Quantum gate for Q switching in monolithic photonic-band-gap cavities containing two-level atoms

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Photonic-band-gap cavities are prime solid-state systems to investigate light-matter interactions in the strong coupling regime. However, as the cavity is defined by the geometry of the periodic dielectric pattern, cavity control in a monolithic structure can be problematic. Thus, either the state coherence is limited by the read-out channel, or in a high-Q cavity, it is nearly decoupled from the external world, making measurement of the state extremely challenging. We present here a method for ameliorating these difficulties by using a coupled cavity arrangement, where one cavity acts as a switch for the other cavity, tuned by control of the atomic transition.

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As the maturity and sophistication of quantum optics progresses, there is a growing movement to translate such effects into practical devices. This impetus suggests, for reasons of scalability and practicality, the need for viable solid-state technologies to produce and distribute single photons as an enabling technology for derivative quantum devices. In particular we are concerned with the role played by cavity quantum electrodynamics (CQED) in such devices.

CQED has been used to great effect in the generation of deterministic, transform limited single (and higher-order Fock states) photon pulses [1], and schemes exist which incorporate CQED for quantum computing [2], and entanglement generation [3]. More recently "hybrid" schemes for quantum computation have been suggested incorporating matter qubits in cavities with single photon generation, linear optics, and high fidelity photon detection [4]. However, many of these schemes (with notable exceptions) will be problematic to scale or to remove from laboratory environments.

Given difficulties with implementing most present schemes in nonresearch environments, significant attention has turned towards photonic-band-gap (PBG) cavities as quantum cavities. This is due to their superb photonic confinement properties and the recent realization of high-Q cavities with small mode volume (of order the wavelength^3) [5,6]. These successes have been fueled by a combination of technological imperatives and advances in fabrication.

A PBG material is created by producing a periodic modulation in the dielectric function of a material so that Bragg interference prevents propagation of certain modes across the structure. Such structures may be two dimensional, with confinement in the third dimension realized by classical waveguiding, or by creating a three-dimensional lattice. We concentrate on the former example as it is easier to produce, and has so far yielded the most dramatic effects. The most popular configuration for 2D PBG structures is a thin membrane with a 2D array of holes (a lattice) drilled in it. A defect (usually an undrilled hole, local variation in lattice spacing, or a combination of both) defines a PBG cavity, as any photon injected into that site cannot propagate laterally away from the defect. In this way, PBG cavities can constitute extremely good cavities with low loss, high coupling and low mode volume, all necessary conditions for probing the strong-coupling limit of CQED.

One problem with high-Q cavities is the difficulty of outcoupling excitations from the cavity [7]. One would like a Q switch, a device that can be modulated in some fashion to change the cavity from high Q to low Q, with the optical intensity dumped from the cavity in a controlled fashion: we will term such a device a "gate." Q switching is well known for classical laser applications [8], but is less easy for PBG cavities, although some recent proposals exist including mechanical switches [9], and nonlinear optical effects [10]. However, in monolithic structures where we cannot use mechanical or thermal effects, and operating at low light levels, there have been no suitable suggestions for an effective Q switch in PBG cavities. This is the problem addressed in this paper.

The structure we consider is a coupled cavity arrangement, similar in spirit to that studied by Waks and Vukovic [11], where two defects in the PBG lattice were placed in close proximity to form evanescently coupled cavities. Our arrangement is shown schematically in Fig. 1(a), where the left-hand cavity is the storage cavity (or simply cavity), the right-hand cavity is the gate, which is in close proximity to a waveguide, or other leaky, classical region. In this limit we can describe the coupling between the distinct regions (cavity, gate, and waveguide), which is due to evanescent leakage of the electromagnetic modes, as being equivalent to photon hopping between the regions [12]. In addition to the previously considered systems, however, we augment this arrangement by placing a single two-state atomic system in the center of the cavity and gate [13], where the transition frequency of the atom can be controlled by some external control field. An example of a system that could realize such an architecture would be a single crystal diamond with photonic crystal drilled using focused ion beam milling and liftoff [14] where single ion implantation techniques [15] are used to locate individual nitrogen-vacancy centers in the center of

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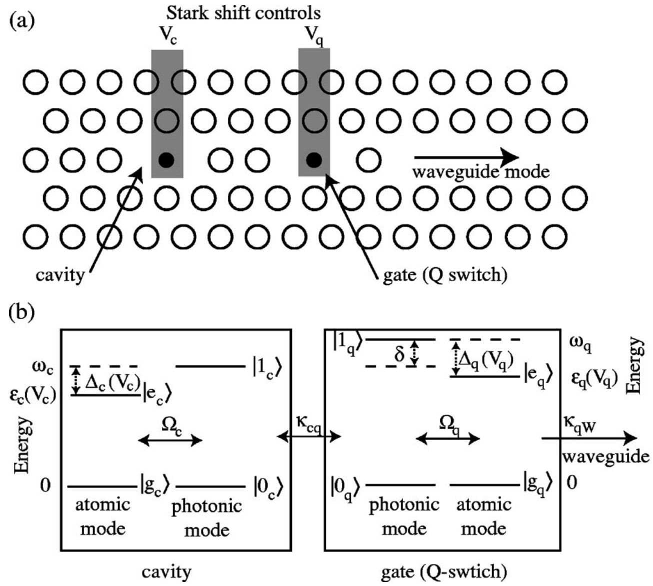

FIG. 1. (a) Schematic of coupled cavity system. The open holes correspond to the PBG lattice, where the missing holes constitute the cavities and waveguide. The filled circles are the atoms, with resonance controlled via top gates (shaded rectangles). The left defect is the "cavity," the right the gate, and outcoupling is via the waveguide mode on the right. (b) Diagram of states, energy levels, energy separations and couplings in the bare basis. Solid lines represent energy levels, dashed lines are used to show the references used for energy separations and arrows to depict coupling, reversible couple by double arrows, irreversible coupling (outcoupling) via a single arrow. All frequencies as shown in the diagram are positive. After initial fabrication, we only have control over the atomic transition energies ε_c and ε_q.

the cavity, controlled via the linear Stark shift [16]. It is this control of the atomic frequency that constitutes our sole *dynamic* (i.e., post fabrication) control of the system parameters, and is responsible for the Q-switching possibilities that we discuss in this paper.

The method for Q switching this system can be understood easily, and is a logical extension of previous work on cavity QED and photon blockade [17]. First the cavity is arranged so that one and only one photon is loaded into the cavity via some external pump, and the gate is in its ground state. The cavity and gate resonances are initially dissimilar, so that light from the cavity cannot leak across to the gate. Secondly the eigenmodes of the gate are varied by changing the resonance frequency of the atom in it, and when one of the gate modes is resonant with a mode of the cavity, photon hopping occurs. The gate is a relatively bad cavity, coupled to the output modes of a waveguide, and so photons leaking into the gate are rapidly outcoupled to the waveguide. As photon hopping is the source of the cavity-gate coupling, it is clear that optimal outcoupling results from balancing the competing needs of large cavity-gate detuning, with photonic population of the resonant mode of the gate at the outcoupling resonance. These points will be made more explicit by considering the model Hamiltonian.

The Hamiltonian for our system is written

$$\mathcal{H} = \mathcal{H}_c + \mathcal{H}_q + \mathcal{I}_c + \mathcal{I}_q + \mathcal{P}, \tag{1}$$

$$\mathcal{H}_\alpha/\hbar = \varepsilon_\alpha(V_\alpha)\sigma_\alpha^- \sigma_\alpha^+ + \omega_\alpha a_\alpha^\dagger a_\alpha, \tag{2}$$

$$\mathcal{I}_\alpha/\hbar = \Omega_\alpha(\sigma_\alpha^- a_\alpha + \sigma_\alpha^+ a_\alpha^\dagger), \tag{3}$$

$$\mathcal{P}/\hbar = \kappa_{cq}(a_q^\dagger a_c + a_q a_c^\dagger), \tag{4}$$

where $\mathcal{H}_\alpha$ and $\mathcal{I}_\alpha$ refer to the bare and interacting parts of the Hamiltonians, respectively, for $\alpha = c, q$ for cavity or gate (Q switch). $\varepsilon_\alpha(V_\alpha)$ is the transition frequency of the atom in α which can be controlled by the Stark shifting gates at some potential V, the exact functional dependence of the Stark shift on gate potential is not important. ω_α is the resonance frequency of the photon in α and Ω_α is the atom-cavity coupling (one-photon Rabi frequency) in α. The σ_α are the usual Pauli operators for the atoms in α and a_α is the usual photon annihilation operator in α. $\mathcal{P}$ describes the photon hopping, with coupling κ_{cq}. Coupling to the external waveguide is described via a non-Hamiltonian term which will be introduced in the density matrix formalism. All these terms are depicted schematically in Fig. 1(b).

In general, the two cavity system with two atoms is a moderately complicated problem to treat exactly, however, by considering just the one quantum manifold (i.e., where only one quantum of excitation is in the system), and assuming that the detuning between the cavities is large, i.e., $\omega_q - \omega_c = \delta \gg \kappa_{cq}, \Omega_\alpha$, we get significant insight. In this limit we can solve each cavity independently (i.e., ignoring κ_{cq} as our zeroth order approximation) to get the approximate eigenstates, which are the well-known dressed states

$$|\pm_c g_q 0_q\rangle = \frac{\left(-\dfrac{\Delta_c(V_c)}{2} \pm \chi_c\right)|g_c 1_c\rangle + \Omega_c |e_c 0_c\rangle}{\sqrt{2\chi_c^2 \pm \chi_c \Delta_c(V_c)}} |g_q 0_q\rangle,$$

$$|g_c 0_c \pm_q\rangle = \frac{\left(-\dfrac{\Delta_q(V_q)}{2} \pm \chi_q\right)|g_q 1_q\rangle + \Omega_q |e_q 0_q\rangle}{\sqrt{2\chi_q^2 \pm \chi_q \Delta_q(V_q)}} |g_c 0_c\rangle, \tag{5}$$

where we have introduced $|g\rangle$ and $|e\rangle$ as the states of the atoms $\Delta_\alpha(V_\alpha) = \omega_\alpha - \varepsilon_\alpha(V_\alpha)$, the detuning, and $\chi_\alpha = \sqrt{[\Delta_\alpha(V_\alpha)/2]^2 + \Omega_\alpha^2}$, the generalized Rabi frequency. The associated eigenenergies are

$$E_{|\pm_c, g_q 0_q\rangle} = \pm\chi_c - \Delta_c(V_c)/2, \tag{6}$$

$$E_{|g_c, 0_c \pm_q\rangle} = \delta \pm \chi_q - \Delta_q(V_q)/2. \tag{7}$$

By setting $\Delta_c = 0$ and $\Delta_q = -\delta < 0$, we can calculate the approximate interaction strength (coupling matrix element) of the gate induced resonance between the cavities, which is (for example, between $|+_c g_q 0_q\rangle$ and $|g_c 0_c -_q\rangle$ at the gate defined resonance $\Delta_q = -\delta + \Omega_c$)

$$\mathcal{J} = \langle g_c 0_c -_q | \mathcal{P} | +_c g_q 0_q\rangle = \frac{1}{\sqrt{2}} \frac{\Omega_q}{\delta} \kappa_{cq}. \tag{8}$$

The value of $\mathcal{J}$ sets the time scale for the interaction photon hopping. In particular, if we wish to adiabatically transfer the

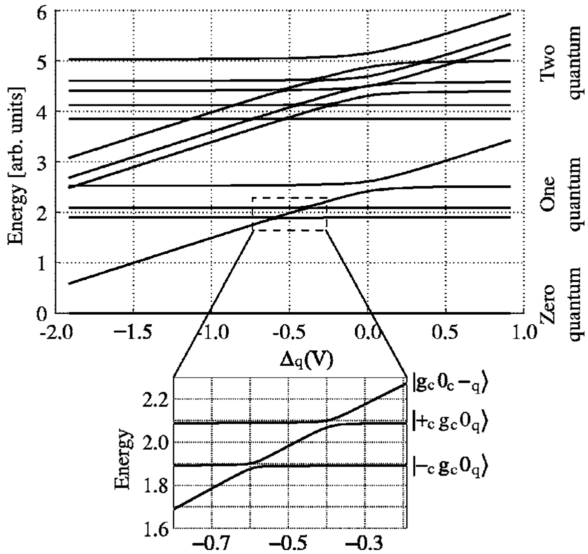

FIG. 2. Eigenvalues for the two-cavity system as a function of Δ_q for $\Omega_c=\Omega_q=\kappa_{cq}=0.1$, $\omega_c=2$, $\omega_q=2.5$, and $\delta=0.5$. The eigenspectrum naturally divides into the component manifolds, we are most interested in the one quantum manifold. The bottom trace shows a closeup of the one quantum manifold, highlighting interactions between the cavities, via the off-resonant dressed state. The resultant anticrossings indicate coupling between the gate and cavity, and hence where switching can occur.

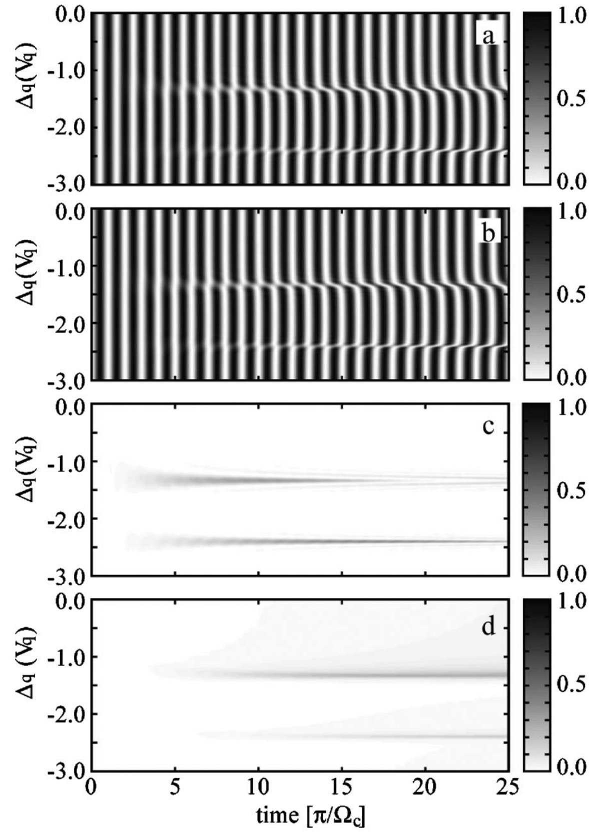

FIG. 3. Transient evolution of the cavity-gate-waveguide system showing populations in the states (a) $|g_c 1_c g_q 0_q 0_W\rangle$, (b) $|e_c 0_c g_q 0_q 0_W\rangle$, (c) $|g_c 0_c g_q 1_q 0_W\rangle$, and (d) $|g_c 0_c g_q 0_q 1_W\rangle$, respectively (the population in $|g_c 0_c e_q 0_q 0_W\rangle$ is never visible for these parameters), as a function of time and Δ_q given initial state $|g_c 1_c g_q 0_q 0_W\rangle$ for $\Omega_c=\Omega_q=0.5$, $\kappa_{cq}=\kappa_{qm}=0.1$ and $\delta=2$. Note the oscillation frequency in the range $-\delta-\Omega_c<\Delta_q<-\delta+\Omega_c$ is found to be $\sqrt{\Omega_c^2+\mathcal{J}^2}$. The maximum of the population transfer peaks in the waveguide in (d) are 0.26 at $\Delta_q\sim-\delta+\Omega_c$, and 0.17 at $\Delta_q\sim-\delta-\Omega_c$.

excitation from the cavity to the gate, the sweep rate of the gate should be slow compared to $1/\mathcal{J}$.

To further explore the coupling between the cavities, we present in Fig. 2 the eigenspectra determined by numerically solving the Hamiltonian in Eq. (1) without further approximation as a function of $\Delta_q(V_q)$, for $\Omega_c=\Omega_q=\kappa_{cq}=0.1$, $\omega_c=2$, $\omega_q=2.5$, and $\delta=\omega_q-\omega_c=0.5$. These parameters were chosen to simply highlight the important system features. The system conveniently breaks into three manifolds, distinguished by the total number of quanta, the zero quantum manifold is the lowest, along the $\Delta_q(V_q)$ axis, then the one quantum and two quantum. As we are most interested in the resonances between the cavity and gate in the one quantum manifold, we present a closeup of this in the inset to Fig. 2, where the anticrossings indicating photon hopping between the cavity and gate are clearly visible. Note that these parameters were merely chosen to demonstrate the relevant processes, and all units are arbitrary.

The previous analysis just treats coupling between the cavity and gate, but to proceed further we need to include the coupling to the waveguide mode (W). This is best done by introducing an irreversible loss term, analogous to spontaneous emission, which models coupling into an extra waveguide mode. We then solve the density matrix equations of motion to examine the transient coupling into the waveguide mode. Concretely we solve for the density matrix ρ using

$$\dot{\rho}=-\frac{i}{\hbar}[\mathcal{H},\rho]+\kappa_{qw}\mathcal{L}[\rho,a_W^\dagger a_q],$$

$$\mathcal{L}[\rho,a_W^\dagger a_q]=a_W^\dagger a_q\rho a_q^\dagger a_W-\frac{a_q^\dagger a_W a_W^\dagger a_q\rho+\rho a_q^\dagger a_W a_W^\dagger a_q}{2}. \qquad (9)$$

κ_{qW} is the gate-waveguide coupling. An example of the evolution obtained is presented in Fig. 3 which shows the populations in the bare state basis, from top to bottom: $|g_c 1_c g_q 0_q 0_W\rangle$, $|e_c 0_c g_q 0_q 0_W\rangle$, $|g_c 0_c e_q 0_q 0_W\rangle$, $|g_c 1_c g_q 0_q 0_W\rangle$, and $|g_c 0_c g_q 0_q 1_W\rangle$, respectively, as a function of time and $\Delta_q(V_q)$ given initial state $|g_c 1_c g_q 0_q 0_W\rangle$. Clearly noticeable are coherent oscillations corresponding to Rabi oscillations in the cavity, and the gradual buildup of population in the waveguide mode. The increase in the oscillation frequency in the region $-\delta-\Omega_c<\Delta_q<-\delta+\Omega_c$ is a consequence of an increased eigenvalue splitting, similar to (but more complicated than) that seen in multiply coupled three state systems, see, for example, Ref. [18]. In the simpler, doubly driven three-state case, the oscillation frequency is given by the sum of the squares of the Rabi frequencies of the driving fields. In this case, the result is similar, with the oscillation frequency given by the sum of the squares of the interaction matrix elements, i.e., $\sqrt{\Omega_c^2+\mathcal{J}^2}$.

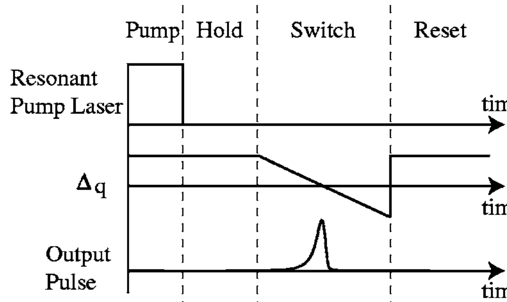

FIG. 4. Steps required to adiabatically generate Q switched single photon pulses as a function of time. Pump: first the cavity is pumped with a laser at frequency $\omega_c+\Omega_c$, and pulse area π, storing one quantum in the cavity. Hold: the system is maintained for a length of time, set by requirements of triggering the rest of the system and acceptable loss probabilities. Switch: Δ_q is adiabatically swept through the resonance, and the single photon pulse outcoupled. Reset: the system is returned to its initial state for repeated operation.

Although illustrating much of the necessary physics, it is clear that the results in Fig. 3 do *not* illustrate an effective mechanism for single photon generation. The reason here is that the outcoupling probability can be no better than 50%, and the Rabi oscillations render the system prone to nonadiabatic errors. Also pertinent is that a complicated set of interference fringes are observed which need to be understood for transient analysis. Given such limitations, it is preferable to initialize the system in the state $|+_c g_q 0_q 0_W\rangle$ and follow an adiabatic transition along the anticrossing between $|+_c g_q 0_q 0_W\rangle$ and $|g_c 0_c -_q 0_W\rangle$. By ensuring that the gate and waveguide mode are strongly coupled, i.e., κ_{qW} is large compared with the coupling matrix elements, population in the state $|g_c 0_c -_q 0_W\rangle$ will be rapidly transferred to $|g_c 0_c g_q 0_q 1_W\rangle$, and hence the cavity-gate resonance will act as an effective Q switch for the cavity. Initialization of the system could be achieved by pumping the cavity with light of frequency $\omega_c+\Omega_c$, which would be resonant with the $|g_c 0_c\rangle - |+_c\rangle$ transition, but not resonant with transitions to the two quantum manifold. A schematic of the steps required to adiabatically outcouple the single photon is shown in Fig. 4.

The results of the adiabatic transfer from cavity to gate are shown in Fig. 5, clearly showing both the population in $|g_c 0_c g_q 0_q 1_W\rangle$, which we denote ρ_{WW}, and the time derivative of this population $\dot{\rho}_{WW}$, which is proportional to the intensity of the resulting photon pulse. In this case we chose $\Omega_c=\Omega_q=\Omega$, $\delta=4\Omega$, $\kappa_{cq}=0.01\Omega$, $\kappa_{qW}=0.1\Omega$, $-3.2\Omega\leqslant\Delta_q\leqslant-2.2\Omega$ and the length of the sweep was $T_{\max}=2\times10^4\pi/\Omega$. Note that because of the difference between κ_{cq} and κ_{qW} the resultant single photon pulse is not a Gaussian. To retrieve a Gaussian pulse, one could either choose a system with equal photon hopping matrix elements or a more complicated gate sweep. Note that the integral of the derivative is unity, as required for a pulse of one photon.

When considering the operating parameters of the Q switch, it is also necessary to determine the quiescent fidelity, i.e., the photon leakage from cavity to the Q-switching gate when the switch is not activated. For simplicity, if we assume $\Delta_c=\Delta_q=0$ and $\delta\gg\kappa_{cq},\Omega_c,\Omega_q$, then the population in

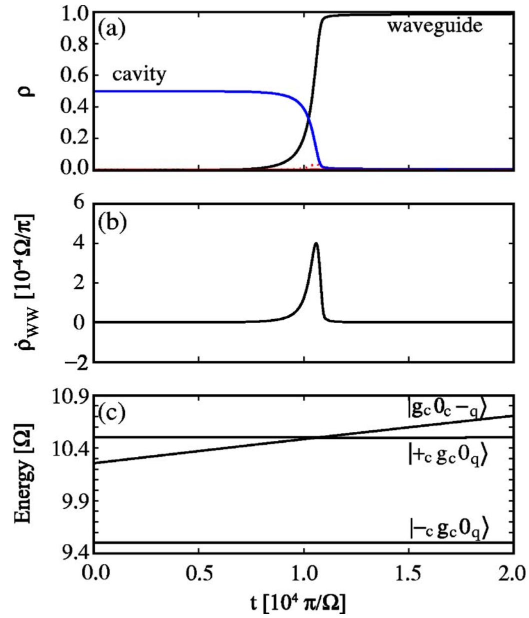

FIG. 5. (Color online) (a) State population (expressed in the bare basis) during the adiabatic sweep as a function of time given initialization in the state $|+_c g_q 0_q\rangle$ at $t=0$. Observe the smooth variation in population from cavity superposition states to the waveguide mode with minor, transient occupation of the gate. This is slightly asymmetric due to the asymmetric coupling to and from the gate. (b) Time derivative of waveguide occupation, $\dot{\rho}_{WW}$, proportional to the output intensity. Again, the pulse is asymmetric, although a tailored bias sweep could symmetrize the pulse. (c) Eigenvalues during the sweep, note that on this scale, the avoided crossing between $|+_c g_q 0_q\rangle$ and $|g_c 0_c -_q\rangle$ is not resolved. For this example, the parameters chosen were $\Omega_c=\Omega_q=\Omega$, $\delta=4\Omega$, $\kappa_{cq}=0.01\Omega$, $\kappa_{qW}=0.1\Omega$, $-3.2\Omega\leqslant\Delta_q\leqslant-2.2\Omega$ and the total sweep length was $T_{\max}=2\times10^4\pi/\Omega$.

state $|+_c g_q 0_q\rangle$, $\rho_{+_c}(t)$ at time t, given initialization in state $|+_c g_q 0_q\rangle$ at $t=0$, is

$$\rho_{+_c}(t)=\exp\left(-\frac{\kappa_{cq}^2}{2\delta^2}\kappa_{qW}t\right), \tag{10}$$

where κ_{cq}^2/δ^2 is the standard, steady state, off-resonant population leaking from the cavity to gate, which is then outcoupled at a rate κ_{qW}. Under the conditions used to generate Fig. 5, this equates to a population of $\rho_{+_c}=0.98$ at $t=2\times10^4\pi/\Omega$, or alternatively, at worst a 2% probability of the photon outcoupling from the cavity.

Finally we comment on the practicality of realizing our scheme in a realistic structure, and for our purposes we assume a PB cavity structure fabricated in diamond containing a single NV^- center at the maximum of the cavity mode. The wavelength of the zero-phonon line resonance of an NV^- center is $\lambda=638\times10^{-9}$ m, with frequency $\omega=2.95\times10^{15}$ Hz, and assuming that each cavity has volume $V=\lambda^3=(638\times10^{-9})^3$ m^3, then the atom cavity coupling will be $\Omega=\mu\sqrt{\omega/(2\hbar\epsilon_0 V)}\sim10^{10}$ Hz (given the electric dipole moment of the NV^- center of $\mu\sim10^{-29}$ C m^{-1}). For this degree of coupling, the tuning range of the centers should be many Ω. The tuning range reported in Ref. [16] is $\sim10^{12}$ Hz, which does not constitute an upper limit on the Stark tuning. Therefore the atomic tuning criterion should be easy to satisfy, and we presume $\delta = 10^{12}$ Hz.

If we assume that the cavities are in the good cavity limit, and that the cavity Q is dominated by photon loss due to the photon hopping between cavities and the waveguide, then the cavity Q must be fairly large to ensure minimal population leakage when we are not at the switching point. The figure of merit here is that the ratio κ_{cq}^2/δ^2 should be small. If we aim for a residual population of 10^{-4}, then $\kappa_{cq}/\delta \leqslant 10^{-2}$, i.e., $\kappa_{cq} = 10^{10}$ Hz and $\kappa_{qW} = 10\kappa_{cq} = 10^{11}$ Hz. If we assume that the cavity Q is dominated by the photon hopping terms, then we have (for the cavity)

$$Q_c = \frac{\omega}{\kappa_{cq}} \sim 10^5 \tag{11}$$

and the Q of the gate will be 10^4. Although technically demanding, we note that Q factors $\sim 10^7$ have been shown to be possible in silicon photonic-band-gap cavities on silica [6]. Furthermore, although we have studied our device in this demanding regime to clarify the effects, proof of principle experiments will be possible with significantly lower Q values by relaxing the requirements for adiabatic transfer and increasing the detuning of the cavities.

With these parameters, the pulse obtained by adiabatically switching the gate will be outcoupled in a time commensurate with $1/\mathcal{J} = 10^{-9}$ s. Without switching, the expected population in the waveguide mode over this time scale would be 0.01. The full set of required parameters are summarized in Table I.

In conclusion, we have presented a scheme for Q switching a photonic-band-gap cavity by controlling the resonance condition of an adjacent cavity. Each cavity contains a single two-level atom, and the transition frequency of the atom can be controlled via a Stark shifting electrode. We refer to the right-hand cavity as the gate which Q switches the cavity.

TABLE I. Nominal parameters for efficient Q-switching for NV^- centers embedded in an all-diamond photonic crystal.

Parameter	Value
Wavelength	638 nm
Transition frequency	2.95×10^{15} Hz
$\Omega_c = \Omega_q$	10^{10} Hz
Q_c	10^5
Q_q	10^4
κ_{cq}	10^{10} Hz
κ_{qw}	10^{11} Hz
δ	$\leqslant 10^{12}$ Hz

The resonance frequencies of the two cavities are initially dissimilar, but by tuning the atomic transition in the gate, a resonance condition between the cavity and gate is obtained, resulting in photon hopping between the cavities. By introducing a waveguide mode adjacent to the gate, photons leak rapidly out of the gate. Such a device constitutes a solid-state source of transform limited single photons on demand. An ideal system to test such concepts would be in micromachined diamond containing the nitrogen-vacancy color center, although our ideas can be applied to any photonic-band-gap cavity containing a two-level atom in the maximum of the cavity mode.

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