Coherently triggered single photons from a quantum-dot cavity system

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(Received 14 June 2010; published 15 July 2010)

We present a scheme to realize a highly efficient solid state source of indistinguishable single photons using cavity-assisted stimulated adiabatic Raman passage in a single quantum dot. The Autler-Townes doublet, generated by using a resonant driving field between biexciton and exciton states is utilized to facilitate a two-photon Raman transition in the quantum dot. An optical field transient then coherently generates a single photon pulse, whose polarization is orthogonal to the polarization of the applied fields and can thus be filtered efficiently from the scattered fields. The triggered generated single photons have greater than 90% efficiency and more than 90% quantum indistinguishability using currently available experimental parameters.

DOI: 10.1103/PhysRevB.82.045308 PACS number(s): 42.50.Dv, 03.65.Ud, 03.67.Bg, 78.67.De

I. INTRODUCTION

A deterministic *on demand* source of single photons is a basic building block for linear quantum computation,[1] quantum cryptography,[2] quantum teleportation,[3] and quantum networks.[4] In all these applications, quantum interference between two single-photon pulses on a symmetric beam splitter has been exploited,[5] which imposes stringent requirement for the implemented single photons to be *indistinguishable* in all degrees of freedom, including their frequencies, spectral widths, pulse shapes, and polarizations. To generate single photons one requires a pumping mechanism to excite a *two-level emitter* and an efficient channeling of the subsequently emitted photons. The efficiency of the source can be enhanced by coupling the emitter to the waveguide[6] or a microcavity mode.[7] Single photon sources have been realized in various quantum systems, including single atom trapped in an optical cavity,[8] trapped ions,[9] molecules,[10] and quantum dots (QD).[11,12]

Semiconductor QDs are attractive for realizing quantum optical phenomena in solid-state systems, offering advantages such as integrability and scalability. As *artificial atoms*, QDs have discrete energy levels due to strong quantum confinement of the electron-hole pairs, and they can be embedded or grown with high precision in different semiconductor microcavities at desired spatial positions. The strong coupling regime, where the emitted photon become entangled with the cavity mode, has also been demonstrated.[13,14] However, QD-cavity-enabled single photon sources rely on *incoherent* pumping of excitons or electron-hole pairs.[15] Usually through incoherent pumping, the QD is excited in a quantum state far above from the desired exciton state, which relaxes quickly to the desired exciton state typically by phonon interactions. Thus the excited state has time uncertainty, resulting in so-called *timing jitter*, leading to a trade off between efficiency and indistinguishability. This is a *major* problem as the incoherent pumping of a QD cannot provide indistinguishable photons at high efficiency.[16] While there have been great achievements in improving the semiconductor cavity coupling and output efficiency of the emitted photons, there has been little work to address coherent *on-demand* loading. The coherent manipulation of energy levels in QDs has been a challenging task because of unavoidable photon scattering from the excitation optical source. In the last year, the coherent manipulation of exciton states in QDs has been demonstrated in a few remarkable experiments,[17–20] using, for example, pump-probe excitation and the Autler-Townes splitting of dressed-exciton states. Flagg *et al.*[21] and Ates *et al.*[22,23] have also demonstrated resonance fluorescence from a QD coupled to a semiconductor microcavity. In their experiment, the QD has been coherently driven by the external laser field and the emitted photons have been collected through the cavity mode, where the latter is geometrically separated from the excitation field.

The above optical control schemes form important breakthroughs in the context of realizing coherent quantum optics phenomena in an integrable and scalable semiconductor system. In a similar spirit, in this work, we introduce a coherent optical excitation scheme that exploits the stimulated Raman adiabatic passage (STIRAP) in a QD embedded in a semiconductor microcavity to coherently generate single photons. We show that the two-photon coherent Raman transitions in QDs[24] can be realized through the Autlet-Townes doublet. Somewhat remarkably, this scheme actually benefits from the natural anisotropic-exchange splitting that occurs between the x and y polarized excitons, which is usually a significant hindrance, e.g., for creating entangled photon pairs.[25] The proposed cavity-assisted STIRAP excitation for generating true single photons overcomes the main time jitter problems that are implicit with incoherent excitation.

Our article is organized as follows. In Sec. II, we introduce the theoretical framework used for modeling the QD-cavity system. A quantum master equation approach is used that allows us to compute the system state populations, as well as the first-order and second-order quantum correlation functions for the cavity emitted photons. Section III presents calculations of a representative cavity system and shows the ensuing populations of the exciton and photon states. Section IV discusses the indistinguishability of the emitted cavity photons and investigates the influence of pure dephasing. Finally, we conclude in Sec. V.

II. THEORETICAL APPROACH TO MODELING THE STIMULATED RAMAN ADIABATIC PASSAGE IN A QD-CAVITY SYSTEM

We consider a QD embedded in a semiconductor microcavity, where energy levels of the system is shown in Fig.

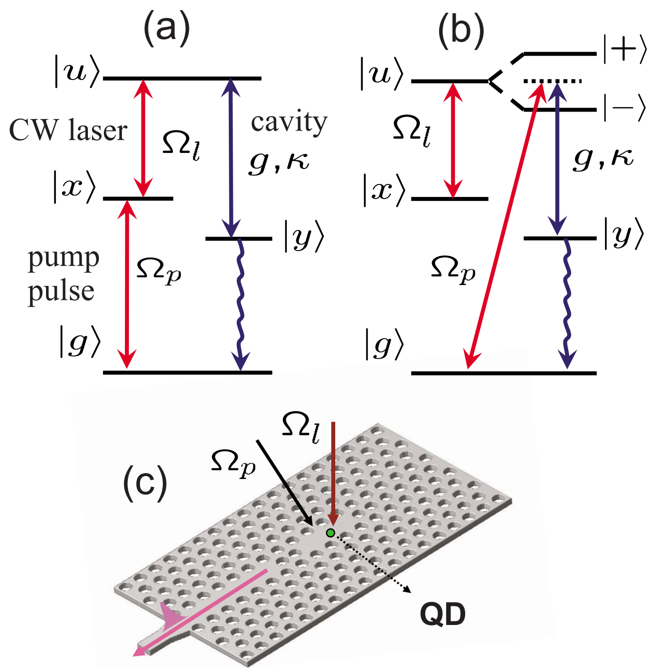

FIG. 1. (Color online) (a) Schematic of the single photon source exploiting the Raman flip process in a QD, using a CW-control laser and a pump pulse laser, and (b) the equivalent scheme in a dressed state picture. (c) Schematic of an example on-chip single photon source using a planar photonic crystal system (see Refs 26 and 27); other examples, e.g., could include the semiconductor micropillar systems (see Refs. 22 and 23).

1(a). The transitions from the biexciton state $|u\rangle$ to the exciton states $|x\rangle$ and $|y\rangle$ are coupled by a x-polarized laser field with Rabi frequency Ω_l, and a y-polarized cavity mode with vacuum Rabi coupling g, respectively. Because of the large biexciton binding energy in QDs, the laser field and the cavity mode effectively remain uncoupled with the transitions from the ground state $|g\rangle$ to the exciton states. The QD is optically pumped from its ground state $|g\rangle$ to the exciton state $|x\rangle$ by applying an x-polarized pump pulse, which has Rabi frequency $\Omega_p(t)$. The Hamiltonian of the system, in the rotating frame, can be written as

$$H=\hbar\Delta_p|x\rangle\langle x|+\hbar(\Delta_p+\Delta_l-\Delta_c)|y\rangle\langle y|-\hbar[\Omega_l|u\rangle\langle x|+\Omega_p(t)|x\rangle \times\langle g|+g|u\rangle\langle y|a+\text{H.c.}]+\hbar(\Delta_p+\Delta_l)|u\rangle\langle u|, \tag{1}$$

where Δ_p, Δ_l, and Δ_c are the detunings of the pump pulse, the laser field, and the cavity mode, and H.c. refers to Hermitian conjugate. For simulating the dynamical system, we perform master equation calculations in the density matrix representation:

$$\frac{\partial\rho}{\partial t}=-\frac{i}{\hbar}[H,\rho]-\frac{1}{2}\sum_{\mu}L_\mu^\dagger L_\mu\rho-2L_\mu\rho L_\mu^\dagger+\rho L_\mu^\dagger L_\mu, \tag{2}$$

where L_μ are the Lindblad operators, $\sqrt{\gamma_1}|x\rangle\langle u|$, $\sqrt{\gamma_1}|y\rangle\langle u|$, $\sqrt{\gamma_2}|g\rangle\langle x|$, and $\sqrt{\gamma_2}|g\rangle\langle y|$ correspond to the spontaneous decays and $\sqrt{2\gamma_d}|u\rangle\langle u|$, $\sqrt{\gamma_d}|x\rangle\langle x|$, and $\sqrt{\gamma_d}|y\rangle\langle y|$ correspond to the dephasing of biexciton and exciton states. The emission of the single photon pulses from the cavity mode is given by the Lindblad operator $\sqrt{\kappa}a$. Initially, the QD is in the ground state $|g\rangle$ and the cavity mode in the vacuum state.

We solve the optical Bloch equations using the density operator, Eq. (2), for $\langle\sigma_{ij}(t)\rangle$, where $\sigma_{ij}=|i\rangle\langle j|$, $i,j=g,x,u,y,Y,G$, with $|Y\rangle\equiv|y,1\rangle$ and $|G\rangle\equiv|g,1\rangle$. From the quantum regression theorem, for $\tau\geq 0$, the two-time correlations $\zeta_{ij}(t,\tau)=\langle a^\dagger(t)\sigma_{ij}(t+\tau)\rangle$ follow the same equation of motion as $\langle\sigma_{ij}\rangle$. We write the two-time correlations for different operators as following

$$\frac{\partial\zeta_{gG}}{\partial\tau}=-i\Omega_p(\tau)\zeta_{xG}-\frac{\kappa}{2}\zeta_{gG}, \tag{3}$$

$$\frac{\partial\zeta_{yY}}{\partial\tau}=ig\zeta_{yu}-\left(\gamma_2+\gamma_d+\frac{\kappa}{2}\right)\zeta_{yY}, \tag{4}$$

$$\frac{\partial\zeta_{xG}}{\partial\tau}=i\Delta_p\zeta_{xG}-i\Omega_l\zeta_{uG}-i\Omega_p^*(\tau)\zeta_{gG}-\frac{1}{2}(\kappa+\gamma_2+\gamma_d)\zeta_{xG}, \tag{5}$$

$$\frac{\partial\zeta_{yu}}{\partial\tau}=-i\Delta_c\zeta_{yu}+i\Omega_l\zeta_{yx}+ig\zeta_{yY}-\frac{1}{2}(\gamma_2+2\gamma_1+3\gamma_d)\zeta_{yu}, \tag{6}$$

$$\frac{\partial\zeta_{yx}}{\partial\tau}=i(\Delta_l-\Delta_c)\zeta_{yx}+i\Omega_l^*\zeta_{yu}+i\Omega_p(\tau)\zeta_{yg}-(\gamma_2+\gamma_d)\zeta_{yx}, \tag{7}$$

$$\frac{\partial\zeta_{uG}}{\partial\tau}=i(\Delta_p+\Delta_l)\zeta_{uG}-i\Omega_l^*\zeta_{xG}-ig\zeta_{YG} -\left(\gamma_1+\gamma_d+\frac{\kappa}{2}\right)\zeta_{uG}, \tag{8}$$

$$\frac{\partial\zeta_{yg}}{\partial\tau}=i(\Delta_p+\Delta_l-\Delta_c)\zeta_{yg}+i\Omega_p^*(\tau)\zeta_{yx}-\frac{1}{2}(\gamma_2+\gamma_d)\zeta_{yg}, \tag{9}$$

$$\frac{\partial\zeta_{YG}}{\partial\tau}=i(\Delta_p+\Delta_l-\Delta_c)\zeta_{YG}-ig\zeta_{uG}-\frac{1}{2}(\gamma_2+\gamma_d+2\kappa)\zeta_{YG}. \tag{10}$$

Equations (3)–(10) are solved with the initial conditions $\zeta_{yY}(t,0)=\langle\sigma_{YY}(t)\rangle$, $\zeta_{gG}(t,0)=\langle\sigma_{GG}(t)\rangle$, $\zeta_{yu}(t,0)=\langle\sigma_{Yu}(t)\rangle$, $\zeta_{gG}(t,0)=\langle\sigma_{GG}(t)\rangle$, and all other $\zeta_{ij}(t,0)=0$. The first-order correlation function is given by $g^{(1)}(t,\tau)\equiv\langle a^\dagger(t)a(t+\tau)\rangle=\zeta_{yY}(t,\tau)+\zeta_{gG}(t,\tau)$. Similarly, we calculate the second order correlation function $g^{(2)}(t,\tau)\equiv\langle a^\dagger(t)a^\dagger(t+\tau)a(t+\tau)a(t)\rangle$.

We have assumed a fixed dephasing rate due to electron-phonon coupling. Although there exists microscopic models for electron-phonon coupling in the weak coupling regime,[28] showing non-Markovian relaxation,[29] no such studies or formalisms exist for the strong coupling and pulsed coherent excitation regime. However, recent experiments at low temperatures, 5–50 K, using coherent excitation of QDs in cavities, have found that phenomenological exponential dephasing and exponential radiative decay fit very well to the data of Rabi oscillations and spectral line widths. Other complexities that occur for incoherent excitation, such as coupling to numerous background excitons and proposed shake-up processes,[30] are negligible for our coherent excitation scheme, a view that is also confirmed in recent experiments.[23]

In the presence of the resonant x-polarized laser field Ω_l, the biexciton state $|u\rangle$ and the exciton state $|x\rangle$ form the Autler-Townes doublet, $|\pm\rangle=(|u\rangle\pm|x\rangle)/\sqrt{2}$, shown in Fig. 1(b). The separation between the states $|\pm\rangle$ depends on Ω_l. The transitions from the states $|\pm\rangle$ to the ground state $|g\rangle$ and the exciton state $|y\rangle$ are dipole allowed, and the dipole couplings with the states $|g\rangle$ and $|y\rangle$ for each Autler-Townes state remain equal for the resonant laser field Ω_l, i.e., for $\Delta_l=0$. This is a necessary requirement for the complete population transfer in STIRAP through multiple intermediate states.[31] Due to the Autler-Townes splitting, the applied pump pulse and the cavity mode get detuned by $\Delta_p\pm\Omega_l/2$ and $\Delta_c\pm\Omega_l/2$, respectively. Therefore, for $\Delta_p=\Delta_c$ the pump pulse and cavity mode together satisfy the two-photon Raman resonance condition. When a slowly varying x-polarized pump pulse $\Omega_p(t)$ is applied in the presence of laser field Ω_l, the population in QD energy levels faithfully follows the cavity-assisted Raman adiabatic passage. In such conditions, the initial state of the cavity-QD system $|g,0\rangle$ is almost adiabatically transferred to the state $|y,1\rangle$, and the populations in the states $|x,0\rangle$ and $|u,0\rangle$ remain negligible throughout the pump pulse. From the state $|y,1\rangle$, a single y-polarized photon is emitted from the cavity mode with emission rate κ and the QD is left in the exciton state $|y\rangle$ which very slowly decays to the ground state $|g\rangle$ with spontaneous decay rate γ_2. After a time $\approx 6/\gamma_2$, when the entire population is returned back to ground state $|g\rangle$, the next pulse is applied.

Following the STIRAP process, in our scheme the population is further recycled back naturally and so no recycling pulses are necessary (cf. the case of trapped atoms[8]). However, if one wanted to reduce the time interval between generated single photon pulses, one could use a y-polarized π-pulse resonant with the $|y\rangle$ to $|g\rangle$ transition after the emission of the single photon. In a micropillar cavity, the emitted photons through the cavity mode are efficiently channeled in one direction using orthogonal excitation detection method.[22] For photonic crystals, as shown in Fig. 1(c), the emitted photon can be channeled with more than 80% efficiency using a coupled waveguide-cavity.[26,27] We remark that only the y-polarized transition satisfying the two-photon Raman resonance condition from the exciton states is possible, and thus, the presence of other background states, e.g., charged excitons, do not affect the evolution of the system. Indeed, the two-photon Raman resonant transitions have been successfully implemented in coherently driving two selected motional states of trapped ion within a manifold of motional states.[32] In what follows, we will consider the initial state of the QD as a neutral ground state. Although real dot systems may have charged states as the initial condition, we will not consider those in this work, though such states can be controlled by using an applied electric field.[33]

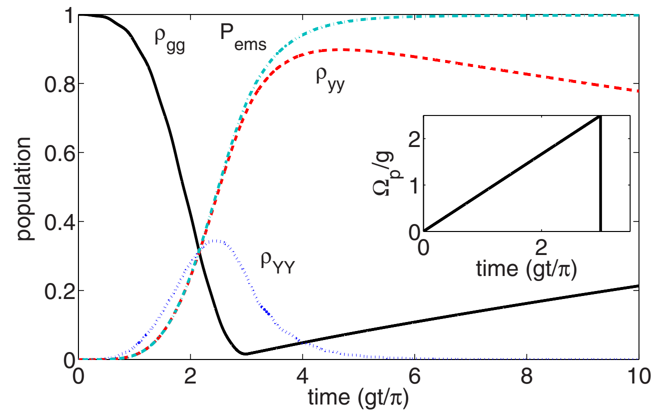

FIG. 2. (Color online) The populations of the states $|g,0\rangle$ (ρ_{gg}: black solid), $|y,0\rangle$ (ρ_{yy}: red dashed), and $|y,1\rangle$ (ρ_{YY}: blue dotted) of the QD-cavity system and the emission probability, P_{ems} (green chain), from the cavity mode for one pump pulse using parameters $\Delta_p=\Delta_l=\Delta_c=0$, $\Omega_l/g=5$, $\gamma_1/g=\gamma_2/g=\gamma_d/g=0.01$, and $\kappa/g=0.5$. The pump pulse is chosen as a sawtooth wave with maximum amplitude $\Omega_{\max}=2.5g$ and pulse widths $gt_p=3\pi$ (see inset).

III. POPULATION DYNAMICS

Next, we carry out calculations using the master equation model discussed in Sect. II. To maximize the STIRAP population transfer, we use a sawtooth wave pulse of pulse-width given by $gt_p=3\pi$ with peak amplitude $2.5g$, which is applied between ground state $|g\rangle$ to $|x\rangle$. In Fig. 2, we show the populations of the quantum states of the QD-cavity system, $\langle i|\rho|i\rangle$ for $|i\rangle=|g,0\rangle,|y,1\rangle$, and $|y,0\rangle$, calculated using Eq. (2). We introduce the following notation: ρ_{gg}—ground-state population, ρ_{YY}—y-polarized exciton and an excited cavity photon, ρ_{yy}—y-polarized exciton after emission of cavity photon; we also define $P_{\mathrm{ems}}(t)=\int_0^t\kappa\langle a^\dagger a(t')\rangle dt'$ as the probability of cavity photon emission at time t. The initial ground state population (ρ_{gg}) decreases and reaches a minimum during the pulse; however, it never becomes zero because of a small population returning back in the ground state from the spontaneous decay of the state $|y,0\rangle$ (ρ_{yy}). The population in the state $|y,1\rangle$ (ρ_{YY}), reaches the maximum before the pump pulse attains its maximum, which then decays to the state $|y,0\rangle$ after emitting a single photon pulse. A small population of the order of 10^{-3} is also generated in the state $|g,1\rangle$ through the spontaneous decay of the state $|y,1\rangle$. The populations in the state $|x,0\rangle$ and $|u,0\rangle$ remain of the order of 10^{-2} during the pulse, thus the evolution of the QD-cavity system efficiently follows the Raman adiabatic passage. The probability of the photon emission from the cavity mode during the pulse, P_{emiss}, reaches 1 when the population in the state $|y,1\rangle$ decays to zero (see Fig. 2).

As discussed above, because of the spontaneous decay of exciton state $|y,0\rangle$, there is a small continuous flow in the population of the ground state $|g,0\rangle$ during the pump pulse. If a large population is returned back in the ground state during the pump pulse, it could lead to the emission of more than one photon from the cavity mode in every pump pulse. To avoid the emission of more than one photon per pulse, the

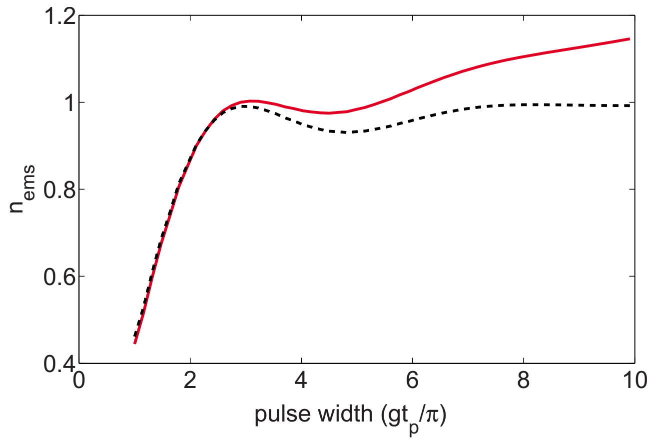

FIG. 3. (Color online) Pulse width dependence of the number of photons emitted from the cavity mode, n_{ems}, for one pump pulse of sawtooth wave shape. The red solid (black dashed) curve is for $\gamma_2/g=0.01$ ($\gamma_2/g=0$). The other parameters are the same as in Fig. 2.

pulse width should be chosen such that $\gamma_2 t_p < 0.1$, which comes from the fact that, for $\gamma_2=0$, the emission probability per pulse remains larger than 0.9 and the contribution of the population in ground state from the spontaneous emission is $\approx \gamma_2 t_p$. For larger peak amplitudes the required pulse widths for complete population transfer are smaller; however, to avoid populating state $|u,0\rangle$ and $|x,0\rangle$, the peak amplitude of the applied pump pulse should be smaller than the detunings from the upper states $|\pm\rangle$, i.e., $\Omega_p(t) < \Omega_l$. For QD embedded in a microcavity the off-resonant exciton have the spontaneous decay rate of around 0.1–1 μeV [13,14] ($10^{-2}g$ if $g=20$ μeV), thus we require the condition $\gamma_2 t_p < 0.1$ or $t_p < 4-40$ ns. We stress that all of these parameters are consistent with present day experiments.

In Fig. 3, we show the number of photons emitted in a pump pulse with increasing pulse widths. The lower red curve is for $\gamma_2=0$ and the upper black curve is for $\gamma_2/g = 0.01$. Following the complete transfer of ground state population in $|y,1\rangle$, and subsequent emission of a single photon pulse, the population in the state $|y,0\rangle$ becomes large and starts contributing to the population in ground state. If the pulse width is larger than the required value for complete population transfer in STIRAP, the population transferred to the ground state in spontaneous emission of $|y,0\rangle$ can contribute to the photon emission per pulse and the number of photons emitted per pulse rises on increasing pulse width. After complete population transfer in STIRAP, for the two-photon Rabi coupling larger than the cavity mode decay rate, i.e., $g\Omega_p/\Omega_l \geq \kappa$, a small population from $|y,1\rangle$ reflects back to the state $|x,0\rangle$ and the emitted number of photons per pulse remains nearly one for a long range of pulse widths (even after a small feedback in the ground-state population from the spontaneous decay of the state $|y,0\rangle$). This can also provide more flexibility in choosing pump pulse widths. However, for $g\Omega_p/\Omega_l < \kappa$, the number of photons emitted per pulse becomes larger than one for increasing pulse widths (namely, after the complete population transfer in the STIRAP process).

IV. INDISTINGUISHABILITY OF THE GENERATED PHOTONS

Next, we analyze the indistinguishability of the generated photons, which is a key figure of merit for quantum information protocols. The indistinguishability of the photons is measured by sending two consecutive generated photon pulses through a beam splitter and detecting Hong-Ou-Mandel type correlation.[34] For perfectly indistinguishable photons, the probability of coincidence detection at the output of a symmetric beam-splitter remains zero because of the Bose-Einstein statistics. In earlier experiments[12] for measuring indistinguishability, the photons are allowed to pass through a Michelson interferometer or through a Mach-Zhender interferometer having path difference between two arms corresponding to a time delay between photons, which equals the time difference between two photon pulses generated from the single photon source. In such a case, the two photons are incident on a symmetric beam splitter at the same time. After passing through the interferometer the probability of the coincidence detection of one photon at each output port can be expressed in terms of the correlations of the cavity field operator, [16]

$$P_c = \frac{1}{2}\left[1 + \frac{\int_0^T dt \int_0^{T-t} d\tau [g^{(2)}(t,\tau) - |g^{(1)}(t,\tau)|^2]}{\int_0^T dt \int_0^{T-t} \langle a^\dagger(t)a(t)\rangle\langle a^\dagger(t+\tau)a(t+\tau)\rangle}\right], \tag{11}$$

where T is the time interval between two pump pulses, $g^{(1)}(t,\tau)=\langle a^\dagger(t)a(t+\tau)\rangle$ and $g^{(2)}(t,\tau)=\langle a^\dagger(t)a^\dagger(t+\tau)a(t+\tau)a(t)\rangle$ are the un-normalized first-order and second-order quantum correlation functions of the cavity field. The first-order correlation contains the interference effects between two photons and the second-order correlation contains the

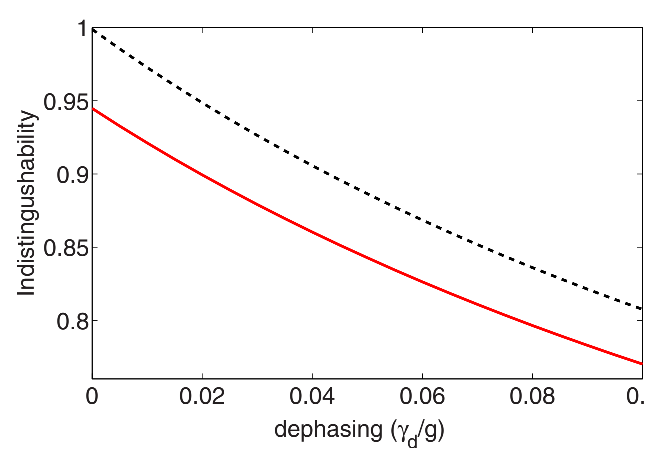

FIG. 4. (Color online) The dependence of the indistinguishability of generated photon pulses on dephasing rate γ_d. The red solid (black) curve is for $\gamma_2/g=0.01$ ($\gamma_2=0.0$), the maximum pulse amplitude $2.5g$ and pulse width $gt_p=3\pi$, respectively. The other parameters are the same as in Fig. 2.

probability of generating more than one photon per pulse from the source.

We calculate the different correlations required in Eq. (11) by solving optical Bloch equations for one time correlations from Eqs. (2)–(10) and using the quantum regression theorem. The indistinguishability of the photons is given by $1-P_c$. In Fig. 4, we show the dependence of the indistinguishability of the generated photons on dephasing of the biexciton and exciton states. The indistinguishability of the photons decreases on increasing dephasing rate. The decrease in indistinguishability is caused by information about the photons being attached to the phonon baths. In our scheme, there are phonon baths that couple to the biexciton ($|u\rangle$) and exciton state ($|y\rangle$), and there are thermal baths that lead to spontaneous emission from $|u\rangle$ to $|y\rangle$ and $|y\rangle$ to $|g\rangle$. In addition, the spontaneous decay of the exciton $|y\rangle$ also contributes to a small probability of generating more than one photon per pulse which adversely affects the indistinguishability of the emitted photons. However, the decay rate γ_1 does not affect the indistinguishability much and only a decrease of the order of 10^{-3} occurs by changing γ_1 from 0 to 0.1.

V. CONCLUSIONS

We have proposed and analyzed a cavity-assisted QD STIRAP scheme that shows good promise as a efficient solid state source of indistinguishable single photons. The two-photon Raman transition in the QD can be realized through the Autler-Townes doublet generated by using a resonant laser field between the biexciton and exciton states. The generated single photon pulses has a linear polarization orthogonal to the driving field, which makes it distinguishable from the background and well suited as a polarization encoded qubit and for pulse shaping.[35] We predict that single photons can be generated with greater than 90% efficiency and more than 90% indistinguishability, for $\gamma_d < 0.02g$, in the Hong-Ou-Mandel interference experiment.

ACKNOWLEDGMENTS

This work was supported by the National Sciences and Engineering Research Council of Canada, and the Canadian Foundation for Innovation. We acknowledge useful discussions with Robin Williams of the National Research Council's Institute for Microstructural Sciences, Canada.

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