Single-Photon Switch Based on Rydberg Blockade

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All-optical switching is a technique in which a gate light pulse changes the transmission of a target light pulse without the detour via electronic signal processing. We take this to the quantum regime, where the incoming gate light pulse contains only one photon on average. The gate pulse is stored as a Rydberg excitation in an ultracold atomic gas using electromagnetically induced transparency. Rydberg blockade suppresses the transmission of the subsequent target pulse. Finally, the stored gate photon can be retrieved. A retrieved photon heralds successful storage. The corresponding postselected subensemble shows an extinction of 0.05. The single-photon switch offers many interesting perspectives ranging from quantum communication to quantum information processing.

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The switch is the device that lies at the heart of digital signal processing which has revolutionized the fields of communication and computation. In both fields, optical techniques are increasingly gaining importance. For example, present-day high-bandwidth internet connections operate optically. In the field of computing, perspectives for optical techniques are being studied, too [1,2]. They rely on all-optical switching and promise high bandwidth and low dissipated power. This creates a generic interest in the fundamental low-power limit of an all-optical switch, which is reached when the incoming gate pulse contains only one photon. Such a single-photon switch operates on the level of a single quantum and is hence well suited for applications in quantum technology. For example, it offers perspectives for heralded quantum memories which will be essential for realizing quantum repeaters [3], for efficiently detecting optical photons in a nondestructive measurement [4], for generating Schrödinger-cat states [5], and for various other applications in the fields of quantum communication and quantum information processing [6–8].

The field of all-optical switching with a huge number of photons per gate pulse had traditionally been dominated by nonlinear optics with techniques such as saturable absorbers and optical bistability. Building a single-photon switch with those techniques would be very difficult because the nonlinearities in nonlinear crystals are tiny at the single-photon level. Lately, however, electromagnetically induced transparency (EIT) [9] enriched the field of all-optical switching. If EIT is combined with Rydberg states [10], one can use Rydberg blockade [11,12] to create very large nonlinearities [13–18]. This triggered a proposal for building single-photon quantum devices [19]. Experiments observed all-optical switching in different systems; see e.g., Refs. [20–27]. However, all these experiments required ~20 or more incoming photons per gate pulse to obtain a clearly visible switching effect. A very recent experiment demonstrated all-optical switching with 2.5 to 5 incoming photons based on normal-mode splitting in a cavity [28].

Here we experimentally demonstrate all-optical switching with a gate pulse that contains only one incoming photon on average, or even fewer. This gate pulse reduces the transmission of a subsequent target pulse by a factor of $\epsilon = 0.812 \pm 0.001$. To achieve this goal, we send the gate pulse into an ultracold atomic gas and store it as a Rydberg excitation using a slow-light technique based on Rydberg EIT. Next, the target pulse is sent through the atomic medium. Without the gate pulse, Rydberg EIT would result in high transmission of the target pulse. With the gate pulse, however, Rydberg blockade suppresses the transmission of the target pulse. After application of the target pulse, we can retrieve the stored gate excitation. This shows that coherence in the stored excitation survives the target pulse. Using the retrieval as a herald to indicate successful storage events, we obtain an extinction of $\epsilon = 0.051 \pm 0.004$ in the postselected subensemble. We study the dependence of ϵ on the numbers of incoming gate and target photons. The Rydberg blockade displays a lifetime of ~60 μs if the target pulse is delayed relative to the gate pulse. The dephasing rate that limits the number of retrieved excitations depends linearly on the atomic density.

Schemes of the experimental setup and the atomic levels are shown in Figs. 1(a) and 1(b). Signal and control light have wavelengths of $\lambda_s = 795$ nm and $\lambda_c = 474$ nm and waists ($1/e^2$ radii of intensity) of $w_s = 8\ \mu$m and $w_c = 12\ \mu$m. The power of the control light is $P_c = 32$ mW for target and retrieval and half as large for the gate. The ultracold gas consists of $N = 2.2 \times 10^5$ atoms at a temperature of $T = 0.43\ \mu$K, which is a factor of ~3 above the critical temperature for Bose-Einstein condensation. The atoms are held in a crossed-beam optical dipole trap at a wavelength of 1064 nm with measured trap frequencies of $(\omega_x, \omega_y, \omega_z)/2\pi = (136, 37, 37)$ Hz. All atoms are prepared in state $|g\rangle$. A magnetic field of ~0.2

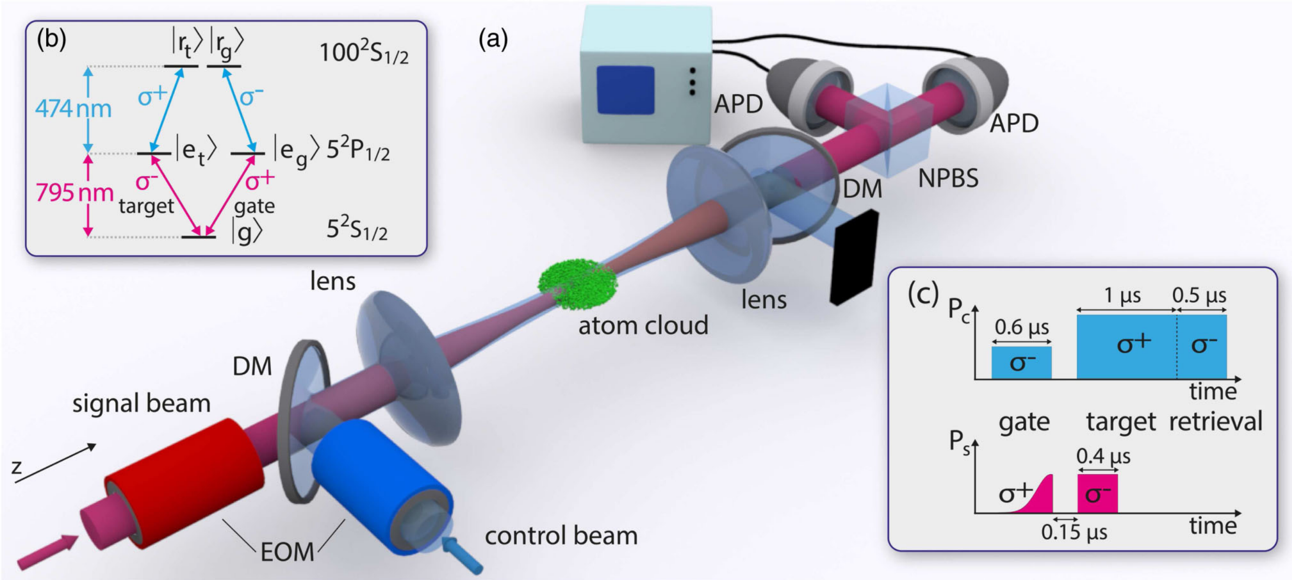

FIG. 1 (color online). (a) Simplified scheme of the optical beam path. Signal and control beams for Rydberg EIT copropagate along the z axis. After propagation through an ultracold gas of ^{87}Rb atoms, a dichroic mirror (DM) splits off the control light, sending it onto a beam dump. The signal light is divided by a nonpolarizing 50:50 beam splitter (NPBS) and detected on two avalanche photodiodes (APDs). Electro-optic modulators (EOMs) are used to set the incoming polarizations to either σ^+ or σ^-. (b) Atomic level scheme. Signal light with polarizations σ^+ and σ^- couples the ground state $|g\rangle = |5^2S_{1/2}, F = 1, m_F = -1\rangle$ with the excited states $|e_g\rangle = |5^2P_{1/2}, F = 2, m_F = 0\rangle$ and $|e_t\rangle = |5^2P_{1/2}, F = 2, m_F = -2\rangle$ for the gate and target pulse, respectively. Control light with polarizations σ^- and σ^+ couples states $|e_g\rangle$ and $|e_t\rangle$ with Rydberg states $|r_g\rangle = |100^2S_{1/2}, m_J = m_I = -1/2\rangle$ and $|r_t\rangle = |100^2S_{1/2}, m_J = 1/2, m_I = -3/2\rangle$ for the gate and target pulse, respectively. (c) Timing of incoming light, see text.

Gauss along the z axis preserves the spin orientation. The efficiency for collecting and detecting a transmitted signal photon is 27%. See Ref. [29] for further details.

Figure 1(c) shows the timing sequence of the incoming light pulses. The gate pulse is followed by a dark time $t_d = 0.15$ μs and then by the target pulse. Both pulses consist of light at the signal and control wavelengths. The signal light is resonant with an atomic transition, causing absorption. The control light creates EIT, thus suppressing the absorption of the signal light. The gate control light is switched off while a large part of the gate signal light is inside the medium due to a small group velocity. This stores gate signal photons in the medium in the form of Rydberg excitations.

To prevent the target control light from reading out these stored Rydberg excitations, the polarization of the control light is switched from σ^- for the gate pulse to σ^+ for the target pulse. Hence, the target control light cannot couple the stored Rydberg excitations to any state in the $5^2P_{1/2}$ manyfold, because such a state would require $m_J = -3/2$, contradicting $J = 1/2$. The signal light polarization is also switched. See Ref. [29] for further details.

The long-range character of the van-der-Waals potential $V(r) = -C_6/r^6$ between Rydberg atoms causes Rydberg blockade. Here r is the interatomic distance and C_6 is the van-der-Waals coefficient. Because of $V(r)$, the presence of a Rydberg excitation shifts the resonance frequency of the EIT feature for other incoming photons. This yields a blockade radius [29] of $r_b = 14$ μm. For $r < r_b$, the resonance shift is larger than the width of the EIT feature and the system is shifted out of the EIT resonance, resulting in absorption. Our experiment is carried out in the regime $w_s \lesssim r_b$, where the blockade sphere surrounding a single Rydberg atom extends over the full transverse profile of the signal beam. Ideally, one would expect that a single Rydberg excitation stored during the gate pulse should reduce the transmission of the target signal beam to near zero. This brings us into a new regime in which we study the absorption that a propagating excitation experiences due to a stationary excitation stored during a previous pulse.

After the target signal pulse has left the medium, we switch the polarization of the control light back from σ^+ to σ^-. This retrieves the excitations stored during the gate pulse. We can use postselection conditioned on the detection of a retrieved photon as a powerful tool for exploring the full potential of Rydberg blockade as a mechanism for all-optical switching, eliminating the reduction of performance due to imperfect storage.

This gate-target pulse sequence is repeated with a cycle repetition time [29] of $t_{\rm cyc} = 100\mu s$. Over the course of several thousand gate-target cycles, the atom number drops, so that a new atomic sample must be loaded.

Experimental results are shown in Fig. 2. The black data show two large peaks. The first peak shows an undesired nonzero transmission during the gate pulse due to imperfect storage. The second large peak shows the number of

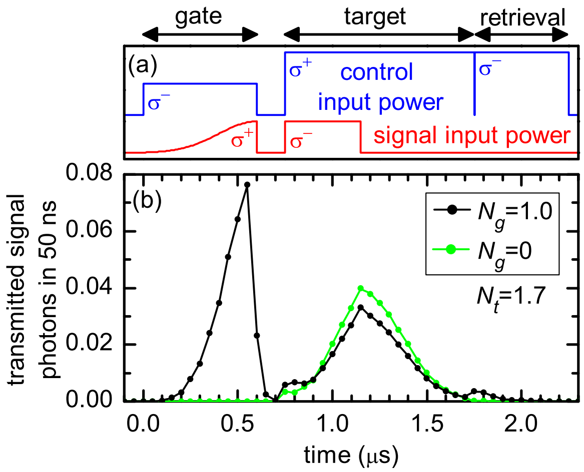

FIG. 2 (color online). (a) Input-power timing sequence. (b) Single-photon switch. Black data show the average number of transmitted signal photons for an average number of incoming signal photons during the gate pulse of $N_g = 1.0$. Green data show a reference with $N_g = 0$. The extinction between black and green target-pulse data is $\epsilon = 0.812 \pm 0.001$. The deviation from $\epsilon = 1$ is clearly observed, thus demonstrating a single-photon switch. The average number of incoming target signal photons is $N_t = 1.7$. The subensemble postselected on the detection of a retrieved photon yields $\epsilon = 0.051 \pm 0.004$.

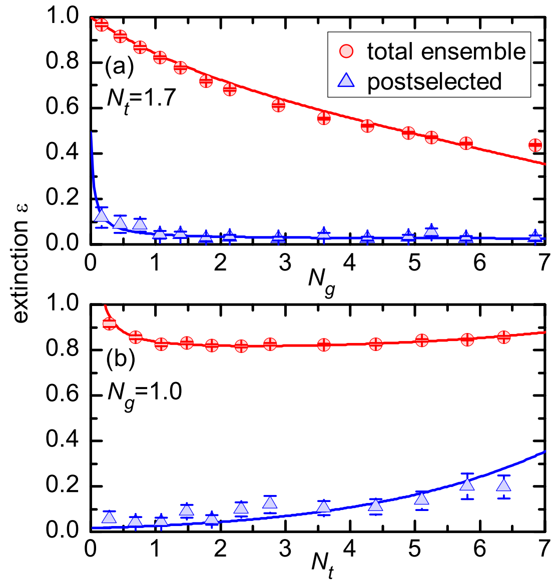

FIG. 3 (color online). (a) Dependence of the extinction ϵ on the incoming average photon number in the gate pulse N_g. Large N_g reduces the probability of storing zero Rydberg excitations, resulting in improved average extinction in the total ensemble. The subensemble that is postselected conditioned on detecting a retrieved gate excitation shows a drastically improved extinction. This proves that the nonideal extinction in the total ensemble is dominantly limited by the storage efficiency. (b) Dependence of the extinction ϵ on the incoming average photon number in the target pulse N_t. ϵ is fairly robust against changing N_t. All lines show fits to models from Ref. [29].

transmitted target photons which is reduced compared to the green reference data. There are also two smaller peaks in the black data: one at the beginning of the target interval, showing undesired partial readout of the stored gate excitation, the other at the beginning of the retrieval interval, showing the desired retrieval signal used for postselection.

To quantify how well the gate pulse reduces the transmission of target signal photons, we use the extinction

$$\epsilon = \frac{N_{\text{trans}}\text{with gate signal pulse}}{N_{\text{trans}}\text{without gate signal pulse}}, \tag{1}$$

where N_{trans} denotes the mean number of transmitted target signal photons in one gate-target cycle. A reduction of ϵ below 1 is clearly observed in Fig. 2, thus realizing an all-optical switch. As the average number of signal photons in the incoming gate pulse N_g is only 1.0, this measurement demonstrates a single-photon switch. The data in Fig. 2 were averaged over $\sim 8 \times 10^6$ gate-target cycles [29].

As the signal light is derived from an attenuated laser beam, the incoming gate photons have a Poissonian number distribution so that there is a noticeable probability that more than one photon enters the medium. However, the probability for storing more than one photon is negligible due to Rydberg blockade among the gate photons before storage, as experimentally confirmed by measuring the pair-correlation function in a retrieval experiment [29].

If the storage of different gate photons is uncorrelated, then the number of excitations stored will be Poissonian, too. This is expected for small N_g where blockade among gate photons has little relevance. Hence, the probability of storing zero excitations is $p_{s,0} = \exp(-\beta N_g)$ for small N_g. Here β is the storage efficiency in the absence of the Rydberg blockade. Obviously, $p_{s,0}$ sets a lower bound on the extinction $p_{s,0} \leq \epsilon$.

Figure 3(a) shows an experimental study of $\epsilon(N_g)$. Note that even for $N_g = 0.17$, we observe a deviation of ϵ from 1 by 4.5 standard errors in the total ensemble and by 20 standard errors after postselection. As the gate photons create a blockade for each other, the simple estimate above is only applicable for small N_g. Hence, β can be obtained from the absolute value of the slope $\beta = |d\epsilon/dN_g|$ at $N_g \to 0$. The lines are fits of models of Ref. [29]. The best-fit value is $\beta = 0.19$. The postselected subensemble shows a drastically improved extinction. For very small N_g, the postselected extinction deteriorates slightly. This is because for small N_g, the heralding probability decreases [29] so that background counts during the retrieval interval contribute an increasing fraction to the heralded events.

Figure 3(b) shows the dependence of ϵ on N_t at fixed target pulse duration. The dependence is rather weak, showing that the single-photon switch is fairly robust.

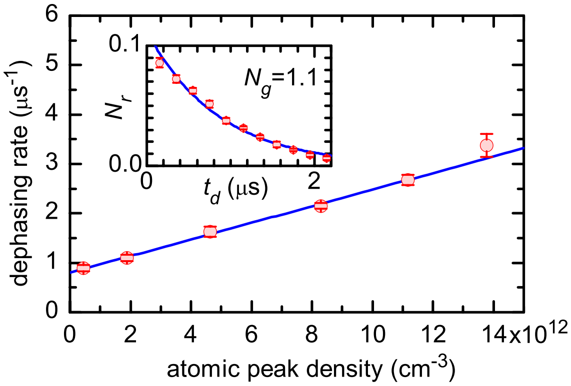

FIG. 4 (color online). The measured dephasing rate increases linearly with atomic density. Inset: measurement of the dephasing rate at a peak density of $\sim 2 \times 10^{12}$ cm^{-3}, where the experiment is normally operated. The retrieved photon number N_r decays as a function of the dark time t_d between gate and retrieval pulse in the absence of a target pulse. An exponential fit yields the $1/e$ dephasing rate.

The lines show fits to models from Ref. [29]. The slight deterioration of ϵ for larger N_t is due to the fact that scattering target signal photons reduces the atomic density, thus reducing the absorption. But this occurs only when averaging over a large number of cycles for each atomic gas. The deterioration of ϵ for small N_t is due to background photo detection events due to an undesired readout of stored gate photons during the target pulse.

All of the above measurements were performed with a dark time between the gate and target pulse of $t_d = 0.15$ μs. If t_d is increased, then the extinction ϵ decays with a $1/e$ time of 60 μs [29], again showing that this all-optical switch is fairly robust.

For comparison, note that the retrieved signal in the absence of a target pulse decays with a $1/e$ time of ~ 0.9 μs. This much shorter time scale is because retrieval is based on a phase-coherent collective directed emission, whereas the blockade merely needs Rydberg population. The ability to perform postselection crucially relies on a sufficiently long coherence time. We find that the coherence time depends linearly on the density of surrounding ground-state atoms, as shown in Fig. 4. We attribute this to a shift of the EIT resonance due to collisions between a Rydberg atom and surrounding ground-state atoms. Such a shift of ~ 10 MHz at a density of $\sim 10^{14}$cm^{-3} was observed in Ref. [30]. The inhomogeneity of the atomic sample converts this shift into a dephasing process. At zero density, the line extrapolates to a dephasing rate of 0.8 μs^{-1}. Because of the thermal motion of the atoms we expect 0.14 μs^{-1}, indicating that further decoherence mechanisms play a role. If future work can identify these mechanisms and remove them, we expect improvements in storage efficiency, heralding probability, and EIT transmission. This offers room for substantial improvements of the overall performance of the all-optical switch.

The work presented here opens the door to the new world of single-photon switching. With better performance, this will bring exciting perspectives on quantum information processing into reach. First, heralding successful storage is interesting for quantum memories. Storage times could be improved by subsequent transfer of the population into long-lived ground states. Second, the presence or absence of one gate photon could be mapped to the absence or presence of many transmitted target photons, respectively. Discriminating between the latter cases is easy, even at low detector efficiency. This could allow for detection of an optical photon with high sensitivity. Third, if the stored photon is eventually retrieved, then the detection of many target photons will represent a nondestructive detection of a single optical photon [31]. Fourth, if the incoming gate pulse contains a coherent superposition of zero and one photon, then the single-photon switch can create a Schrödinger-cat type coherent superposition of states with macroscopically different target photon numbers. Fifth, a photonic quantum-logic gate could be built based on this single-photon switch. For applications four and five, dissipation and decoherence must be kept low, which at first glance seems to contradict switching between transmission and absorption. However, if our switch is placed inside an optical resonator, resonant with the signal light, then transmission inside the atomic gas will lead to transmission through the resonator, whereas absorption inside the atomic gas will lead to reflection from the first mirror. This will convert the transmission-absorption switch into a transmission-reflection switch which could operate at low dissipation and decoherence.

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