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Control of the electromagnetic environment of a quantum emitter by shaping the vacuum field in a coupled-cavity system

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We propose a scheme for the ultrafast control of the emitter-field coupling rate in cavity quantum electrodynamics. This is achieved by the control of the vacuum field seen by the emitter through a modulation of the optical modes in a coupled-cavity structure. The scheme allows the on-off switching of the coupling rate without perturbing the emitter and without introducing frequency chirps on the emitted photons. It can be used to control the shape of single-photon pulses for high-fidelity quantum state transfer, to control Rabi oscillations, and as a gain-modulation method in lasers. We discuss two possible experimental implementations based on photonic crystal cavities and on microwave circuits.

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I. INTRODUCTION

Spontaneous emission (SE) is at the heart of quantum optics and quantum photonics. Tremendous progress has been achieved in optimizing the SE of quantum emitters (QEs) in atomic systems [1] and artificial atoms such as quantum dots [2–4] and superconducting circuits [5] by placing them into resonators. These coupled-cavity–QE systems have been established as nonclassical light sources [6] and may serve as light-matter interfaces [1,7–9].

In cavities, the presence of an increased optical density of states enhances the emission and absorption properties of QEs. Typically, the QE-cavity interaction is governed by a coupling constant g given by the dipole moment of the QE and the vacuum electric field associated to the cavity mode. This interaction constant is thus given by intrinsic properties of the system, which are difficult to modify. However, g is a crucial parameter, since it determines the relevant time scale of the interaction. Indeed, a QE in the excited state decays into the cavity mode with a decay time $\tau_{dec} = (\frac{4g^2}{\kappa})^{-1}$ [10] in the weak-coupling regime ($g \ll \kappa$), where κ is the cavity loss rate. On the other hand, in the strong-coupling regime, where the coupling constant g exceeds the loss of the cavity κ, a coherent and reversible energy exchange between the cavity and the QE takes place with a characteristic time scale $\tau_r \propto 1/g$.

So far, the control of the QE-cavity interaction in the solid state has been performed mostly by tuning their spectral overlap [11–16] and, in the large majority of experiments, on time scales much longer than the interaction time. Dynamic control is, however, needed for the control of the photon waveform and the establishment of QE-photon entanglement. Such dynamic control has been demonstrated recently by using a combined variation of the loss rate and the cavity field by ultrafast carrier injection [17] and by ultrafast detuning in photonic crystal diodes [18]. However, both these techniques produce a temporal variation of the cavity frequency, resulting in a frequency chirp of the emitted photons, which limits the fidelity of quantum state transfer [19].

In this letter, we propose the concept of pure vacuum-field modulation as a method to control the QE-cavity interaction in real time. We demonstrate that the vacuum field of an optical mode in a given cavity can be completely suppressed by varying the frequency of two coupled lateral cavities, without producing any frequency chirp of the target mode. This enables the on-off switching of the QE-cavity interaction rate g, which is fundamentally different from control techniques based on controlling the QE-cavity detuning. As an example, we theoretically demonstrate the shaping of a single-photon pulse into a symmetric wave packet as a prerequisite for high-fidelity quantum state transfer [7]. Finally, we discuss two experimentally feasible platforms to implement the proposed scheme. The full and direct control of the light-matter interaction constant g represents a powerful tool to develop advanced applications in quantum information science, e.g., switching Rabi oscillations, and it can serve as the basis for a new class of gain-modulated lasers.

The paper is organized as follows. The underlying principle and the coupled mode theory is presented in Sec. II. The shaping of photon pulses is illustrated in Sec. III. In Sec. IV we discuss two different implementations of the proposed system, and a summary and conclusions are given in Sec. V.

II. SYSTEM AND MODEL

We first consider the coupling of three in-line cavities with a QE placed in the central cavity (called the target cavity) with frequency ω_t, as shown in Fig. 1. Furthermore, we assume that the frequency of the two outer cavities $\omega_{l,r} = \omega_t \pm \Delta$, named left and right control cavity in the following, can be tuned at will under the assumption that the detuning Δ is the same for both but with a different sign. The Hamiltonian of the empty three-cavity system reads ($\hbar = 1$)

$$H_{cc} = \begin{pmatrix} (\omega_t + \Delta) - i\kappa_l & \eta & 0 \\ \eta & \omega_t - i\kappa_t & \eta \\ 0 & \eta & (\omega_t - \Delta) - i\kappa_r \end{pmatrix}, \quad (1)$$

where $\kappa_{r,l,t}$ are the loss rate of the control cavities and target cavity, respectively, and η denotes the coupling rate between the adjacent cavities. Neglecting the loss rates, the Hamiltonian can be exactly diagonalized, which yields three nondegenerate eigenvalues ω_i ($i = 1,2,3$):

$$\begin{aligned} \omega_1 &= \omega_t, \\ \omega_{2,3} &= \omega_t \pm \sqrt{2\eta^2 + \Delta^2}. \end{aligned} \quad (2)$$

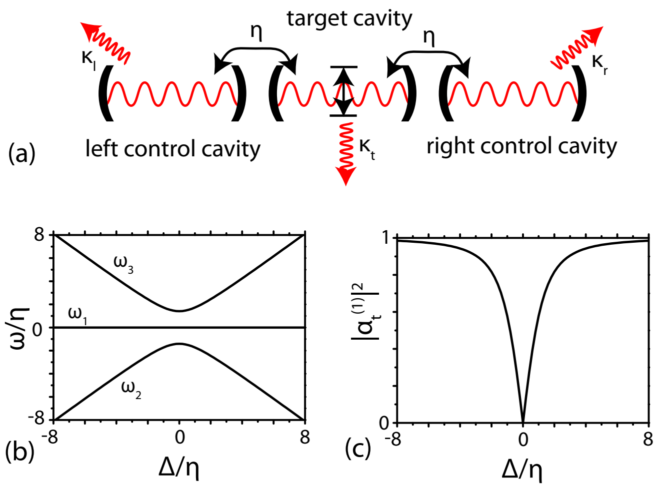

FIG. 1. (Color online) (a) Illustration of the coupled-cavity scheme. (b) Calculated eigenfrequencies ω_i/η as a function of the dimensionless cavity detuning Δ/η. (c) Calculated $|\alpha_t^{(1)}|^2$, which governs the modulation of the QE-cavity coupling constant g. The used parameters are $\kappa_{t,r,l} = 0.1\eta$.

The most interesting eigenvalue, which we will consider in the following, is $\omega_1 = \omega_t$ because its frequency is independent from the detuning of the control cavities. The corresponding eigenvector reads

$$\vec{\alpha}^{(1)} = \begin{pmatrix} \alpha_l^{(1)} \\ \alpha_t^{(1)} \\ \alpha_r^{(1)} \end{pmatrix} = \frac{1}{\sqrt{2 + (\Delta/\eta)^2}} \begin{pmatrix} -1 \\ \frac{\Delta}{\eta} \\ 1 \end{pmatrix}. \quad (3)$$

Figures 1(b) and 1(c) show the frequencies ω_i versus detuning and the target cavity fraction $|\alpha_t^{(1)}|^2$ of the mode ω_1 calculated by solving the eigenvalue problem for H_{cc}. Interestingly, by changing the detuning of the control cavities, the target cavity fraction $\alpha_t^{(1)}$ of the eigenmode can be changed from $\alpha_t^{(1)}(\Delta = 0) = 0$ to $\alpha_t^{(1)}(\Delta \gg \eta) \approx 1$. These properties provide already the main ingredient for the proposed system: the eigenmode at frequency $\omega_1 = \omega_t$ can be freely engineered to have no electric field component in the target cavity or the full field intensity corresponding to a decoupled target cavity. This has drastic consequences for a QE with frequency $\omega_e = \omega_t$ coupled to the target cavity with strength $g < \kappa_t$. Due to its spatial position it can couple to the mode in case $\alpha_t^{(1)} > 0$ or is completely decoupled if $\alpha_t^{(1)} = 0$ (i.e., when the three cavities are in resonance).

Moreover, also the effective loss rate of the modes can be changed due to the mixing of the cavity components. An approximate diagonalization ($\eta \gg \kappa_t,\kappa_c$) of Eq. (1) including the loss terms yields similar eigenvectors as given in Eq. (3). The effective loss of the coupled eigenmode ω_1 can then be written as $\kappa_1 = |\alpha_t^{(1)}|^2\kappa_t + |\alpha_l^{(1)}|^2\kappa_l + |\alpha_r^{(1)}|^2\kappa_r$. Thus, if the loss rates are equal, the tuning of the left and right control cavities allows for a pure g modulation.

In order to describe the interaction of the coupled-cavity system with the QE, we write the full Hamiltonian including the QE with frequency ω_e and the QE-target cavity coupling g using the basis of the coupled modes:

$$H_{\text{eff}} = \begin{pmatrix} \omega_e - i\gamma & g_1 & g_2 & g_3 \\ g_1 & \omega_1 - i\kappa_1 & 0 & 0 \\ g_2 & 0 & \omega_2 - i\kappa_2 & 0 \\ g_3 & 0 & 0 & \omega_3 - i\kappa_3 \end{pmatrix}, \quad (4)$$

where the coupling rate $g_i = \alpha_t^{(i)} g$ $(i = 1,2,3)$ and $\alpha_t^{(i)}$ describe the target cavity fraction of each eigenmode. Assuming that $\eta \gg g$ and $\omega_e = \omega_t$, the QE is spectrally decoupled from the modes $\omega_{2,3}$ (as we will assume in the following). It is immediately obvious that by changing the detuning of the control cavities Δ one can directly modulate the effective interaction strength between the QE and the cavity mode $\alpha_t^{(1)} g$, without changing the frequency of the coupled-cavity mode ($\omega_1 = \omega_t$). The SE rate of the QE (neglecting the spontaneous decay into leaky modes given by γ) can be described in the uncoupled-cavity case by $\gamma_0 = \frac{4g^2}{\kappa_t}$ [10]. This obviously changes in the case of the coupled cavities to $\gamma_{\text{eff}} = \frac{4|\alpha_t^{(1)}|^2 g^2}{\kappa_1}$. One obtains for the ratio of the SE rates

$$\frac{\gamma_{\text{eff}}}{\gamma_0} = |\alpha_t^{(1)}|^2 \frac{\kappa_t}{\left(|\alpha_t^{(1)}|^2\kappa_t + |\alpha_l^{(1)}|^2\kappa_l + |\alpha_r^{(1)}|^2\kappa_r\right)}, \quad (5)$$

which can be tuned from zero (all cavities in resonance and $|\alpha_t^{(1)}| = 0$) to one ($\Delta \gg \eta$ and $|\alpha_t^{(1)}| \approx 1$). We note that the stimulated emission rate and the modal gain in a microcavity laser are also directly related to the amplitude of the electric field at the QE's position. This vacuum-field modulation technique can therefore be used to modulate the gain without directly affecting the carrier population.

III. SHAPING OF PHOTON PULSES

The above-described system can be generally used to control all spontaneous and stimulated emission processes in a microcavity. In the following, as an example, we illustrate the shaping of an emitted photon pulse from an initially inverted QE. To this aim, one has to implement a dynamic tuning of the control cavities $\Delta \rightarrow \Delta(t)$ in such a way that the typically sharp rise of the emitted single-photon wave packet can be slowed down to exactly match the time-inverted decay tail.

We simulate the dynamics using a wave-function approach and replace the loss of the target cavity by the coupling to a quasicontinuum of modes representing a waveguide coupled to the target cavity [19,20]. The Hamiltonian of the system reads

$$H = \sum_{i=t,r,l} \omega_i a_i^+ a_i + \sum_{k=1}^{N} \omega_k b_k^+ b_k - \sum_{i=r,l} i\eta(a_t^+ a_i - a_i^+ a_t) + ig(a_t^+\sigma - \sigma^+ a_t) - \sqrt{\frac{\kappa_t \Delta\omega_k}{2\pi}} \sum_{k=1}^{N} (a_t^+ b_k - b_k^+ a_t). \quad (6)$$

The first term is the free evolution of the three-cavity modes with operators a_i and the second term describes the evolution of the quasicontinuum with mode frequencies ω_k and operators b_k. The remaining terms describe the cavity-cavity coupling, the QE-target cavity coupling, and the target cavity-continuum coupling, respectively. By plugging the expansion of the wave function (limiting ourself to only a single excitation in the system)

$$|\Psi\rangle = (c_e|e,0,0,0\rangle + c_t|g,1,0,0\rangle + c_l|g,0,1,0\rangle + c_r|g,0,0,1\rangle) \otimes |vac\rangle + |g,0,0,0\rangle \otimes \sum_k c_\kappa^{(k)} b_k^+ |vac\rangle \tag{7}$$

in the time-dependent Schrödinger equation $-i\partial_t|\Psi\rangle = H_{\text{eff}}|\Psi\rangle$ using an effective Hamiltonian, including the loss of the cavity modes and the decay of the QE, we obtain the evolution equations for the state amplitudes in a frame rotating at ω_t:

$$\dot{c}_e = gc_t - \gamma c_e, \tag{8}$$

$$\dot{c}_t = -gc_e - \eta(c_l + c_r) + \kappa' \sum_{k=1}^{N} c_\kappa^{(k)}, \tag{9}$$

$$\dot{c}_l = \eta c_t - i\Delta(t)c_l - \kappa_l c_l, \tag{10}$$

$$\dot{c}_r = \eta c_t + i\Delta(t)c_r - \kappa_r c_r, \tag{11}$$

$$\dot{c}_\kappa^{(k)} = -i\Delta_k c_\kappa^{(k)} - \kappa' c_t, \tag{12}$$

where $\kappa' = \sqrt{\frac{\kappa_t \Delta\omega_k}{2\pi}}$ is the target-cavity-continuum coupling and $\Delta\omega_k$ is the spacing of the quasicontinuum modes. The QE dynamics are given by the amplitude c_e, while the target and the two control cavities are described by amplitudes c_t and $c_{l,r}$, respectively. Finally the waveguide quasicontinuum is described by the amplitudes $c_\kappa^{(k)}$.

Starting from an inverted QE [$c_e(0) = 1$], we calculate the evolution of the state amplitudes as well as the output pulse, which is given by the inverse Fourier transform of the amplitudes of the continuum $c_\kappa^{(k)}(T)$ at a given time T much larger than the duration of the effective interaction of the QE with the cavity system [19,20]. We explicitly take into account the losses of the control cavities, while the main loss channel of the target cavity is the quasicontinuum coupling. Furthermore, the QE is weakly coupled to the cavity mode $g < \kappa_t$.

We consider three scenarios for the detuning. In case all three cavities are in resonance ($\Delta = 0$), the QE is completely decoupled from the cavity and does not decay into the cavity mode (not shown). For a constant detuning $\Delta \gg \eta$ [Fig. 2(a)], the QE interacts with the uncoupled target cavity, resulting in an exponential decay with the standard spontaneous emission rate $4g^2/\kappa$, as shown in Fig. 2(b), where we plot the QE population $|c_e(t)|^2$ and the target cavity photon population $|c_t(t)|^2$. Due to the active modulation of the QE-cavity coupling, in the case of a time-dependent detuning $\Delta(t)$ one can shape the emitted wave packet, in principle, arbitrarily. Here, we illustrate the shaping into a symmetric pulse, which can be absorbed with in principle unit fidelity by a similar system with time-inverted control cavity detuning. The time-dependent detuning and the corresponding dynamics of the QE and cavity populations are displayed in Figs. 2(c) and 2(d), respectively. The shaped output pulse emitted into the waveguide is shown in Fig. 3(a) together with the natural pulse shape without active control in the case of $\Delta \gg \eta$. A Gaussian fit of the shaped photon pulse reveals a nearly perfect time-symmetric wave packet. The phase of the symmetric photon pulse [Fig. 3(b)] is constant in time (apart from the small deviations due to numerical error caused by small amplitudes), since there is no frequency tuning involved, in contrast to recent proposals using ultrafast electrical control of the QE energies [19]. The fidelity F of the absorption of this symmetric photon pulse by a similar system with time-inverted operation can be obtained by calculating the overlap integral of the incident pulse with the time-inverted pulse [19]. This yields, for the present case, $F = 0.997$ and can be further improved by optimization of the dynamic tuning.

The dynamic tuning of QE-cavity coupling requires that the coupled-cavity modes follow the adiabatic eigenstates given by Eq. (2). Assuming a linear time dependence of the detuning $\Delta(t) = \beta t$, the problem can be described by a generalized Landau-Zener model [21,22]. The resulting condition $\sqrt{\beta} \ll \eta$ needs to be satisfied to ensure adiabaticity. This sets an upper speed limit on the dynamic tuning. Including the QE-cavity coupling, the condition for shaping photon pulses can be written as $4g^2/\kappa \ll \sqrt{\beta} \ll \eta$. In case

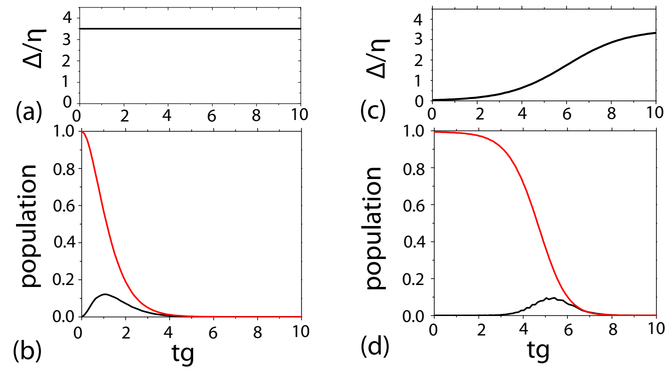

FIG. 2. (Color online) Temporal dependence of (a) the detuning Δ/η and (b) the QE population (red line) and the intracavity photon population (black line) for the case of constant detuning. (c) and (d) show the same quantities for the case of time-varying detuning $\Delta(t)$. The parameters used are $\{g,\eta,\kappa_{r,l}\} = \{0.1,10,1\}\kappa_t$.

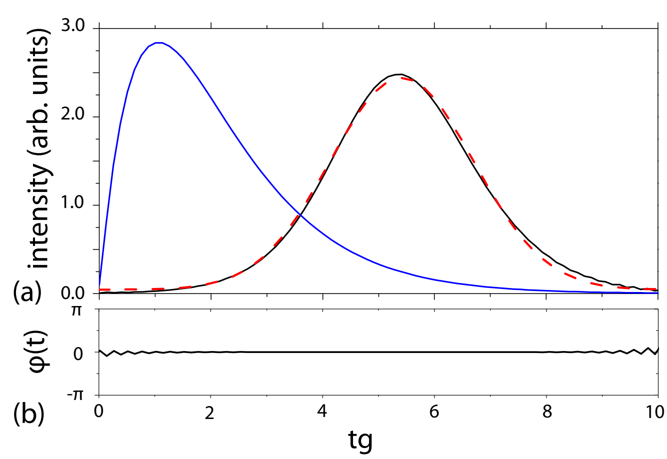

FIG. 3. (Color online) (a) Pulses emitted into the waveguide for $\Delta = const$ (blue) and for $\Delta(t)$ (black). The red dashed line is a Gaussian fit of the emitted symmetric photon pulse. (b) Phase $\varphi(t)$ of the shaped photon pulse versus time. The parameters used are $\{g,\eta,\kappa_{r,l}\} = \{0.1,10,1\}\kappa_t$.

one opts for the switching of Rabi oscillations. the more stringent condition reads $\kappa < g \ll \sqrt{\beta} \ll \eta$. On the other hand, detrimental off-resonant coupling of the QE to modes $\omega_{2,3}$ in case of $\Delta \approx 0$ gives a lower speed limit governed by the ratio g/η, which should be small. However, there also the spontaneous decay of the QE in leaky modes γ eventually may affect the performance. For both the lower and upper tuning speed limits, a large cavity coupling η is desired and should be carefully engineered during the implementation. In the case of the nonlinear detuning used in the simulations shown in Fig. 2, β needs to be replaced by $\dot{\Delta}(t)$. For the chosen detuning $\Delta(t)$, the adiabatic condition is very well satisfied.

To give an example, in a recent experiment [17] the tuning rate is of the order of $\sqrt{\beta} \approx 1/100ps^{-1}$, $\eta \approx 1/5ps^{-1}$ and the quantum dot decay rate about $\gamma_0 \approx 1/1000ps^{-1}$, which satisfies the adiabatic condition. Note that while we focus here on the adiabatic tuning, diabatic effects may open up additional applications of the coupled-cavity system, e.g., for unconventional beam splitters and for the generation of entanglement [23].

IV. IMPLEMENTATION

The proposed scheme can be, in principle, implemented in any coupled-cavity system. Here we discuss two possibilities. In the solid state, various coupled-cavity systems have been realized, such as microdisks [24,25], nanowires [26,27], ring resonators [28], and photonic crystals [29–32]. Here we consider a coupled photonic crystal cavity-quantum dot system. The desired tuning can be implemented by modulation of two laser beams impinging on the control cavities injecting free carriers and thus providing the local refractive index change. The symmetric positive or negative detuning of the two control cavities can be achieved simply by applying a complementary modulation to the two control beams around a central bias value, with commercial optical modulators. The modulation speed is only limited by the free-carrier lifetime in the material and can well exceed the SE time of a weakly coupled quantum dot [17]. By applying electric fields to reduce the free-carrier lifetime down to a few picoseconds, a modulation speed of about 100 GHz can be reached.

To be more precise, we consider a system consisting of three in-line coupled L3 photonic crystal cavities experimentally implemented, e.g., in Ref. [33]. Three-dimensional finite-element simulations are used to determine the local density of optical states (LDOS) D experienced by a QE in the central target cavity depending on its wavelength for different detuning of the control cavities, as shown in Figs. 4(a)–4(c), together with the electric field profiles of the mode with frequency ω_1. The refractive index change can be translated into a frequency or wavelength change by using the expression $\Delta\omega/\omega = -\Delta n/n$. For $\Delta = 0$ the mode ω_1 has no target cavity fraction and hence the dipole in the central cavity interacts only with the modes $\omega_{2,3}$ [Fig. 4(a)]. For intermediate detuning, all three coupled modes have electric field distributions in the target cavity mode, each at its own frequency [Fig. 4(b)]. In case of very large detunings the central cavity mode is completely decoupled and only a single peak is visible in the LDOS spectrum [Fig. 4(c)]. The difference in the loss rates of modes $\omega_{2,3}$ arises from different diffractive out-of-plane

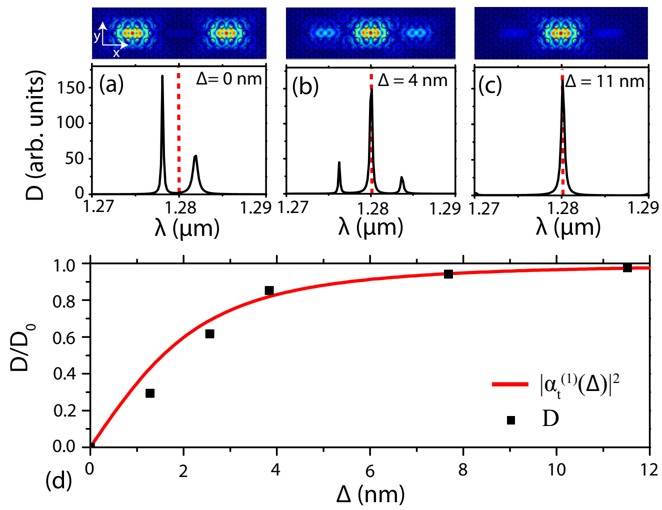

FIG. 4. (Color online) Electric field distributions of the mode at ω_1 and spectra of the local density of optical states D experienced by a dipole in the target cavity for a detuning (a) $\Delta = 0$ nm, (b) $\Delta = 4$ nm, and (c) $\Delta = 11$ nm. The dashed red line indicates the position of the mode with frequency ω_1. (d) Normalized LDOS D/D_0 calculated by finite-element simulations (squares) in comparison to the results from the coupled oscillator model $|\alpha_t^{(1)}|^2$ (red line).

losses [34]. The dependence of the LDOS at frequency ω_1 on the detuning is shown in Fig. 4(d). It is in good agreement with the simple coupled oscillator model $D_t = |\alpha_t^{(1)}|^2 D_0$, where D_0 is the LDOS of the decoupled target cavity mode determined by the finite-element calculations. Once again, the LDOS experienced by the QE can be fully controlled by tuning the control cavities.

An alternative implementation is possible in superconducting circuits [5] due to the ultrahigh-quality resonators as well as long coherence times of qubits. The development of tunable superconducting resonators [35–37] has been the basis for tunable couplers with an unprecedented level of control [38]. This approach is also based on a coupled-cavity system, where the positive or negative frequency detuning of the resonators is controlled by a magnetic flux applied to a superconducting quantum interference device. Thus superconducting circuits represent an ideal platform to realize the present proposal in the microwave regime.

V. SUMMARY AND CONCLUSIONS

Summarizing, the present proposal illustrates that a coupled-cavity system can be used to fully and deterministically control the QE-cavity coupling without inducing a spectral detuning between the QE and the cavity mode. This enables active shaping of single-photon pulses in ultrafast experiments. Finally, we propose two experimental platforms, which can be used to implement the tuning scheme. The present proposal provides a powerful technique to control the SE of QEs and paves the way towards a fully controllable single-photon source. Furthermore, the results show the potential of the coupled-cavity approach for actively controlling the light-matter interaction in the solid state, enabling more advanced application in quantum information processing.

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[1] J. M. Raimond, M. Brune, and S. Haroche, Rev. Mod. Phys. **73**, 565 (2001).

[2] J. Reithmaier, G. Sek, A. Löffler, C. Hofmann, S. Kuhn, S. Reitzenstein, L. Keldysh, V. Kulakovskii, T. Reinecke, and A. Forchel, Nature (London) **432**, 197 (2004).

[3] T. Yoshie, A. Scherer, J. Hendrickson, G. Khitrova, H. Gibbs, G. Rupper, C. Ell, O. Shchekin, and D. Deppe, Nature (London) **432**, 200 (2004).

[4] K. Hennessy, A. Badolato, M. Winger, D. Gerace, M. Atatüre, S. Gulde, S. Fält, E. Hu, and A. Imamoglu, Nature (London) **445**, 896 (2007).

[5] M. Devoret and R. Schoelkopf, Science **339**, 1169 (2013).

[6] A. Shields, Nat. Photonics **1**, 215 (2007).

[7] J. I. Cirac, P. Zoller, H. J. Kimble, and H. Mabuchi, Phys. Rev. Lett. **78**, 3221 (1997).

[8] T. Wilk, S. C. Webster, A. Kuhn, and G. Rempe, Science **317**, 488 (2007).

[9] S. Ritter, C. Nölleke, C. Hahn, A. Reiserer, A. Neuzner, M. Uphoff, M. Mücke, E. Figueroa, J. Bochmann, and G. Rempe, Nature (London) **484**, 195 (2012).

[10] L. C. Andreani, G. Panzarini, and J.-M. Gérard, Phys. Rev. B **60**, 13276 (1999).

[11] A. Högele, S. Seidl, M. Kroner, K. Karrai, R. J. Warburton, B. D. Gerardot, and P. M. Petroff, Phys. Rev. Lett. **93**, 217401 (2004).

[12] A. Faraon, D. Englund, I. Fushman, J. Vučković, N. Stoltz, and P. Petroff, Appl. Phys. Lett. **90**, 213110 (2007).

[13] R. Patel, A. Bennett, I. Farrer, C. Nicoll, D. Ritchie, and A. Shields, Nat. Photonics **4**, 632 (2010).

[14] L. Midolo, F. Pagliano, T. Hoang, T. Xia, F. van Otten, L. Li, E. Linfield, M. Lermer, S. Höfling, and A. Fiore, Appl. Phys. Lett. **101**, 091106 (2012).

[15] R. Trotta, P. Atkinson, J. Plumhof, E. Zallo, R. Rezaev, S. Kumar, S. Baunack, J. Schroeter, A. Rastelli, and O. Schmidt, Adv. Mater. **24**, 2668 (2012).

[16] H. Thyrrestrup, A. Hartsuiker, J.-M. Gérard, and W. L. Vos, Opt. Express **21**, 23130 (2013).

[17] C.-Y. Jin, R. Johne, M. Y. Swinkels, T. B. Hoang, L. Midolo, P. J. van Veldhoven, and A. Fiore, Nat. Nanotechnol. **9**, 886 (2014).

[18] F. Pagliano, Y. Cho, T. Xia, F. van Otten, R. Johne, and A. Fiore, Nat. Commun. **5**, 5786 (2014).

[19] R. Johne and A. Fiore, Phys. Rev. A **84**, 053850 (2011).

[20] R. Johne and A. Fiore, Phys. Rev. A **86**, 063815 (2012).

[21] C. Carroll and F. Hioe, J. Phys. A: Math. Gen. **19**, 2061 (1986).

[22] G.-F. Wang, D.-F. Ye, L.-B. Fu, X.-Z. Chen, and J. Liu, Phys. Rev. A **74**, 033414 (2006).

[23] C. M. Quintana, K. D. Petersson, L. W. McFaul, S. J. Srinivasan, A. A. Houck, and J. R. Petta, Phys. Rev. Lett. **110**, 173603 (2013).

[24] J.-W. Ryu, S.-Y. Lee, and S. W. Kim, Phys. Rev. A **79**, 053858 (2009).

[25] M. Benyoucef, J.-B. Shim, J. Wiersig, and O. Schmidt, Opt. Lett. **36**, 1317 (2011).

[26] M. Bayer, T. Gutbrod, J. P. Reithmaier, A. Forchel, T. L. Reinecke, P. A. Knipp, A. A. Dremin, and V. D. Kulakovskii, Phys. Rev. Lett. **81**, 2582 (1998).

[27] M. Galbiati, L. Ferrier, D. D. Solnyshkov, D. Tanese, E. Wertz, A. Amo, M. Abbarchi, P. Senellart, I. Sagnes, A. Lemaître *et al.*, Phys. Rev. Lett. **108**, 126403 (2012).

[28] D. G. Rabus and M. Hamacher, Photon. Technol. Lett., IEEE **13**, 812 (2001).

[29] M. Notomi, T. Tanabe, A. Shinya, E. Kuramochi, H. Taniyama, S. Mitsugi, and M. Morita, Opt. Exp. **15**, 17458 (2007).

[30] S. Vignolini, F. Riboli, D. S. Wiersma, L. Balet, L. H. Li, M. Francardi, A. Gerardino, A. Fiore, M. Gurioli, and F. Intonti, Appl. Phys. Lett. **96**, 141114 (2010).

[31] F. Intonti, F. Riboli, N. Caselli, M. Abbarchi, S. Vignolini, D. Wiersma, A. Vinattieri, D. Gerace, L. Balet, L. Li *et al.*, Phys. Rev. Lett. **106**, 143901 (2011).

[32] R. Konoike, Y. Sato, Y. Tanaka, T. Asano, and S. Noda, Phys. Rev. B **87**, 165138 (2013).

[33] N. Caselli, F. Riboli, F. La China, A. Gerardino, L. Li, E. H. Linfield, F. Pagliano, A. Fiore, F. Intonti, and M. Gurioli, ACS Photonics **2**, 565 (2015).

[34] K. A. Atlasov, K. F. Karlsson, A. Rudra, B. Dwir, and E. Kapon, Opt. Express **16**, 16255 (2008).

[35] M. Wallquist, V. S. Shumeiko, and G. Wendin, Phys. Rev. B **74**, 224506 (2006).

[36] M. Sandberg, C. Wilson, F. Persson, T. Bauch, G. Johansson, V. Shumeiko, T. Duty, and P. Delsing, Appl. Phys. Lett. **92**, 203501 (2008).

[37] A. Palacios-Laloy, F. Nguyen, F. Mallet, P. Bertet, D. Vion, and D. Esteve, J. Low Temp. Phys. **151**, 1034 (2008).

[38] M. Pierre, I.-M. Svensson, S. R. Sathyamoorthy, G. Johansson, and P. Delsing, Appl. Phys. Lett. **104**, 232604 (2014).