Transient behaviors of photon switching by quantum interference

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(Received 24 November 2003; published 2 June 2004)

We have experimentally studied the transient behaviors of photon switching by quantum interference in laser-cooled ^{87}Rb atoms. The experimental results are consistent with theoretical predictions. We show that the rise time of the probe absorption cannot be less than $2/\Gamma$ where Γ is the spontaneous decay rate of the excited states. In the low-intensity limit, the rise time is equal to $2\Gamma/(\Omega_c^2+\Omega_s^2)$ for the scheme used here, where Ω_c and Ω_s are the Rabi frequencies of the coupling and switching fields. A simple picture based on the dark and nondark states is provided to explain these transient behaviors.

DOI: 10.1103/PhysRevA.69.063801 PACS number(s): 42.50.Gy, 42.62.Fi, 32.80.Pj

The enhancement of nonlinearity in optical media by quantum interference has attracted much attention [1–8]. This phenomenon has potential applications in the manipulation of quantum information; for example, electromagnetically induced transparency (EIT) could possibly be exploited in the storage of quantum information and in quantum memory [9–11], and the cross-phase modulation (XPM) scheme, which exhibits a giant Kerr nonlinearity along with vanishing linear susceptibilities [12,13], could potentially be employed in quantum logic gates, quantum nondemolition measurements, and quantum teleportation [14–16]. Recently, Harris and Yamamoto described a four-state atomic system based on the EIT scheme that inhibits linear one-photon absorption but enhances nonlinear two-photon absorption [7]. The vanishing linear absorption and the large enhancement of nonlinear absorption achievable using this system has prompted the consideration of nonlinear optical processes at low-light levels [17]. In Ref. [7], Harris and Yamamoto obtained the important result that one switching photon is sufficient to cause one probe photon to be absorbed under ideal conditions in this four-state system. This result leads to the very interesting idea of the threefold entangled state, which has potential applications in the manipulation of quantum information [18–20]. Yan *et al.* [21] and Braje *et al.* [22] have observed the phenomenon of photon switching by quantum interference in a four-state atomic system and have demonstrated switching of the nonlinear absorption with laser pulses; however, no researcher has yet achieved the single-photon switching predicted by Harris and Yamamoto.

In the four-state system of photon switching by quantum interference, a weak probe field and a strong coupling field form the three-state Λ-type configuration of EIT; in this system, a switching field drives the transition between the ground state of the coupling field and another excited state. Without the switching field, the probe absorption is inhibited or suppressed due to the EIT effect. The presence of the switching field enables the absorption of the probe field and induces the three-photon transition from the ground state of the probe field to the excited state of the switching field. As shown in Ref. [7], a switching pulse containing one photon is able to reduce the transmission of a probe pulse to $1/e$ under ideal conditions. This single-photon switching is essential to the realization of the threefold entangled state. The calculation presented in Ref. [7] is based on the steady-state solutions of the absorption, phase shift, and group velocity delay time of the probe field. The aim of the present study was to further elucidate the transient behaviors of the probe absorption and rise time of the switching event.

In this study, we examined the transient behavior of photon switching by quantum interference in cold ^{87}Rb atoms. The cold atoms were produced by a vapor-cell magneto-optical trap (MOT). The MOT used in these experiments is described in detail elsewhere [23,24]. Typically, we trap 5×10^7 atoms with a temperature of about 250 μK in the MOT. All the laser and magnetic fields of the MOT are turned off during the measurement of the probe absorption. The coupling and probe fields both drive the $|5S_{1/2}, F=2\rangle \rightarrow |5P_{3/2}, F'=2\rangle$ transition resonantly. They are circularly polarized with right (σ_+ polarization) and left (σ_- polarization) helicities, respectively. The two laser beams propagate in almost the same direction. The switching field drives the $|5S_{1/2}, F=2\rangle \rightarrow |5P_{3/2}, F'=1\rangle$ transition resonantly. This field is linearly polarized in the direction parallel to the propagation direction of the coupling and probe fields. The probe, coupling, and switching fields come from three laser diodes that are injection-locked by the same external-cavity diode laser. Their spectral widths are irrelevant in this study. The beam widths of the coupling and switching fields are larger than the cold atom cloud. The probe beam has a width of 1 mm and is completely inside the coupling beam. Each of the three beams passes through an acousto-optic modulator (AOM) before interacting with the atoms. A laser beam of spectral width 8Γ serves as a repumping field to excite the population in the $|5S_{1/2}, F=1\rangle$ state to the $|5P_{3/2}, F'=1\rangle$ state, where $\Gamma=2\pi\times5.9$ MHz is the spontaneous decay rate of the $|5P_{3/2}\rangle$ excited states. This beam has a width of 10 mm and a power of 10 mW. The repumping field is linearly polarized in the direction parallel to the propagation direction of the coupling and probe fields. The scheme of the relevant energy levels and laser excitations is depicted in Fig. 1(a).

Since the probe field is much weaker than the coupling field, most of the population in $|5S_{1/2}, F=2\rangle$ is optically pumped to the $|m=2\rangle$ Zeeman state. Only the $|5S_{1/2}, F=2, m=2\rangle \rightarrow |5P_{3/2}, F'=2, m=1\rangle$ transition contributes sig-

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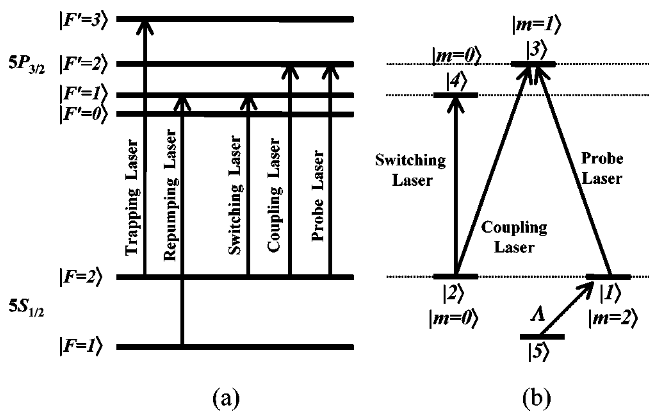

FIG. 1. (a) Relevant energy levels of ^{87}Rb atoms and laser excitations in the experiment. (b) The five-state system considered in the theoretical calculation.

nificantly to the observation of the probe absorption. Therefore, we have a four-state system comprising the $|5S_{1/2}, F=2, m=2\rangle$, $|5S_{1/2}, F=2, m=0\rangle$, $|5P_{3/2}, F=2, m=1\rangle$, and $|5P_{3/2}, F=1, m=0\rangle$ states (denoted $|1\rangle$, $|2\rangle$, $|3\rangle$, and $|4\rangle$, respectively), as shown in Fig. 1(b). This four-state system is not a closed system; we treat all states other than the above four as the fifth state ($|5\rangle$) and consider that the population of this fifth state is returned to the $|1\rangle$ state via an incoherent pump.

The timing sequence of the probe absorption measurement is described below. We first turn off the magnetic field of the MOT and then, after 1.5-ms delay, we shut off the trapping beams and switch on the coupling field at $t=0$. This 1.5-ms delay is sufficient to greatly reduce the influence of the magnetic field of the MOT. We then switch on the probe field at $t=50$ μs and the switching field at $t=300$ μs. The transient behavior of the switching field conforms to an exponential function with a time constant of about 20 ns. A 125-MHz photodiode (New Focus 1801) detects the probe transmission and its output is directly sent to a digital oscilloscope. After the measurement is complete, the coupling, probe, and switching fields are turned off and the MOT is turned back on. The repumping field is present throughout the procedure. The above sequence is repeated with a period of 100 ms; data are averaged 256 times by the oscilloscope before being transferred to the computer.

We also measured the steady-state spectra of the probe absorption. The repumping beam is retroreflected in the spectroscopic measurement to prevent the population from being trapped in the $|5S_{1/2}, F=1, m=0\rangle$ state. It is linearly polarized in the direction normal to the propagation direction of the coupling and probe fields by inserting a quarter-wave plate in its optical path. This reflected repumping beam does not appear in the measurement of the transient probe absorption. Figure 2(a) shows a typical EIT spectrum. The probe frequency was swept at a rate of $2.5\Gamma/\text{ms}$. This sweep rate was slow enough that the depth of the narrow transparency window could be accurately measured. Since the atom cloud gradually expands in the absence of the trap or the cold atoms get pushed away by the laser fields during the spectroscopic measurement, the spectrum is slightly asymmetric. From the separation of the two absorption peaks in the spectrum, the Rabi frequency Ω_c of the coupling field can be directly determined. This value was consistent with the value estimated from the laser power and beam width. At $\Omega_c=0.24\Gamma$, the absorption in the center of the transparency window is 10% of the maximum absorption, while at $\Omega_c=0.16\Gamma$, it is 20%. Comparing these observations with theoretical predictions, we estimate that the relaxation rate of the ground-state coherence is around 0.002Γ in our system. Figure 2(b) shows a typical probe absorption spectrum in the presence of both the coupling and switching fields. When recording this spectrum, the probe frequency was swept at a rate of $16\Gamma/\text{ms}$; this faster sweep rate makes the measured signal larger and prevents the spectrum from being distorted by the decay of the atomic density.

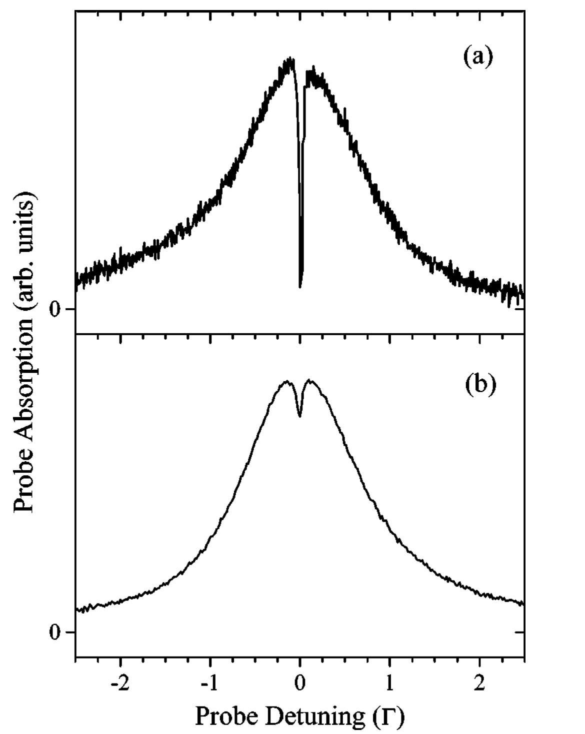

FIG. 2. The spectra of the probe absorption. The horizontal axis is in units of the spontaneous decay rate of the excited states. $\Omega_p=0.06\Gamma$, $\Omega_c=0.24\Gamma$, and $\Omega_s=0$ in (a) and $\Omega_s=0.40\Gamma$ in (b).

Figure 3 shows the measured and predicted behavior of the transient probe absorption versus time at four Rabi frequencies of the coupling and switching fields. The Rabi frequency Ω_s of the switching field was determined from the laser power and beam width. In these measurements, the system was initially in the EIT condition. After the switching field is suddenly turned on, the probe absorption first increases at a rate that depends on Ω_c and Ω_s. In Figs. 3(a)–3(c), the probe absorption exhibits oscillations rather than a simple exponential rise and the frequency of these oscillations increases with increasing Rabi frequency. However, when both Ω_c and Ω_s are small [Fig. 3(d)], the initial increase of the probe absorption is just an exponential function. The behavior described above is due to the photon switching effect of the four-state system. After the initial rapid increase, the probe absorption then slowly decreases to a steady-state value. If we turned off the repumping field

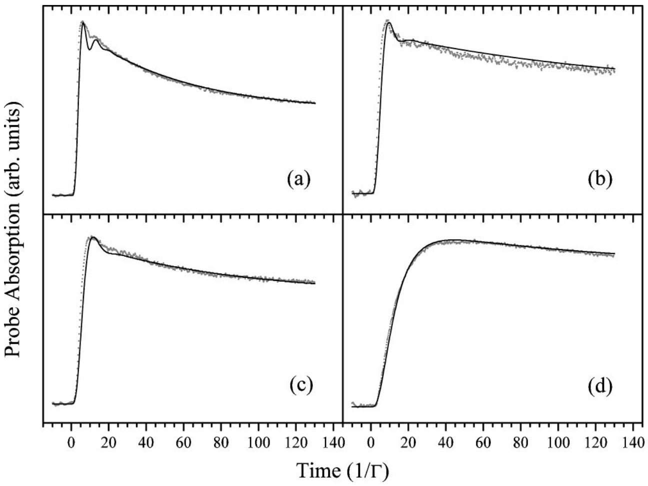

FIG. 3. The probe absorption as a function of time at four different Rabi frequencies of the coupling and switching fields. Gray circles represent experimental data and black lines represent theoretical predictions. The horizontal axis is in units of the excited-state lifetime τ (27 ns in this experiment). The values listed below are in units of the spontaneous decay rate Γ. $\Omega_p=0.12$ in all cases. (a) $\Omega_c=1.0$, $\Omega_s=1.5$, and $\Lambda=0.0090$. (b) $\Omega_c=1.0$, $\Omega_s=0.50$, and $\Lambda=0.0045$. (c) $\Omega_c=0.75$, $\Omega_s=0.50$, and $\Lambda=0.0080$. (d) $\Omega_c=0.35$, $\Omega_s=0.20$, and $\Lambda=0.0067$.

during the probe absorption measurement, the decay became more severe. This decay after the initial rapid increase is due to population loss from the four-state system, which occurs because this system is not closed. The time constants of the rapid increase and slow decay in the probe absorption curve differ considerably. We determined the rise time of the photon switching effect by measuring the $1/e$ time constant of the initial rise. The rise times of the experimental data in Figs. 3(a)–3(d) are 3.3τ, 4.5τ, 5.3τ, and 13τ, respectively, where $\tau=27$ ns is the excited-state lifetime.

We calculate the probe absorption as a function of time by solving the optical Bloch equation of the density-matrix operator of the five-state system in Fig. 1(b). The equation is given by

$$\frac{d\rho}{dt}=i\left[-\frac{H_0}{\hbar}+\left(\frac{\Omega_p}{2}e^{i\omega_{31}t}|1\rangle\langle 3|+\frac{\Omega_c}{2}e^{i\omega_{32}t}|2\rangle\langle 3|+\frac{\Omega_s}{2}e^{i\omega_{42}t}|2\rangle \times\langle 4|+\text{H.c.}\right),\rho\right]+\left\{\frac{d\rho}{dt}\right\}, \quad (1)$$

where H_0 is the atomic Hamiltonian; Ω_p, Ω_c, and Ω_s are the Rabi frequencies of the probe, coupling, and switching fields, respectively; and ω_{31}, ω_{32}, and ω_{42} are the resonance frequencies of the $|1\rangle\to|3\rangle$, $|2\rangle\to|3\rangle$, and $|2\rangle\to|4\rangle$ transitions. The term $\{d\rho/dt\}$ describes the relaxation of ρ; its matrix elements are as follows:

$$\left\{\frac{d\rho_{11}}{dt}\right\}=C_{31}^2\Gamma\rho_{33}+\Lambda\rho_{55},$$

$$\left\{\frac{d\rho_{22}}{dt}\right\}=C_{32}^2\Gamma\rho_{33}+C_{42}^2\Gamma\rho_{44},$$

$$\left\{\frac{d\rho_{12}}{dt}\right\}=-\gamma\rho_{12},$$

$$\left\{\frac{d\rho_{ij}}{dt}\right\}_{i=1,2;j=3,4}=-\frac{\Gamma}{2}\rho_{ij},$$

$$\left\{\frac{d\rho_{jj\prime}}{dt}\right\}_{j,j\prime=3,4}=-\Gamma\rho_{jj\prime},$$

$$\left\{\frac{d\rho_{55}}{dt}\right\}=(1-C_{31}^2-C_{32}^2)\Gamma\rho_{33}+(1-C_{42}^2)\Gamma\rho_{44}-\Lambda\rho_{55},$$

where Λ is the incoherent pump rate that moves population from $|5\rangle$ to $|1\rangle$, Γ is the spontaneous decay rate of the $|3\rangle$ and $|4\rangle$ excited states, γ is the relaxation rate of the ground-state coherence, andC_{31}, C_{32}, and C_{42} are the Clebsch-Gordan coefficients with respect to the electric-dipole-moment matrix elements of $\langle 3|$ and $|1\rangle$, $\langle 3|$ and $|2\rangle$, and $\langle 4|$ and $|2\rangle$, respectively. The equation of the density-matrix operator was solved numerically. The imaginary part of ρ_{31} determines the absorption cross section of the probe field. In the calculation, we used the experimental conditions of Ω_p, Ω_c, and Ω_s and adjusted Λ as a free parameter. Varying Λ in the magnitude over the range less than 0.01Γ only affects the behavior of the slow decay in the absorption curve and does not change the transient behavior due to the photon switching effect. We set $\gamma=0.002\Gamma$ to correspond with our experimental system; this term is so small that it has little influence in the calculation and can be neglected. In the experiments, the switching field is switched on exponentially with a time constant of 0.74τ; hence this experimental characteristic was incorporated into the calculations. The results of the theoretical calculation are plotted in Fig. 3 as black lines. The experimental data and theoretical predictions show satisfactory agreement.

To gain insight into the transient behaviors due to the photon switching effect, we consider a closed four-state system in the calculation and treat the probe field as a perturbation. Additionally, to simplify the problem, we assume that the switching field is switched on as a step function and that γ is negligible. Using Eq. (1), we carry out the calculation to the first order. The optical coherence with respect to the probe absorption obeys the following equation:

$$\frac{d^3}{dt^3}\sigma_{31}+\Gamma\frac{d^2}{dt^2}\sigma_{31}+\frac{\Gamma^2+\Omega_c^2+\Omega_s^2}{4}\frac{d}{dt}\sigma_{31}+\frac{\Gamma(\Omega_c^2+\Omega_s^2)}{8}\sigma_{31} -\frac{i\Omega_p\Omega_s^2}{8}=0, \quad (2)$$

where σ_{31} is the amplitude or slowly varying part of ρ_{31}, which is defined as $\rho_{31}=\sigma_{31}\mathrm{e}^{-i\omega_{31}t}$. The steady-state probe absorption is obtained by dropping all time-derivative terms in the above equation, giving

$$\sigma_{31}^{ss}=i\frac{\Omega_s^2}{\Omega_c^2+\Omega_s^2}\frac{\Omega_p}{\Gamma}. \quad (3)$$

The medium is completely transparent for the probe field at $\Omega_s=0$. The presence of the switching field makes the me-

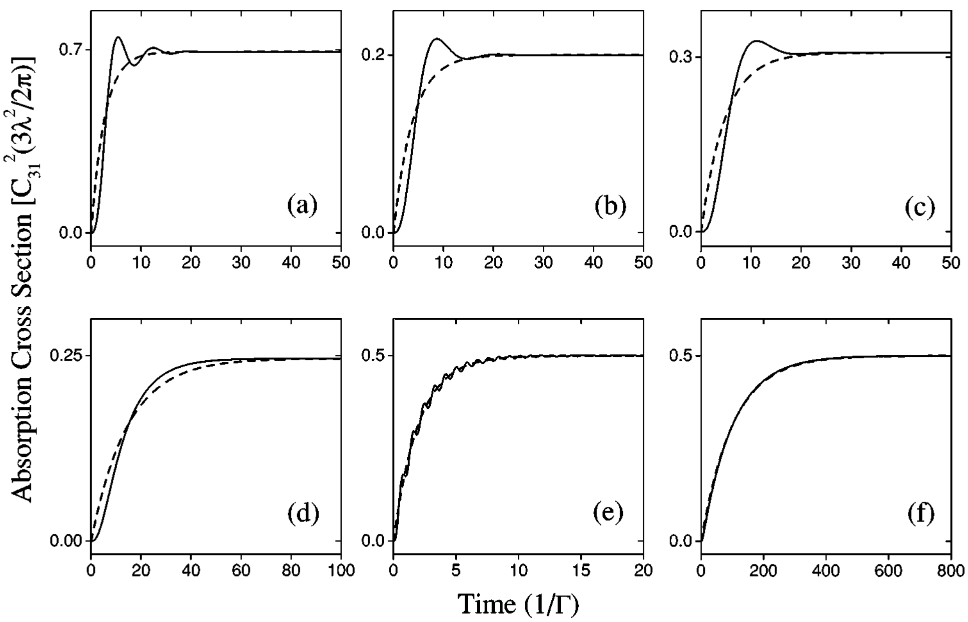

FIG. 4. The solutions of Eq. (2) (shown as solid lines). They are fitted with an exponential function of time constant τ_r. Dashed lines are the best fits. The values of Ω_c and Ω_s given below are in units of the spontaneous decay rate and τ_r is in units of the excited-state lifetime. (a) $\Omega_c=1.0$, $\Omega_s=1.5$, and $\tau_r=2.7$. (b) $\Omega_c=1.0$, $\Omega_s=0.50$, and $\tau_r=3.8$. (c) $\Omega_c=0.75$, $\Omega_s=0.50$, and $\tau_r=4.8$. (d) $\Omega_c=0.35$, $\Omega_s=0.20$, and $\tau_r=15$. (e) $\Omega_c=10$, $\Omega_s=10$, and $\tau_r=2.0$. (f) $\Omega_c=0.1$, $\Omega_s=0.1$, and $\tau_r=100$.

dium opaque, with the optical density depending on the ratio Ω_s/Ω_c. In the low-intensity limit—i.e., $\Omega_c, \Omega_s \ll \Gamma$—we expect the rate of change of σ_{31} to also be much smaller than Γ. Then, Eq. (2) becomes

$$\frac{d}{dt}\sigma_{31} = -\frac{\Omega_c^2+\Omega_s^2}{2\Gamma}(\sigma_{31}-\sigma_{31}^{ss}). \tag{4}$$

The transient probe absorption increases exponentially to the steady state with a time constant of $2\Gamma/(\Omega_c^2+\Omega_s^2)$. In the high-intensity limit—i.e., $\Omega_c, \Omega_s \gg \Gamma$—Eq. (2) becomes

$$\frac{d}{dt}\sigma_{31} = -\frac{\Gamma}{2}(\sigma_{31}-\sigma_{31}^{ss}). \tag{5}$$

This equation indicates that the rise time of the probe absorption is never less than $2/\Gamma$. We calculate the probe absorption as a function of time for a given Ω_c and Ω_s. The solution of Eq. (2) is fitted with an exponential function, and the rise time of the probe absorption is determined as the time constant of the best fit. Examples of the solutions of Eq. (2) and their best fits are shown in Fig. 4. The rise times of the probe absorption curves in Figs. 4(a)–4(d) are consistent with the experimental observations, and the rise times in Figs. 4(e) and 4(f) exactly coincide with the high- and low-intensity limits of $2/\Gamma$ and $2\Gamma/(\Omega_c^2+\Omega_s^2)$, respectively. In addition, we note that the rise time is unaffected by swapping the values of Ω_c and Ω_s; this is because all the coefficients of the time derivatives in Eq. (2) depend only on Ω_{sc}, which is defined as $\sqrt{\Omega_c^2+\Omega_s^2}$. The solid line in Fig. 5 shows the rise time of the probe absorption as a function of Ω_{sc}.

We further explain the transient behavior of the photon switching by quantum interference based on the dark and nondark states. In the interaction picture, the excited states can be rearranged as

$$|D\rangle = \frac{\Omega_c}{\Omega_{sc}}|4\rangle - \frac{\Omega_s}{\Omega_{sc}}|3\rangle, \tag{6}$$

$$|N\rangle = \frac{\Omega_s}{\Omega_{sc}}|4\rangle + \frac{\Omega_c}{\Omega_{sc}}|3\rangle, \tag{7}$$

$$\Omega_{sc} \equiv \sqrt{\Omega_c^2+\Omega_s^2}, \tag{8}$$

where $|D\rangle$ and $|N\rangle$ are the dark and nondark states. The coupling and switching fields do not interact with the dark state—i.e., $\langle 2|H_c+H_s|D\rangle=0$, where H_c and H_s are the coupling and switching Hamiltonians. The transition diagram is that depicted in Fig. 6(a). Because the probe field is considered as a perturbation, the four-state system can be decomposed into two subsystems, as shown in Fig. 6(b). One subsystem has two states $|1\rangle$ and $|D\rangle$ coupled only by the probe field. The absorption cross section is given by

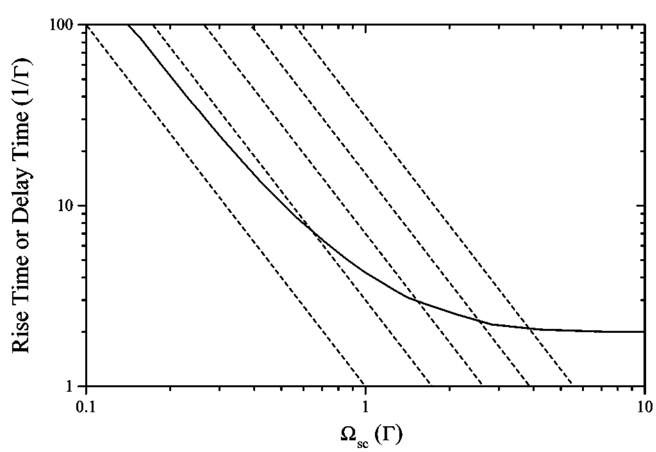

FIG. 5. The rise time of the probe absorption and the delay time of the probe propagation versus Ω_{sc}. Solid line represents the rise time of the probe absorption. Dashed lines represent the propagation delay time at several optical densities (X_{OD}). From left to right, the dashed lines correspond to $X_{OD}=2$, 4, 8, 16, and 32.

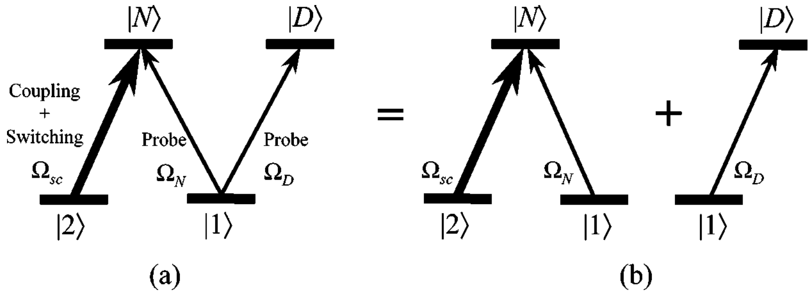

FIG. 6. (a) Transition diagram of laser fields in the basis of dark and nondark states. (b) The two subsystems equivalent to the four-state system in (a).

$$\frac{3\lambda^2}{2\pi}C_{31}^2\left(\frac{\Omega_s}{\Omega_{sc}}\right)^2(1-\mathrm{e}^{-(\Gamma/2)t}).$$

The transient probe absorption in the two-state subsystem increases exponentially at the rate of $\Gamma/2$. The other subsystem is in the Λ-type EIT configuration formed by $|1\rangle$, $|2\rangle$, and $|N\rangle$. Its ground-state coherence is established by the coupling and probe fields before the switching field is switched on. The transient absorption of the probe field in the EIT subsystem is similar to a damped oscillator. For large Ω_{sc}, the oscillator is underdamped, and the absorption cross section is

$$-\frac{3\lambda^2}{2\pi}C_{31}^2\left(\frac{\Omega_s}{\Omega_{sc}}\right)^2\frac{\Gamma}{2\omega}\mathrm{e}^{-(\Gamma/4)t}\sin(\omega t),$$

where $\omega=\sqrt{\Omega_{sc}^2/4-\Gamma^2/16}$. The transient probe absorption in the EIT subsystem oscillates at a frequency of ω and its amplitude decays exponentially at a rate of $\Gamma/4$. In the high-intensity limit, the amplitude of the absorption cross section of the EIT subsystem is smaller than that of the two-state system by a factor of Γ/Ω_{sc}. For small Ω_{sc}, the oscillator is overdamped, and the absorption cross section is

$$-\frac{3\lambda^2}{2\pi}C_{31}^2\left(\frac{\Omega_s}{\Omega_{sc}}\right)^2\frac{\Gamma}{\Gamma'}(\mathrm{e}^{-[(\Gamma-\Gamma')/4]t}-\mathrm{e}^{-[(\Gamma+\Gamma')/4]t}),$$

where $\Gamma'=\sqrt{\Gamma^2-4\Omega_{sc}^2}$. The transient probe absorption in the EIT subsystem behaves as the superposition of two exponential functions with decay rates of $(\Gamma\pm\Gamma')/4$. In the low-intensity limit, the absorption cross sections of the two subsystems are comparable. The smallest decay rate of the three exponential functions is $(\Gamma-\Gamma')/4\approx\Omega_{sc}^2/2\Gamma$. A more detailed description of the above analysis can be found in Ref. [25]. The transient behaviors of the photon switching, including the rise times in the high- and low-intensity limits and the oscillation in the absorption curve, are clearly explained by the sum of the two absorption cross sections.

In Ref. [7], one-photon switching is achieved under the condition that the pulse length of the switching field is set equal to the delay time of the probe propagation in a medium. The delay time is given by $T_D=(\Gamma\Omega_c^2/\Omega_{sc}^4)X_{\mathrm{OD}}$, where X_{OD} is the optical density of the medium when the probe field propagates alone. To satisfy the condition that the transmission of the probe pulse is reduced to $1/e$ in the presence of the switching pulse, it is required that $\Omega_{sc}^2=\Omega_c^2X_{\mathrm{OD}}/(X_{\mathrm{OD}}-1)$. Consequently, $T_D=(X_{\mathrm{OD}}-1)\Gamma/\Omega_{sc}^2$. The dashed lines in Fig. 5 show T_D as a function of Ω_{sc} at several optical densities. The rise time of the probe absorption should be much less than the propagation delay time of the probe pulse or the length of the switching pulse. From Fig. 5, we can determine the value of X_{OD} or, equivalently, the ratio of Ω_s to Ω_c, which will ensure that the switch for the probe pulse is closed. For the example of $X_{\mathrm{OD}}=10$ and $\Omega_c=\Gamma$ used in Ref. [7], we find that the delay time is about 2 times larger than the rise time.

In conclusion, we have experimentally studied the transient behaviors of photon switching by quantum interference. Our experimental data were shown to be consistent with theoretical predictions, and we carried out a systematic analysis of the transient behaviors of the probe absorption. A simple picture based on the dark and nondark states was provided to explain these behaviors. We show that the rise time of the probe absorption cannot be less than $2/\Gamma$ and is equal to $2\Gamma/(\Omega_c^2+\Omega_s^2)$ in the low-intensity limit [26]. The four-state photon switching system based on the EIT scheme has potential applications in quantum optics and the manipulation of quantum information. The present findings show that the single-photon switching in the four-state system predicted by Harris and Yamamoto [7] is feasible with respect to the rise time of the probe absorption. This single-photon switching leads to the intriguing possibility that, in a medium with large optical density and negligible ground-state relaxation rate, a threefold entangled state may be realizable. Such a threefold entangled state could potentially be used in the manipulation of quantum information.

This work is supported by the National Science Council under NSC Grant No. 92-2112-M-007-021.

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