguide mode with Purcell factor P_{1d} that we set to be equal for simplicity. We require the use of other hyperfine, auxiliary levels, a_1, a_2, s_1. The transition $a_1 \leftrightarrow e_2$ is connected by a laser,

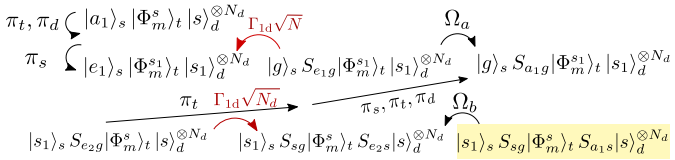

Figure 2. Steps for the second protocol, with $\pi_{c,t,d}$ denoting the π-pulses for the population transfers within the source/target/detector atoms. The driven transitions to the excited states are indicated by the Rabi frequencies $\Omega_{a,b}$. The final state is highlighted in yellow.

whereas the $a_1 \leftrightarrow e_1$ is a two-photon transition mediated by a_2, with effective Rabi frequency Ω_a, so that direct spontaneous emission from $e_1 \to a_1$ is forbidden [15, 24, 27, 28]. The level s_1 is used to store excitations and decouple them from the dynamics induced by the waveguide. We require that s_1 is not connected to neither e_1 nor e_2 by a dipole transition, so that it is only connected to s through microwave fields.

This protocol starts with the target ensemble $|\Phi_m^s\rangle_t$, and with the source/detector atoms in a_1/s respectively. The heralded transfer of a single collective excitation consists of several steps [Fig. 2]. The first one, coherently transfers the excitations from $s \to s_1$ in the target and detector ensemble to protect them from the waveguide dynamics. The second step uses a short laser pulse in the source atom to excite it to state e_1, and then switches on the lasers driving $e_1 \leftrightarrow a_1$ via a two-photon transition in the target ensemble with (effective) Rabi frequency Ω_a for a time T_a. Ideally, the source atom exchanges the excitation with the target ensemble via the waveguide, thus generating a collective excitation in a_1-state of the target ensemble. After that, the laser Ω_a is turned off and one waits for a time $t \gg (\Gamma^*)^{-1}$ such that any remaining population in the excited state decays. Thirdly, we apply π-pulses to decouple the source atom, putting it in s_1, and moving the target and detector ensemble excitations back to s. Another short pulse is applied to move the collective excitation in the target ensemble from a_1 to e_2 and we switch on the laser Ω_b driving the $e_2 \leftrightarrow a_1$ transition in the detector ensemble for a time T_b with Rabi frequency Ω_b. This transfers the collective excitation in the target ensemble to s and creates a collective excitation in a_1 in the detector ensemble. At the end, a measurement of state a_1 of the detector atoms (through fluorescence in the waveguide by exciting it to an optically excited state) heralds the successful preparation of a collective excitation in the target ensemble, i.e., $|\Phi_{m+1}^s\rangle_t$, with respect to which the infidelity $I_{m\to m+1}$ is defined.

Let us analyze the protocol in detail. In the second step, the evolution of the source/target atoms is described by a master equation given by both the collective dissipation of Eq. 1, individual dissipation from leaky modes, and the unitary dynamics induced by the laser Ω_a, which can be analytically solved in the limit $NP_{1d} \gg \Omega_a$ [24]. By choosing the Rabi frequency $\Omega_a \approx \sqrt{N\Gamma_{1d}\Gamma^*}$, and $T_a = \pi/\sqrt{\Gamma_{1d}\Gamma^*}$, the probability for the ensemble to end up in the desired state after the second step, $|g\rangle_s \otimes S_{a_1 g}|\Phi_m^{s_1}\rangle_t$, is maximized $p_a \approx e^{-\pi/\sqrt{P_{1d}}}$. Similarly, in the third step, the evolution of the source/detector atoms can be analytically solved in the limit $N_d P_{1d} \gg \Omega_b$, obtaining a probability to end up in the desired state $S_{sg}|\Phi_m^s\rangle_t \otimes S_{a_1 s}|s\rangle_d^{\otimes N_d}$ given also by $p_b \geq p_a$. Thus, the total probability of success of the protocol and infidelity are given by:

$$p \gtrsim p_a^2 \approx e^{-2\pi/\sqrt{P_{1d}}}, I_{m\to m+1} = 0, \tag{5}$$

as we rule out all the possible errors through post-selection. The only way the detector atoms arrive to a_1 is that the steps two and three have occurred as desired. Any spontaneous emission in any of the atoms or photon absorption is incompatible with that event. In case of unsuccessful detection we pump all the target atoms back to g and restart the process, such that to accumulate m excitations, one needs an average number of operations $R_m \propto p^{-m}$, but still with $I_m = 0$. As p is close to 1 when $N, N_d, P_{1d} \gg 1$, we can achieve large number of excitations despite the exponential scaling of R_m.

After this analysis the question arises whether the exponential number of operations is a fundamental problem of probabilistic protocols. This leads to our third and final scheme in which we design a protocol which circumvents this exponential scaling with a judicious modification of the previous protocol by using several additional atomic states s_n [see Fig. 1(d)] and atomic number resolved detection. The idea is that with each successful detection, we transfer the heralded collective excitation in s not to the same level s_1, but to one of several states $\{s_n\}$ to then *merge* them a posteriori with an adequate protocol that we explain below. In contrast to our previous schemes, in case of unsuccessful detection, i.e., the detector atoms were not found in a_1, the m excitations stored in $\{s_n\}_n$ are not destroyed, but only the one we wanted to add. We can pump back the target atoms in s, a_1 to g and try again. The price one has to pay is that errors appear since one may not recover the original collective state in the ensemble, i.e., because a spontaneous emission event occurred within the target ensemble. One can show that the main source of errors occurs when a collective excitation was indeed produced in the ensemble, but not detected (because, e.g., of spontaneous emission in the detector ensemble). To reduce this error, we have to ensure that undetected collective states return back to g coherently which can be done by exploiting the collective coupling to the waveguide [24]. Moreover, to obtain the desired scaling of infidelity $(1/(NP_{1d}))$, we also need to modify the third step of the protocol for which then, only a single detector atom is needed, $N_d = 1$. It can be shown that once we are in $S_{a_1 g}|\Phi_m^{\{s_n\}_n}\rangle|s\rangle_d$, the optimal strategy is to apply a fast π-pulse in the target atoms such that $S_{e_2 g}|\Phi_m^{\{s_n\}_n}\rangle_t$ is prepared. Then, the system is left free to interact only through the collective dissipation of Eq. 1 for a time $T_b = 1/\Gamma_{1d}$. This dynamics is terminated by applying another π-pulse on the target and detector atoms, which puts any possible excited state back in a_1. We find that the optimal probability for the system to be in $S_{sg}|\Phi_m^{\{s_n\}_n}\rangle_t \otimes |a_1\rangle_d$ is $p_b \approx 0.1\,(1 - 1/P_{1d})$. Finally, if we fail, we pump any possible excitation of the target ensemble coherently back to g and repeat the process until success, of

the order of $1/p$ times with $p = p_a p_b$. Integrating the corresponding master equation [24], identifying all errors and how they accumulate in the repetitions, we arrive to an (averaged) infidelity with the desired scaling:

$$I_{m\to m+1} \lesssim m/(pNP_{1d}) \quad (6)$$

Thus, the problem reduces to *merging* the atomic excitations distributed in the levels $\{s_n\}_n$. For example, if instead of adding excitations one by one, we use $\log_2 m$ metastable states s_n, and adopt a tree-like structure, we obtain a superpolynomial scaling [24, 29] in $R_m \propto m^{\log_2 m}/p$ to merge m collective excitations in a single atomic level. The key step is to double the number of excitations in each step using beam splitter transformations and post-selection conditioned on atomic detection in one of the levels.

Finally, we go one step further by using number-resolved atomic detection to obtain a polynomial scaling. The key point is to realize that if we have n excitations stored in two atomic levels, after applying a beam splitter transformation, the probability to obtain exactly $2n$ excitations in one of them (by detecting no excitation in the other) decays with n; however, the probability of obtaining more than $3n/2$ in one level (by detecting $p < n/2$ in the other) is actually bounded below by $1/3$, independent of n [24]. Assuming the worst case scenario in which after detecting $q < m/2$ excitations in one of the two levels, we assume that the other state only goes up to $3n/2$ excitations, that gives us an upper bound for the number of operations [24] $R_m \lesssim m^{4.41}\log_{3/2}(m)/p$, that is already polynomial. The number of atomic levels, s_i, required to reach m excitations scales logarithmically $\log_{3/2} m$, and the final infidelity scales as $I_m \propto \text{poly}(m)/(NP_{1d})$. We note that atomic detection itself introduces errors, however, the aim of this scheme is to show that polynomial scaling is possible, despite not being currently the most efficient method.

State-of-the art technologies provide systems with $N \sim 1-2$ emitters coupled to engineered dielectrics [16–19, 30, 31] with $P_{1d} \sim 1-100$. Advances in both fabrication and trapping techniques foresee implementations with $N \sim 100$ atoms and $P_{1d} \gtrsim 100$ in the near future. Atomic internal levels may be replaced by motional levels of each of the atoms, if they are trapped in pseudo-harmonic potentials [32], since there can be many of them at our disposal. Finite atomic detection efficiency (η_d) adds up errors proportional to the number of measurements $\sim R_m(1-\eta_d)$. There are other sources of errors not considered here and that will give limitations to our proposal, such as decoherence of hyperfine levels, laser fluctuations, retardation, which can be neglected if $\Gamma_{1d} N \ll v_g/(N\lambda_a)$, imperfect atomic cooling or positioning, which can be controlled to a large extent but has to be considered to achieve large fidelities (see [14] where part of these errors were considered).

In conclusion, we present three probabilistic protocols to generate heralded entangled atomic states that afterwards can be mapped to photonic states at will with very high fidelities. In particular, we show how to accumulate m collective atomic excitations with infidelity $I_m = 0$ and an exponential number of operations $R_m = p^{-m}$, being p the heralding probability for adding a single excitation. We design a protocol where p can be close to unity for systems with $N, P_{1d} \gg 1$, which allow to accumulate $m \sim 20-50$ excitations with systems $P_{1d} \sim 10^2 - 10^3$ using $R_m \sim 10^4$ operations with unit fidelity. Moreover, we present a protocol with polynomial scaling in R_m by using number resolved atomic detection and overall low infidelity $I_m \propto \text{poly}(m)/(NP_{1d})$. With suitable modification of the level structure and using two guided modes, our protocols can be extended to generate NOON-like states [33] and can also be exported to other systems such as low-mode cavities [30, 31, 34, 35], superconducting circuits [36] or optical fibers [37–41].

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[1] Murao, M., Jonathan, D., Plenio, M. B. & Vedral, V. Quantum telecloning and multiparticle entanglement. *Phys. Rev. A* **59**, 156–161 (1999).

[2] Briegel, H.-J., Dür, W., Cirac, J. I. & Zoller, P. Quantum Repeaters: The Role of Imperfect Local Operations in Quantum Communication. *Phys. Rev. Lett.* **81**, 5932 (1998).

[3] Gisin, N., Ribordy, G., Tittel, W. & Zbinden, H. Quantum cryptography. *Rev. Mod. Phys.* **74**, 145 (2002).

[4] Durkin, G. A., Simon, C. & Bouwmeester, D. Multiphoton entanglement concentration and quantum cryptography. *Phys. Rev. Lett.* **88**, 187902 (2002).

[5] Giovannetti, V., Lloyd, S. & Maccone, L. Quantum metrology. *Phys. Rev. Lett.* **96**, 010401 (2006).

[6] Lounis, B. & Orrit, M. Single-photon sources. *Reports on Progress in Physics* **68**, 1129 (2005).

[7] Xu, Y. *et al.* Finite-difference time-domain calculation of spontaneous emission lifetime in a microcavity. *JOSA B* **16**, 465–474 (1999).

[8] Painter, O. *et al.* Two-dimensional photonic band-gap defect mode laser. *Science* **284**, 1819–1821 (1999).

[9] Hughes, S. Enhanced single-photon emission from quantum dots in photonic crystal waveguides and nanocavities. *Optics letters* **29**, 2659–2661 (2004).

[10] Rao, V. M. & Hughes, S. Single quantum-dot purcell factor and β factor in a photonic crystal waveguide. *Physical Review B* **75**, 205437 (2007).

[11] Arcari, M. *et al.* Near-unity coupling efficiency of a quantum emitter to a photonic crystal waveguide. *Phys. Rev. Lett.* **113**, 093603 (2014).

[12] Somaschi, N. *et al.* Near-optimal single-photon sources in the solid state. *Nature Photonics* (2016).

[13] Dell'Anno, F., De Siena, S. & Illuminati, F. Multiphoton quantum optics and quantum state engineering. *Physics reports* **428**, 53–168 (2006).

[14] González-Tudela, A., Paulisch, V., Chang, D. E., Kimble, H. J. & Cirac, J. I. Deterministic generation of arbitrary photonic states assisted by dissipation. *Phys. Rev. Lett.* **115**, 163603 (2015).

[15] Porras, D. & Cirac, J. Collective generation of quantum states of light by entangled atoms. *Physical Review A* **78**, 053816 (2008).

[16] Laucht, A. *et al.* A Waveguide-Coupled On-Chip Single-Photon Source. *Phys. Rev. X* **2**, 011014 (2012).

[17] Lodahl, P., Mahmoodian, S. & Stobbe, S. Interfacing single photons and single quantum dots with photonic nanostructures. *Rev. Mod. Phys.* **87**, 347–400 (2015).

[18] Yu, S.-P. *et al.* Nanowire photonic crystal waveguides for single-atom trapping and strong light-matter interactions. *Appl. Phys. Lett.* **104**, 111103 (2014).

[19] Goban, A. *et al.* Superradiance for atoms trapped along a photonic crystal waveguide. *Phys. Rev. Lett.* **115**, 063601 (2015).

[20] Chang, D., Jiang, L., Gorshkov, A. & Kimble, H. Cavity qed with atomic mirrors. *New Journal of Physics* **14**, 063003 (2012).

[21] All results can be extended to separations of $\lambda_a/2$ if the phases of the laser amplitudes are adjusted correctly.

[22] Gardiner, G. W. & Zoller, P. *Quantum Noise* (Springer-Verlag, Berlin, 2000), 2nd edn.

[23] Myerson, A. H. *et al.* High-fidelity readout of trapped-ion qubits. *Phys. Rev. Lett.* **100**, 200502 (2008).

[24] See Supplementary material for full details of the calculations of probabilities and fidelities of the protocols discussed in the main text.

[25] Duan, L.-M., Lukin, M. D., Cirac, J. I. & Zoller, P. Long-distance quantum communication with atomic ensembles and linear optics. *Nature* **414**, 413 (2001).

[26] Dakna, M., Clausen, J., Knöll, L. & Welsch, D.-G. Generation of arbitrary quantum states of traveling fields. *Phys. Rev. A* **59**, 1658–1661 (1999).

[27] van Enk, S. J., Cirac, J. I. & Zoller, P. Ideal quantum communication over noisy channels: A quantum optical implementation. *Phys. Rev. Lett.* **78**, 4293–4296 (1997).

[28] Borregaard, J., Kómár, P., Kessler, E. M., Sørensen, A. S. & Lukin, M. D. Heralded quantum gates with integrated error detection in optical cavities. *Phys. Rev. Lett.* **114**, 110502 (2015).

[29] Fiurášek, J., García-Patrón, R. & Cerf, N. J. Conditional generation of arbitrary single-mode quantum states of light by repeated photon subtractions. *Phys. Rev. A* **72**, 033822 (2005).

[30] Thompson, J. D. *et al.* Coupling a single trapped atom to a nanoscale optical cavity. *Science* **340**, 1202–1205 (2013).

[31] Tiecke, T. *et al.* Nanophotonic quantum phase switch with a single atom. *Nature* **508**, 241–244 (2014).

[32] Parkins, A. S., Marte, P., Zoller, P., Carnal, O. & Kimble, H. J. Quantum-state mapping between multilevel atoms and cavity light fields. *Phys. Rev. A* **51**, 1578–1596 (1995).

[33] Paulisch, V., Gonzalez-Tudela, A., Kimble, H. J. & Cirac, J. I. *In preparation* .

[34] Specht, H. P. *et al.* A single-atom quantum memory. *Nature* **473**, 190–193 (2011).

[35] Haas, F., Volz, J., Gehr, R., Reichel, J. & Estève, J. Entangled states of more than 40 atoms in an optical fiber cavity. *Science* **344**, 180–183 (2014).

[36] Mlynek, J., Abdumalikov, A., Eichler, C. & Wallraff, A. Observation of dicke superradiance for two artificial atoms in a cavity with high decay rate. *Nat. Communications* **5**, 5186 (2014).

[37] Paulisch, V., González-Tudela, A., Kimble, H. J. & Cirac, J. I. Heralded multiphoton states with coherent spin interactions in waveguide qed. *New Journal of Physics* **19**, 043004 (2017).

[38] Vetsch, E. *et al.* Optical interface created by laser-cooled atoms trapped in the evanescent field surrounding an optical nanofiber. *Phys. Rev. Lett.* **104**, 203603 (2010).

[39] Goban, A. *et al.* Demonstration of a State-Insensitive, Compensated Nanofiber Trap. *Phys. Rev. Lett.* **109**, 033603 (2012).

[40] Petersen, J., Volz, J. & Rauschenbeutel, A. Chiral nanophotonic waveguide interface based on spin-orbit interaction of light. *Science* **346**, 67–71 (2014).

[41] Béguin, J.-B. *et al.* Generation and detection of a sub-poissonian atom number distribution in a one-dimensional optical lattice. *Phys. Rev. Lett.* **113**, 263603 (2014).