Cavity-assisted spontaneous emission of a single three-level atom as a single-photon source

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Stimulated Raman interaction of a classically pumped single three-level A-type atom in a resonator cavity featuring both radiative and unwanted losses is studied. It is shown that in the regime of stimulated adiabatic Raman passage the excited outgoing wave packet of the cavity-assisted electromagnetic field can be prepared in a one-photon Fock state with high efficiency. In this regime, the spatio-temporal shape of the wave packet does not depend on the interaction shape of the pump field, provided that the interaction of the atom with the pump is time-delayed with respect to the interaction of the atom with the cavity-assisted field. Therefore, the scheme can be used to generate a sequence of identical radiation pulses each of almost one-photon Fock state. It is further shown that the spatio-temporal shape of the outgoing wave packet can be manipulated by controlling the time of interaction between the atom and the cavity-assisted electromagnetic field.

INTRODUCTION

The interaction of a single atom with a quantized radiation-field mode in a high-Q cavity serves as a basic ingredient in various schemes in quantum information science (for a review see, e.g., Refs. [1, 2]). In this context, quantum control of single-photon emission from an atom in a cavity for generating one-photon Fock states on demand has been an essential prerequisite [3]. Single-photon sources operating on the basis of vacuum-simulated adiabatic passage with just a single atom in a high-Q optical cavity has been realized, where adjustment of the spatio-temporal shapes of the outgoing wave packets associated with the emitted photons has been achieved by means of driving laser pulses [4, 5].

More recently, we have discussed the possibility to use a single two-level atom in a high-Q cavity as a single-photon emitter, with the wave packet associated with the emitted photon being shorter than the cavity decay time and its spatio-temporal shape being time-symmetric [4, 6].

In the following we consider, within the frame of exact quantum electrodynamics, the interaction of a pumped three-level A-type atom with a realistic cavity-assisted electromagnetic field, with the aim of one-photon Fock-state emission. In particular, we study in detail both the properties of the excited outgoing wave packet and the quantum state it is prepared in. Further, the general explicit expression obtained for the spatio-temporal shape of the outgoing wave packet allows us to study different coupling regimes of atom–pump and atom–cavity-field interactions.

BASIC EQUATIONS

We consider an atom (position $\mathbf{r}_A$) that interacts with the electromagnetic field in the presence of a dispersing and absorbing dielectric medium. Applying the multipolar-coupling scheme in electric dipole approximation, we may write the Hamiltonian that governs the temporal evolution of the overall system, which consists of the electromagnetic field, the dielectric medium (including the dissipative degrees of freedom), and the atom coupled to the field, in the form of [7, 8]

$$
\hat{H} = \int d^3r \int_0^\infty d\omega \hbar \omega \hat{f}^d(r, \omega) \cdot \hat{f}(r, \omega) + \sum_k \hbar \omega_k S_{kk} - \hat{d}_A \cdot \hat{E}(\mathbf{r}_A).
$$

(1)

In this equation, the first term is the Hamiltonian of the field–medium system, where the fundamental bosonic fields $\hat{f}(r, \omega)$ and $\hat{f}^d(r, \omega)$,

$$
[\hat{f}_{\mu}(r, \omega), \hat{f}^d_{\nu}(r', \omega')] = \delta_{\mu\nu} \delta(\omega - \omega') \delta^{(3)}(r - r'),
$$

(2)

$$
[\hat{f}_{\mu}(r, \omega), \hat{f}_{\nu}(r', \omega')] = 0,
$$

(3)

play the role of the canonically conjugate system variables. The second term is the Hamiltonian of the atoms, where the $\hat{S}_{kk'}$ are the atomic flip operators for the atom,

$$
\hat{S}_{kk'} = |k'\rangle_A \langle k|,
$$

(4)

with the $|k\rangle_A$ being the energy eigenstates of the atom. Finally, the last term is the atom–field coupling energy, where

$$
\hat{d}_A = \sum_{kk'} \hat{d}_{kk'} \hat{S}_{kk'}
$$

(5)

is the electric dipole-moment operator of the atom ($\hat{d}_{kk'} = \langle k | \hat{d}_A | k' \rangle_A$), and the operator of the medium-assisted electric field $\hat{E}(\mathbf{r})$ can be expressed in terms of the variables $\hat{f}(r, \omega)$ and $\hat{f}^d(r, \omega)$ as follows:

$$
\hat{E}(\mathbf{r}) = \hat{E}^+(\mathbf{r}) + \hat{E}^-(\mathbf{r}),
$$

(6)

$$
\hat{E}^+(\mathbf{r}) = \int_0^\infty d\omega \hat{E}(\mathbf{r}, \omega), \quad \hat{E}^-(\mathbf{r}) = [\hat{E}^+(\mathbf{r})]^\dagger.
$$

(7)
where the classical (retarded) Green tensor \( G(r, r', \omega) \) is the solution to the equation

\[
\nabla \times \nabla \times G(r, r', \omega) - \frac{\omega^2}{c^2} \varepsilon(r, \omega) G(r, r', \omega) = \delta^{(3)}(r - r')
\]

and satisfies the boundary condition at infinity, i.e., \( G(r, r', \omega) \rightarrow 0 \) if \(|r - r'| \rightarrow \infty|\). In particular, we assume that the atomic transition frequencies \( \Delta \) and satisfies the boundary condition at infinity, i.e., \( G(r, r', \omega) \rightarrow 0 \) if \(|r - r'| \rightarrow \infty|\).

**PUMPED THREE-LEVEL \( \Lambda \)-TYPE ATOM IN A CAVITY**

Let us consider a three-level \( \Lambda \)-type atom (Fig. 1) in a cavity bounded by a perfectly reflecting mirror and a fractionally transparent (coupling) mirror and model the cavity by a one-dimensional dielectric layer system (see, e.g., Ref. [6]). In particular, we assume that the atomic transition \( |2 \rangle \rightarrow |3 \rangle \) is strongly coupled to the cavity field. Further, restricting our attention to (quasi-)resonant atom–field interaction, we may start from the (one-dimensional version of the) Hamiltonian (11) and apply the rotating-wave approximation to the atom–field interaction. Moreover, we assume that an external (classical) pump field with frequency \( \omega_p \) is applied to the \(|1 \rangle \rightarrow |2 \rangle \) transition of the atom. In this way we may write the Hamiltonian in the form of

\[
\hat{H} = \int dz \int_0^\infty d\omega \ h \omega \hat{f}(z, \omega) \hat{f}(z, \omega) + h \omega_0 \hat{S}_{22} \]  

\[ - g_\epsilon(t) \left[ d_{23} \hat{S}_{32} \hat{E}^*(z_A) + \text{H.c.} \right] \]  

\[ - \frac{\hbar}{2} \Omega_\epsilon(t) \left[ \hat{S}_{12} e^{i\omega_p t} + \text{H.c.} \right]. \]  

Here, \( \Omega_\epsilon(t) \) is the (real, time-dependent) Rabi frequency of the pump field, and the (real) function \( g_\epsilon(t) \) defines the (time-dependent) shape of the interaction of the atom with the cavity field, which can be realized by (quasistatic) motion of the atom through the cavity in the direction perpendicular to the cavity axis.

In what follows we assume that the atom is initially (at time \( t = 0 \)) prepared in the state \(|1 \rangle \) and the rest of the system, i.e., the part of the system that consists of the electromagnetic field and the cavity, is prepared in the ground state \(|\{0 \rangle \). In order to avoid lengthy formulae we assume that the frequencies of the atomic transitions \(|1 \rangle \leftrightarrow |2 \rangle \) and \(|3 \rangle \leftrightarrow |2 \rangle \) are equal to each other \((\omega_{21} = \omega_{23} = \omega_0)\). We may then expand the state vector of the overall system at later times \( t (t \geq 0) \) as

\[
|\psi(t)\rangle = C_1(t)|1\rangle|\{0\rangle\rangle + C_2(t) e^{-i\omega_p t}|2\rangle|\{0\rangle\rangle \]  

\[ + \int dz \int_0^\infty d\omega \ C_3(z, \omega, t) e^{-i\omega \cdot z} f(z, \omega)|\{0\rangle\rangle, \]

where \( f(z, \omega)|\{0\rangle\rangle \) is an excited single-quantum state of the combined field–cavity system. It is not difficult to prove that the Schrödinger equation for \(|\psi(t)\rangle\) leads to the following system of (in-tgro)-differential equations for the probability amplitudes \( C_1(t) \), \( C_2(t) \) and \( C_3(z, \omega, t) \):

\[ C_1' = \frac{i}{2} \Omega_\epsilon(t) e^{i\Delta_p t} C_2(t), \]  

\[ C_2' = \frac{i}{2} \Omega_\epsilon(t) e^{-i\Delta_p t} C_1(t) \]  

\[ - \frac{d_{23}}{\sqrt{\pi \hbar c A}} \int_0^\infty d\omega \ \frac{\omega^2}{c^2} \int dz \sqrt{\varepsilon''(z, \omega)} \]  

\[ \times G(z_A, z, \omega) C_3(z, \omega, t) e^{-i(\omega - \omega_0) t}, \]  

\[ C_3(z, \omega, t) = \frac{d_{23}^2}{\sqrt{\pi \hbar c A}} \frac{\omega^2}{c^2} \sqrt{\varepsilon''(z, \omega)} \]  

\[ \times G^*(z_A, z, \omega) C_2(t) e^{i(\omega - \omega_0) t}, \]

where \( A \) is the area of the coupling mirror of the cavity, and \( \Delta_p \) is the detuning of the pump frequency from the atomic transition frequency \((\Delta_p = \omega_p - \omega_0)\). The spectral response of the cavity field is known to be determined by the Green function \( G(z, z', \omega) \). For a sufficiently high-Q cavity, the excitation spectrum effectively turns into a quasi-discrete set of lines of mid-frequencies \( \omega_k \) and widths \( \Gamma_k \), according to the poles of the Green function at the complex frequencies

\[ \Omega_k = \omega_k - \frac{1}{2} i \Gamma_k, \]

FIG. 1. Scheme of relevant energy levels and transitions for the three-level \( \Lambda \)-type atom.
In this case, substituting the formal solutions to Eqs. (13) and (15) [with the initial condition $C_1(0) = 1, C_2(0) = 0$ and $C_3(z,ω,0) = 0$] into Eq. (14), we derive the integrodifferential equation

$$\dot{C}_2 = \frac{i}{2} \Omega_p(t) e^{-i\Delta_p t} + \int_0^t dt' K(t,t') C_2(t'),$$

where the kernel function $K(t,t')$ reads (for details, see Ref. [3])

$$K(t, t') = -\frac{1}{4} \Omega_p(t) \Omega_p(t') e^{-i\Delta_p (t-t')} - \frac{1}{4} \alpha_k \Omega_k g_c(t) g_c(t') e^{-i((\Omega_k - \omega_0)(t-t'))},$$

$$\alpha_k = \frac{4|d_{21}|^2}{\hbar \varepsilon_A |n_1(\Omega_k)|^2} \sin^2[\omega_k |n_1(\Omega_k)|z_A/c],$$

where $l$ is the length of the cavity, $n_1(\omega)$ is the (complex) refractive index of the medium inside the cavity, and $\omega_k$ is now the fixed quasi-discrete line frequency that is (quasi-)resonant to the atomic transition frequency.

Following Ref. [4], it can be shown that when the Hilbert space of the system is effectively spanned only by a single excitation, on a time-scale that is short compared to the inverse spontaneous emission rate of the atom the quantum state of the outgoing field is given by means of the multimode Wigner function

$$W_{\text{out}}(\alpha_i, t) = W_1(\alpha_1, t) \prod_{i \neq 1} W_i^{(0)}(\alpha_i, t),$$

where

$$W_1(\alpha, t) = [1 - \eta(t)] W_i^{(0)}(\alpha) + \eta(t) W_i^{(1)}(\alpha),$$

with $W_i^{(0)}(\alpha)$ and $W_i^{(1)}(\alpha)$ being the Wigner functions of the vacuum state and the one-photon Fock state, respectively, for the $i$th nonmonochromatic mode. As we see, the mode labeled by the subscript $i = 1$, i.e., the mode associated with the excited outgoing wave packet, is basically prepared in a mixed state of a one-photon Fock state and the vacuum state, due to unavoidable existence of unwanted losses. The other nonmonochromatic modes of the outgoing field with $i \neq 1$ are in the vacuum state and, therefore, remain unexcited. The Wigner function $W_1(\alpha, t)$ reveals that $\eta(t)$ can be regarded as being the efficiency to prepare the excited outgoing wave packet in a one-photon Fock state,

$$\eta(t) = \int_0^\infty d\omega |F(\omega, t)|^2 \simeq \int_{-\infty}^\infty d\omega |F(\omega, t)|^2,$$

where

$$F(\omega, t) = \frac{d_{21}}{\sqrt{\pi \hbar \varepsilon_A}} \sqrt{\frac{c}{\omega^2}} e^{\frac{2\Omega_0}{\hbar \varepsilon_A}} \int_0^t dt' G^* (0^+, \omega, \omega) C_2^2 (t') e^{i\omega(t-t')} e^{i\omega_0 t'}. \tag{24}$$

The excited outgoing wave packet is characterized by the mode function

$$F_1(\omega, t) = \frac{F(\omega, t)}{\sqrt{\eta(t)}}, \tag{25}$$

and its spatio-temporal shape reads

$$\phi_1(z, t) = \frac{\kappa_k}{2} \sqrt{\frac{\pi \hbar \omega_k}{\varepsilon_A \hbar c}} \int_0^t dt' g_c(t') C_2^2 (t') e^{-i((\Omega_k - \omega_0) t - \omega_0 t')} e^{i\omega_0 (t-z/c)} \times [\Theta(z+l) \Theta(-z) \Theta(t - t') + \Theta(z) \Theta(t - t' - z/c)], \tag{26}$$

(for details, see Ref. [3]). Here, the term with $\Theta(-z)$ corresponds to the part of the excited outgoing mode that is still inside the cavity, while the term with $\Theta(z)$ stands for the part which is already escaped from the cavity.

**RESULTS AND DISCUSSION**

Let us first consider the case when the atom interacts with the pump field and the cavity-assisted electromagnetic field in the "counter-intuitive" order, where the atom–cavity-field interaction is well-established before the pump field is turned on. Then, provided that the pump field grows sufficiently slowly, the evolution of the quantum state of the combined atom–cavity-field system closely follows the one of a dark state, i.e., the dressed state of a single-mode–atom system which has no contribution from the upper atomic state [2]. Thus, stimulated Raman adiabatic passage is realized, thereby a single photon being generated in the (excited) outgoing field.

For a numerical evaluation of the equations given above, we have modeled the (time-dependent) Rabi frequency of the pump field, $\Omega_p(t)$, and the (time-dependent) atom–cavity-field interaction shape function $g_c(t)$ by Gaussian functions (see Fig. 2).
Further, we have assumed that, apart from the atom, the cavity is empty (no medium inside the cavity).

The behavior of the absolute value of the spatio-temporal shape $\phi_1(z,t)$ of the excited outgoing wave packet is illustrated in Figs. 2 and 3 where the maximum single-mode vacuum Rabi frequency of the atom–cavity-field interaction is assumed to be $R_k = \sqrt{\alpha_k \omega_k} = 5 \times \Gamma_k$, and $(\Omega_p, \Delta_p, \omega_k - \omega_0) = (10, 10^{-3}, 10^{-3}) \times \Gamma_k$. The (numerical) results show that the spatio-temporal shape of the outgoing wave packet is determined by the cavity decay rate as well as the pump process, in particular by the delay between the interaction of the atom with the pump field and the cavity-assisted field, $\Omega_p(t)$ and $g_c(t)$ respectively. The efficiency of one-photon Fock state preparation $\eta(t)$ is an almost monotonically increasing function that reaches a constant value equal to the ratio of the cavity radiative decay rate to the total decay rate, which is assumed to be 90% in the following. Note that the cavity total decay rate is the sum of the decay rates due to radiative input–output coupling and due to unwanted losses such as scattering and medium absorption. In the same way, a sequence of light pulses each of almost one-photon Fock state can be generated by means of a beam of atoms that cross the cavity one at a time or by the same atom introducing recycling field pulses, which bring the atomic population to the initial level $|1\rangle$.

It should be pointed out that in the case of stimulated Raman adiabatic passage the spatio-temporal shape of the excited outgoing wave packet is independent of the shape of the pump field, as can be seen from a comparison of Fig. 2 with Fig. 3.

The interaction of a classically pumped single three-level atom in a cavity also allows the generation of an outgoing wave packet carrying a one-photon Fock state with given spatio-temporal shape, adjusted by means of the pump field shape. The control of the spatio-temporal shape of the outgoing wave packet can be achieved when the pump field is intense enough and very short compared to the cavity decay time. As illustrated in Figs. 4–6.
and the square-like pulse in Fig. 6 reveals the efficiency of one-photon Fock state preparation is low of stimulated adiabatic Raman passage is not realized, for a twin-peaked pump field shape and (10) reveals that the leading edge of the spatio-temporal shape of the excited outgoing wave packet precisely reproduces the shape of the pump field. However, since the regime of stimulated adiabatic Raman passage is not realized, the efficiency of one-photon Fock state preparation is low of the cavity field can be regarded as being constant on this time scale and is turned off before the pump field is turned off. For a twin-peaked pump field shape and \((R_k, \Omega_p, \Delta_p, \omega_k - \omega_0) = (30, 700, 0.001, 0.001) \times \Gamma_k\), Fig. 4 reveals that the leading edge of the spatio-temporal shape of the excited outgoing wave packet precisely reproduces the shape of the pump field. 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For the twin-peaked pulse in Fig. 5 and the square-like pulse in Fig. 6 respectively, efficiencies of \(\eta(t \to \infty) = 0.014\) and \(\eta(t \to \infty) = 0.022\) are observed. Note that with increasing intensity of the pump field, the oscillations of the spatio-temporal shape become very fast and may be therefore not resolvable. To conclude, we have shown that the interaction of a pumped three-level A-type atom with the electromagnetic field in a cavity can be used to generate one-photon Fock states in well-defined outgoing wave packets. In particular in the case when stimulated Raman adiabatic passage is realized, then the efficiency of one-photon Fock state generation may be close to 100%; it is only limited by the unavoidable unwanted losses in the system such scattering and medium absorption. In particular, for large times the efficiency equals to the ratio of the cavity radiative decay rate to the total decay rate. The spatio-temporal shape of the outgoing wave packet is robust against fluctuations of the shape of the atom–pump interaction, and, therefore, the scheme can be used to produce identical wave packets on demand carrying one-photon Fock states each. For sufficiently intense and short pumping, i.e., beyond the adiabatic regime, the scheme enables control of the spatio-temporal shape of the excited outgoing wave packet by means of the atom–
pump interaction shape. However in this case, the efficiency of one-photon Fock state generation is rather low in general.

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[1] H. Walther, Phys. 54, 617 (2006)
[2] H. Walther, B. T. H. Varcoe, B. G. Englert, and T. Becker, Rep. Prog. Phys. 69, 1325 (2006)
[3] C. Monroe, Nature 416, 238 (2002)
[4] A. Kuhn, M. Hennrich, and G. Rempe, Phys. Rev. Lett. 89, 067901 (2002)
[5] M. Keller, B. Lange, K. Hayasaka, W. Lange, and H. Walther, Nature 431, 1075 (2004)
[6] M. Khanbekyan, D.-G. Welsch, C. D. Fidio, and W. Vogel, Phys. Rev. A 78, 013822 (2008)
[7] C. D. Fidio, W. Vogel, M. Khanbekyan, and D.-G. Welsch, Phys. Rev. A 77, 043822 (2008)
[8] L. Knöll, S. Scheel, and D.-G. Welsch, “Coherence and statistics of photons and atoms,” (Wiley, New York, 2001) p. 1, quant-ph/0003121
[9] W. Vogel and D.-G. Welsch, Quantum Optics, 3rd ed. (Wiley-VCH, Weinheim, 2006)