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Quantum storage of single photons with unknown arrival time and pulse shapes*

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We present a scheme for the quantum storage of single photons using electromagnetically induced transparency (EIT) in a low-finesse optical cavity, assisted by state-selected spontaneous atomic emission. Mediated by the dark mode of cavity EIT, the destructive quantum interference between the cavity input–output channel and state-selected atomic spontaneous emission leads to strong absorption of single photons with unknown arrival time and pulse shapes. We discuss the application of this phenomenon to photon counting using stored light.

Keywords: quantum information and processing, quantum electrodynamics

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1. Introduction

Quantum information is easily transmitted by traveling photons, but the photons need to be stored locally to process the information.[1] Motivated by this idea, techniques that facilitate the quantum storage of single photons are actively being explored.[2–5] A promising approach to the quantum storage of single photons uses classical laser fields to manipulate pulses of light in optically dense media, such as atomic ensembles, using the electromagnetically induced transparency (EIT) technique in free space[6] or in an optical cavity.[7] Nowadays, EIT-based light storage is extensively being studied, due to its high storage efficiency,[8–10] long lifetime,[11–16] spatial multiplexing,[17] broadband character[18] and its ability to store the quantum states of light.[19–27] The key to quantum storage that uses EIT is to control the power and shape of the control field in such a way that it depends on the arrival time and pulse shapes of the input photons.[28,29] However, due to the intrinsic randomness of a single photon in real-world applications, the arrival time and pulse shape of a single photon are always unknown.

In this paper, we present an alternative scheme for quantum storage using EIT in a low-finesse optical cavity assisted by state-selected spontaneous atomic emission.[30] Under certain conditions, the resonant frequencies of the dark and bright modes of cavity EIT are clearly separated.[31] Under these conditions, input photons whose frequencies are near the resonant frequency of the dark mode are strongly absorbed by the destructive quantum interference between the cavity input–output channel and the state-selected spontaneous atomic emission channel. Compared with previous quantum storage schemes[6–29] that used the EIT technique, the control field in our proposed scheme is independent of the arrival time and pulse shapes of the input photons. Thus, we can continuously activate the coupling field to allow the atom–cavity system to store input photons with unknown arrival time and pulse shapes. We also discuss the application of this technique to photon counting using stored light.[32,33]

2. Quantum storage of single photons

The basic component of our model is an ensemble of N cold atoms trapped in a low-finesse one-sided optical cavity, as illustrated in Fig. 1(a). The atomic transition from a stable ground state $|g\rangle$ to an excited state $|e\rangle$ is resonantly coupled to the cavity mode a with an effective single-atom coupling strength of g, while a coupling laser with a Rabi frequency of Ω resonantly drives the atomic transition from the excited state $|e\rangle$ to a special state $|f\rangle$. Just as in state-selective fluorescence experiments,[30] we choose a state $|f\rangle$ that only decays by spontaneous emission back to a ground state $|s\rangle$. The interaction Hamiltonian for this system is given by

$$H_{\rm I} = H_{\rm EIT} + H_{\rm spon} + H_{\rm in-out}. \tag{1}$$

Here, $H_{\rm EIT} = \sum_{j=1}^{N} \left(g\left|e_j\right\rangle\left\langle g_j\right|a + \Omega\left|f_j\right\rangle\left\langle e_j\right| + \text{H.c.}\right)$ accounts for coherent atom–light coupling in the Λ-type

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EIT configuration, $H_{\rm spon}=\sum_{j=1}^{N}{\rm i}\gamma_f\left|f_j\right\rangle\left\langle f_j\right|/2$ represents the spontaneous emission process, where γ_f is the decay rate of state $|f\rangle$, and $H_{\rm in-out}=\int\omega{\rm d}\omega\Theta^{\dagger}(\omega)\Theta(\omega)+\left[{\rm i}\int{\rm d}\omega\sqrt{\kappa/2\pi}\Theta^{\dagger}(\omega)a+{\rm H.c.}\right]$ describes the cavity input–output process, where κ and $\Theta(\omega)$ are the bare cavity decay rate and the one-dimensional free-space mode, respectively. Assuming that all atoms are initially in the ground state $|g\rangle$ and using the collective operators $C_{\mu}^{\dagger}=\frac{1}{\sqrt{N}}\sum_{j=1}^{N}\left|\mu_j\right\rangle\left\langle g_j\right|$ ($\mu=e,f,s$), the Hamiltonians $H_{\rm EIT}$ and $H_{\rm spon}$ can be rewritten as $H'_{\rm EIT}=\sqrt{N}gC_e^{\dagger}a+\Omega C_f^{\dagger}C_e+{\rm H.c.}$[28] and $H'_{\rm spon}={\rm i}\gamma_f C_f^{\dagger}C_f/2$, respectively.

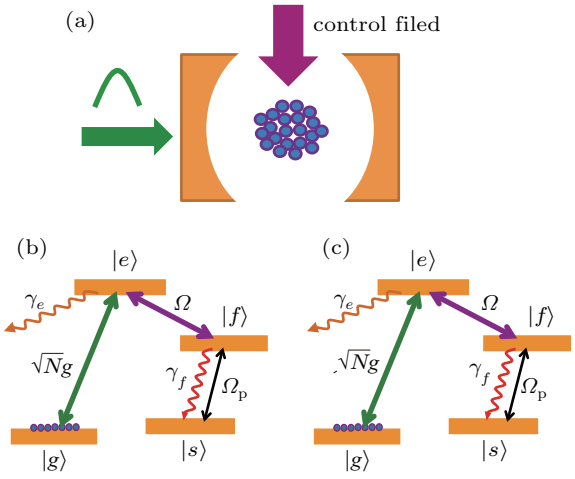


Fig. 1. (a) Schematic setup for the quantum storage of single photons with unknown arrival time and pulse shapes. (b) The relevant atomic-level structure and transitions for the Λ-type EIT configuration assisted by state-selective spontaneous emission $|f\rangle\rightarrow|s\rangle$. (c) The atomic transitions relevant to photon counting. After the photons have been stored in state $|s\rangle$, the control field Ω for EIT is turned off and a pump field $\Omega_{\rm p}$ that is resonant with the atomic transition $|s\rangle\rightarrow|f\rangle$ is turned on. By fluorescence detection, one can realize photon counting for small photon numbers.

To clearly see the photon absorption assisted by state-selected atomic spontaneous emission, one can define two polariton modes for the EIT configuration: a dark-state polariton mode $b_0=a\cos\theta-C_f\sin\theta$ and a bright-state polariton mode $b_1=a\sin\theta+C_f\cos\theta$, where $\cos\theta=\Omega(\Omega^2+Ng^2)^{-1/2}$ and $\sin\theta=\sqrt{N}g(\Omega^2+Ng^2)^{-1/2}$.[6] In terms of these two polariton modes, the Hamiltonians $H'_{\rm EIT}$, $H'_{\rm spon}$, and $H_{\rm in-out}$ can be rewritten as[31]

$$H''_{\rm EIT}=\sqrt{Ng^2+\Omega^2}\left(b_1C_e^{\dagger}+{\rm H.c.}\right), \tag{2}$$

$$H''_{\rm spon}={\rm i}\gamma_f(b_0^{\dagger}\sin\theta-b_1^{\dagger}\cos\theta)(b_0\sin\theta-b_1\cos\theta)/2, \tag{3}$$

$$H'_{\rm in-out}=\int\omega{\rm d}\omega\Theta^{\dagger}(\omega)\Theta(\omega)+\left[{\rm i}\int{\rm d}\omega\sqrt{\kappa/2\pi}\Theta^{\dagger}(\omega)\times(b_0\cos\theta+b_1\sin\theta)+{\rm H.c.}\right]. \tag{4}$$

The bright-state polariton mode b_1 in the Hamiltonian $H''_{\rm EIT}$ resonantly coupled to the collective state C_e will induce an effect analogous to vacuum Rabi splitting, which splits mode b_1 into two normal modes $A_{\pm}$, with frequencies given by $\omega_{\pm}=\pm[\Omega^2+Ng^2]^{-1/2}$. As shown in an experiment using the optomechanical dark mode,[31] when $|\omega_{\pm}|\gg\kappa,\gamma_f,\gamma_e$ (γ_e is the decay rate of state $|e\rangle$), the dark mode b_0 is clearly separated from the bright mode b_1, and thus we can neglect the terms including mode b_1 if the frequencies of the input photons are all near the resonant frequency of the dark mode b_0, thus leading to the reduced form of the Hamiltonian $H_{\rm I}$,

$$H'_{\rm I}={\rm i}\gamma_0b_0^{\dagger}b_0/2+\int\omega{\rm d}\omega\Theta^{\dagger}(\omega)\Theta(\omega)+\left[{\rm i}\int{\rm d}\omega\sqrt{\frac{\kappa_0}{2\pi}}\Theta^{\dagger}(\omega)b_0+{\rm H.c.}\right]. \tag{5}$$

Here, $\gamma_0=\gamma_f\sin^2\theta$ and $\kappa_0=\kappa\cos^2\theta$ are the effective decay rates of the dark mode due to the spontaneous emission of state $|f\rangle$ and cavity decay, respectively.[34]

The Hamiltonian in Eq. (5) is similar to that for a one-sided cavity.[35] According to the quantum Langevin equation, the evolution equation of the dark mode b_0, governed by the Hamiltonian $H=H'_{\rm I}+\omega_0b_0^{\dagger}b_0$, is given by[35]

$$\dot{b}_0=-{\rm i}\left[b_0,(\omega_0+\frac{{\rm i}\gamma_0}{2})b_0^{\dagger}b_0\right]-\frac{\kappa_0}{2}b_0+\sqrt{\kappa_0}a_{\rm in}, \tag{6}$$

where ω_0 is the resonant frequency of the dark mode b_0, and $a_{\rm in}=\frac{1}{\sqrt{2\pi}}\int{\rm d}\omega{\rm e}^{-{\rm i}\omega(t-t_0)}\Theta_0(\omega)$ is the input field operator, where $\Theta_0(\omega)$ is the value of $\Theta(\omega)$ at $t=t_0$ (the input). The relationship between the input and output field operators is $a_{\rm in}(t)+a_{\rm out}(t)=\sqrt{\kappa_0}b_0(t)$,[35] where the output field operator $a_{\rm out}(t)=\frac{1}{\sqrt{2\pi}}\int{\rm d}\omega{\rm e}^{-{\rm i}\omega(t-t_1)}\Theta_1(\omega)$ and $\Theta_1(\omega)$ is the value of $\Theta(\omega)$ at $t=t_1$ (the output). Using Fourier transformations on $\Lambda(t)=\sqrt{\frac{1}{2\pi}}\int_{-\infty}^{+\infty}{\rm d}\omega\Lambda(\omega){\rm e}^{-{\rm i}\omega(t-t_0)}$ with $\Lambda=b_0,a_{\rm in},a_{\rm out}$, we can obtain

$$a_{\rm out}(\omega)=\frac{(\kappa_0-\gamma_0)/2+{\rm i}(\omega-\omega_0)}{(\kappa_0+\gamma_0)/2-{\rm i}(\omega-\omega_0)}a_{\rm in}(\omega). \tag{7}$$

If the frequencies of the input photons are near the resonant frequency of the dark mode, i.e., $|\omega-\omega_0|\ll\kappa_0$, then the expression in Eq. (7) reduces to

$$a_{\rm out}(\omega)=\frac{\kappa_0-\gamma_0}{\kappa_0+\gamma_0}a_{\rm in}(\omega), \tag{8}$$

when $\gamma_0=0$, and $a_{\rm out}(\omega)\approx a_{\rm in}(\omega)$, which means that the input photons are not absorbed. However, when $\gamma_0\neq0$, then $a_{\rm out}(\omega)<a_{\rm in}(\omega)$, i.e., the input photons are absorbed. Furthermore, if $\kappa_0\approx\gamma_0$, we can obtain $a_{\rm out}(\omega)\approx0$, i.e., the probability for photon absorption approaches unity. Therefore, the spontaneous radiation of state $|f\rangle$ is the determining factor for photon absorption. Since the state $|f\rangle$ only decays by spontaneous emission back to the state $|s\rangle$, the input photon is then stored in the long-term storage state $|s\rangle$.

The physical reason for the above results can be explained by quantum interference, as shown by both theory[36,37] and experiment,[31] where the destructive quantum interference between two different cavity input–output channels, mediated by the optomechanical dark mode, induces the transfer of traveling photons from one cavity input–output channel to the other. In our scheme, mediated by the dark mode of cavity EIT, the destructive quantum interference between the input–output channel and the state-selected spontaneous atomic emission channel leads to the transfer of single input

photons from the cavity input–output channel to the state-selected atomic spontaneous emission channel. It should be noted that due to the use of quantum interference, the coupling field in our scheme is not directly coupled to the storage state $|s\rangle$; furthermore, it is independent of the arrival time and pulse shapes of the input photons. In this way, we can continuously activate the coupling field to prepare the atom–cavity system for the storage of the input photons; therefore, our scheme can be distinguished from the scheme that uses the conventional EIT technique,[6] in which the controlling power and the shape of the control field strongly depend on the arrival time and pulse shapes of the input photons.[28] To confirm the results, we can perform exact numerical simulations through the Schrödinger equation $\mathrm{d}|\Psi(t)\rangle/\mathrm{d}t = -\mathrm{i}(H_{\mathrm{EIT}} + H_{\mathrm{in-out}} - \mathrm{i}\gamma_e C_e^{\dagger}C_e/2 - \mathrm{i}\gamma_f C_f^{\dagger}C_f/2)|\Psi(t)\rangle$, where $|\Psi(t)\rangle = \left[\int_{-\infty}^{+\infty}\mathrm{d}\omega A_\omega(t)\Theta^{\dagger}(\omega) + B(t)a^{\dagger} + C(t)C_e^{\dagger} + D(t)C_f^{\dagger}\right]|\mathrm{vac}\rangle\otimes_{j=1}^{N}|g_j\rangle$ is the stochastic wave function describing the system; here, $|\mathrm{vac}\rangle$ denotes that all the free-space modes and the cavity mode are in their vacuum states. For the numerical simulations, we follow the method described in Ref. [38] to discretize the free-space field $\Theta(\omega)$ by introducing a finite but small frequency interval $\delta\omega$ between two adjacent modes. For convenience, the input pulse is taken to be Gaussian with $\varepsilon(t) \propto \exp[-(t-T/2)^2/(T/5)^2]$;[39] here, T is the pulse duration.

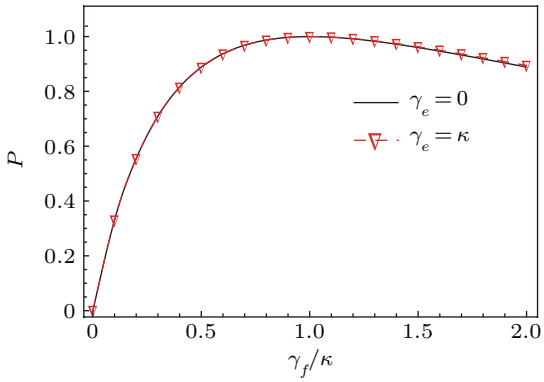


Fig. 2. The probability P for photon absorption versus the spontaneous emission rate γ_f with $\gamma_e = 0$ (solid) and κ (dashed), on the assumption that $\sqrt{N}g = \Omega = 5\kappa$, and $\kappa T = 120$.

Our numerical simulations show that our scheme works remarkably well. In Fig. 2, we plot the probability P for photon absorption versus the spontaneous emission rate γ_f with $\gamma_e = 0$ and κ, on the assumption that $\sqrt{N}g = \Omega$. From Fig. 2, one can see that (i) with an increase in γ_f, the probability P first increases and then decreases, reaching a maximum when $\gamma_f = \kappa$; (ii) the influence of the decay rate γ_e can be ignored, so that the bright-state polariton mode b_1 is strongly inhibited under the condition $\sqrt{Ng^2+\Omega^2} \gg \kappa, \gamma_e, \gamma_s$[31] and the absorption of single photons is almost stored in the state $|s\rangle$ through the decay process from state $|f\rangle$ to the stored state $|s\rangle$. The coupling field in our scheme is not directly coupled to the storage state $|s\rangle$, and thus it can be continuously active, allowing the atom–cavity system to be ready to store a input photon.

Figure 3 shows the probability P for photon absorption versus the normalized time t/κ, for different arrival time and pulse durations: $(t_0, T) = \{(0, 90/\kappa), (0, 210/\kappa), (150/\kappa, 150/\kappa)\}$. It is clear that even though the photons have different arrival time and pulse shapes, they are absorbed by quantum interference. Due to fluctuations in the number of atoms, the number of the atoms trapped in the cavity cannot be exactly controlled in experiments. Figure 4 shows the probability P versus number-of-atom fluctuations ΔN. We note that there are high photon-absorption probabilities (above 93%), even with large number-of-atom fluctuations $|\Delta N| \approx 0.4N$. To return the stored state to the photonic state, the coupling field Ω is turned off, and another laser field $\Omega_R(t)$ coupled to the atomic transition $|s\rangle \leftrightarrow |e\rangle$ is turned on. Adiabatically changing $\Omega_R(t)$ will then map the stored states onto photon states. By adiabatic evolution with an atomic ensemble in a low-finesse optical cavity, experiments have demonstrated the retrieval of the stored states with high efficiencies of up to 84%.[40] One potential application of the scheme analyzed in this paper could be photon counting for a small photon number. One can turn on the coupling field Ω to make the atom–cavity system continuously store input photons. Following storage of the photons, the control field Ω is turned off and a pump field Ω_{p} resonant with the atomic transition $|s\rangle \to |f\rangle$ is turned on, as shown in Fig. 1(c). The atoms will then continuously fluoresce, which

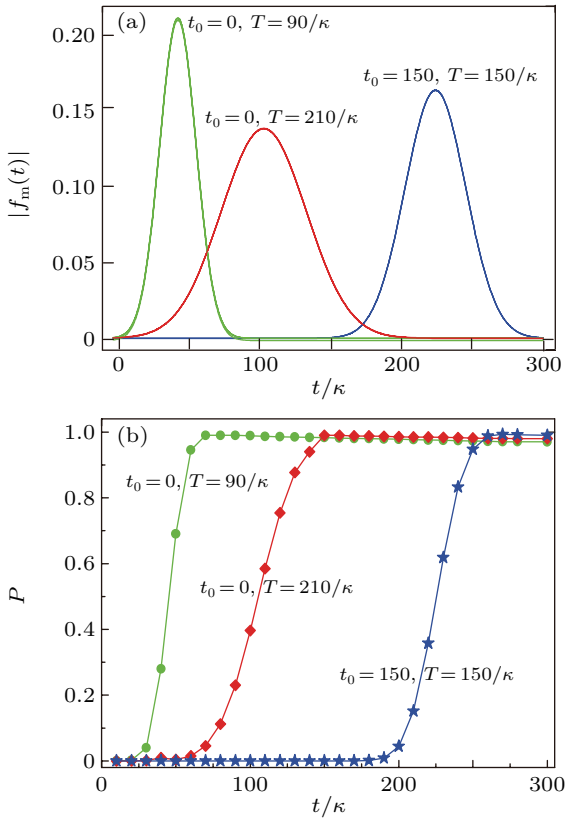


Fig. 3. (a) Gaussian shape for the input single-photon pulse: $\varepsilon(t) \propto \exp[-(t-t_0-T/2)^2/(T/5)^2]$, with different arrival time and pulse durations: $(t_0, T) = \{(0, 90/\kappa), (0, 210/\kappa), (150/\kappa, 150/\kappa)\}$. (b) The probability P for photon absorption versus the normalized time t/κ , on the assumption that $\gamma_f = \gamma_e = \kappa$, and $\sqrt{N}g = \Omega = 5\kappa$.

can be detected with a probability approaching unity. Furthermore, by detecting this fluorescence, one can realize photon counting for a small photon number. We note that using fluorescence detection, James *et al.*[32] and Imamoğlu[33] have proposed schemes for photon counting with far-from-resonant Raman and conventional EIT techniques, respectively. In contrast to these approaches,[32,33] the quantum storage used here even allows the storage of input photons with unknown arrival time and pulse shapes.

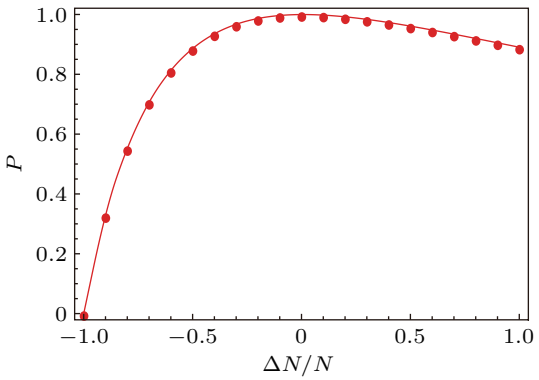


Fig. 4. The probability P for photon absorption with varying number-of-atom fluctuations ΔN. The other parameters are: $\gamma_f = \gamma_e = \kappa$, $\sqrt{N}g = \Omega = 5\kappa$, and $\kappa T = 120$.

3. Feasibility and conclusion

Finally, we address the experimental feasibility of the proposed scheme. The relevant atomic configuration can be chosen from the hyperfine states of ^{87}Rb atoms. For instance, the hyperfine levels $|F = 1, m = 1\rangle$ and $|F = 2, m = 2\rangle$ of $5S_{1/2}$ could correspond to the stable ground states $|g\rangle$ and $|s\rangle$, respectively. The excited state $|F = 2, m = 2\rangle$ of $5P_{1/2}$ corresponds to $|e\rangle$, while the state $|F = 3, m = 3\rangle$ of $5P_{3/2}$ corresponds to $|f\rangle$. This choice guarantees that the state-selected spontaneous emissions $|f\rangle \to |s\rangle$. The coupling of atomic transitions $|e\rangle \leftrightarrow |f\rangle$ can be realized by using a Raman transition via an intermediate level. Even in a low-finesse cavity, a large ensemble of N cold atoms can be in the collective strong-coupling regime $Ng^2/(\kappa\gamma_e) \approx 27 \gg 1$.[41] Hence, our scheme is a good fit for the status of current experimental technology. In summary, we have proposed a scheme for photon storage that uses quantum interference with an atomic ensemble in a low-finesse optical cavity. For this scheme, we have explored the destructive quantum interference between the cavity input–output channel and a state-selected spontaneous atomic emission, which leads to the strong absorption of input photons. We have shown that this effect can be used for the storage of single photons with unknown arrival time and pulse shapes. Our scheme could find application in the photon counting of small photon numbers using stored light. The quantum interference between two different dissipated channels would provide an alternative way to study dissipative quantum phenomena, as well as a powerful strategy for exploring them.

References

[1] Kimble H J 2008 *Nature* **453** 1023
[2] Lukin M D 2003 *Rev. Mod. Phys.* **75** 457
[3] Hammerer K, Sørensen A S and Polzik E S 2010 *Rev. Mod. Phys.* **82** 1041
[4] Specht H P, Nölleke C, Reiserer A, Uphoff M, Figueroa E, Ritter S and Rempe G 2011 *Nature* **473** 190
[5] Boozer A D, Boca A, Miller R, Northup T E and Kimble H J 2007 *Phys. Rev. Lett.* **98** 193601
[6] Fleischhauer M and Lukin M D 2000 *Phys. Rev. Lett.* **84** 5094
[7] Gorshkov A V, Calarco T, Lukin M D and Sørensen A S 2008 *Phys. Rev. A* **77** 043806
[8] Schraft D, Hain M, Lorenz N and Halfmann T 2016 *Phys. Rev. Lett.* **116** 073602
[9] Hsiao Y F, Tsai P J, Chen H S, Lin S X, Hung C C, Lee C H, Chen Y H, Chen Y F, Yu I A and Chen Y C 2018 *Phys. Rev. Lett.* **120** 183602
[10] Deng R J, Yang Z H and Jia X J 2017 *Acta Phys. Sin.* **66** 074201 (in Chinese)
[11] Dudin Y O, Li L and Kuzmich A 2013 *Phys. Rev. A* **87** 031801
[12] Katz O and Firstenberg O 2018 *Nat. Commun.* **9** 2074
[13] Qi Y R, Gao H and Zhang S G 2009 *Chin. Phys. Lett.* **26** 114211
[14] Heinze G, Hubrich C and Halfmann T 2013 *Phys. Rev. Lett.* **111** 033601
[15] Chen L R, Xu Z X, Li P, Wen Y F, Zeng W Q, Wu Y L, Tian L, Li S J and Wang H 2016 *Chin. Phys. B* **25** 024203
[16] Li W, Islam P and Windpassinger P 2020 *Phys. Rev. Lett.* **125** 150501
[17] Vernaz-Gris P, Huang K, Cao M, Sheremet A S and Laurat J 2018 *Nat. Commun.* **9** 363
[18] Wei Y C, Wu B H, Hsiao Y F, Tsai P J and Chen Y C 2020 *Phys. Rev. A* **102** 063720
[19] Peters T, Wang T P, Neumann A, Simeonov L S and Halfmann T 2020 *Opt. Express* **28** 5340
[20] Gouraud B, Maxein D, Nicolas A, Morin and Laurat J 2015 *Phys. Rev. Lett.* **114** 180503
[21] Sayrin C, Clausen C, Albrecht B, Schneeweiss P and Rauschenbeutel A 2015 *Optica* **2** 353
[22] Appel J, Figueroa E, Korystov D, Lobino M and Lvovsky A I 2008 *Phys. Rev. Lett.* **100** 093602
[23] Wang K, Zhang W, Zhou Z Y, Dong M X, Shi S, Liu S L, Ding D S and Shi B S 2017 *Chin. Opt. Lett.* **15** 060201
[24] Eisaman M, André, Massou F, Fleischhauer M, Zibrov A and Lukin M 2005 *Nature* **438** 837
[25] Nicolas A, Veissier L, Giner L, Giacobino E, Maxein D and Laurat J 2014 *Nat. Photon.* **8** 234
[26] Ding D S, Zhou Z Y, Shi B S and Guo G C 2013 *Nat. Commun.* **97** 062318
[27] Wang Y F, Li J F, Zhang S C, Su K Y, Zhou Y R, Liao K Y, Du S W, Yan H and Zhu S L 2019 *Nat. Photon.* **13** 346
[28] Gorshkov A V, André A, Fleischhauer M, Sørensen A S and Lukin M D 2007 *Phys. Rev. Lett.* **98** 123601
[29] Novikova I, Phillips N B and Gorshkov A V 2008 *Phys. Rev. A* **78** 021802
[30] Rowe M A, Kielpinski D, Meyer V, Sackett C A, Itano W M, Monroe C and Wineland D J 2001 *Nature* **409** 791
[31] Dong C H, Fiore V, Kuzyk M and Wang H 2012 *Science* **338** 1609
[32] James D F V and Kwiat P G 2002 *Phys. Rev. Lett.* **89** 183601
[33] Imamoğlu A 2002 *Phys. Rev. Lett.* **89** 163602
[34] Lin G W, Qi Y H, Lin X M, Niu Y P and Gong S Q 2015 *Phys. Rev. A* **92** 043842
[35] Walls D F and Milburn G J 1994 *Quantum Optics* (Berlin: Springer-Verlag)
[36] Tian L 2012 *Phys. Rev. Lett.* **108** 153604
[37] Wang Y D and Clerk A A 2012 *Phys. Rev. Lett.* **108** 153603
[38] Duan L M, Kuzmich A and Kimble H J 2003 *Phys. Rev. A* **67** 032305
[39] Duan L M and Kimble H J 2004 *Phys. Rev. Lett.* **92** 127902
[40] Simon J, Tanji H, Thompson J K and Vuletic V 2007 *Phys. Rev. Lett.* **98** 183601
[41] Wu H, Gea-Banacloche J and Xiao M 2008 *Phys. Rev. Lett.* **100** 173602