

Accessing strongly coupled systems without compromising them

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The last decades have seen a burst of experimental platforms reaching the so-called strong-coupling regime, where quantum coherent effects dominate over incoherent processes such as dissipation and thermalization. This has allowed us to create highly nontrivial quantum states and put counterintuitive quantum-mechanical effects to test well beyond what the founding fathers of quantum mechanics ever thought possible. The strong-coupling regime comes with certain challenges though: The need for a large isolation makes it difficult to access the system for control or monitoring purposes. In this Letter we propose a way to access such systems through an engineered environment that does not compromise their strong-coupling effects. As a proof of principle, we apply the approach to the photon-blockade effect present in nonlinear resonators or Kerr cavities, optomechanical systems, and the open Jaynes-Cummings model, but argue that the mechanism is quite universal. We also propose an architecture based on superconducting circuits where the required unconventional environment can be implemented, opening the way to the experimental analysis of our ideas.

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Introduction. The precise control of interactions is arguably the landmark of quantum optical systems. It has allowed us to handcraft quantum superposition states able to test the laws of quantum mechanics well beyond what the founding fathers of quantum mechanics ever thought possible. The generation of such states is only possible by accessing the so-called strong-coupling regime, meaning that the quantum-coherent part of the evolution (governed by the Schrödinger equation) induced by the interaction between the system components occurs within a timescale where dissipation and other mechanisms responsible for decoherence are still not relevant. Hence, since the pioneering experiments in cavity quantum electrodynamics [1–3] and trapped ions [4,5], researchers have worked hard to achieve the strong-coupling regime in many other platforms, including cold atoms [6,7], superconducting circuits [8,9], and mechanical devices [10], among others.

Naturally, two strategies can be followed to achieve the strong-coupling regime: enhancing the interactions or

decreasing the decoherence rates. One can try to enhance the interaction by confining the system to smaller volumes, an approach that has led to the advent of modern nanocavities [11,12]. Alternatively, one can find systems with larger dipole moments, a path followed by the artificial atom community of superconducting circuits [9], for example. However, these are features that are not readily available in most experimental platforms. The second route consists in isolating the system incredibly well, but this comes at a high price: It becomes then difficult to access it without compromising its strong coupling.

In this Letter we introduce a way to overcome this last limitation, that is, of connecting the system to an environment that we can use for driving and monitoring purposes, but keeping the system in its strong-coupling regime. Instead of a single environment, our idea uses two environments, and an additional system (twin to the main one, but not necessarily in the strong-coupling regime) that couples to these two environments as well. When the systems couple symmetrically to the environments, a perfectly isolated dark mode appears [13], which inherits the interactions of the system under the right conditions. By allowing for a slightly asymmetric coupling to the environments, this mode acquires a small coupling into our controlled environment, but effectively remaining in the strong-coupling regime. This mechanism is universal, that is, it works for any type of interaction. In order to study it in detail, we focus on the case of a nonlinear or Kerr resonator [14], but analyze also the Jaynes-Cummings and optomechanical

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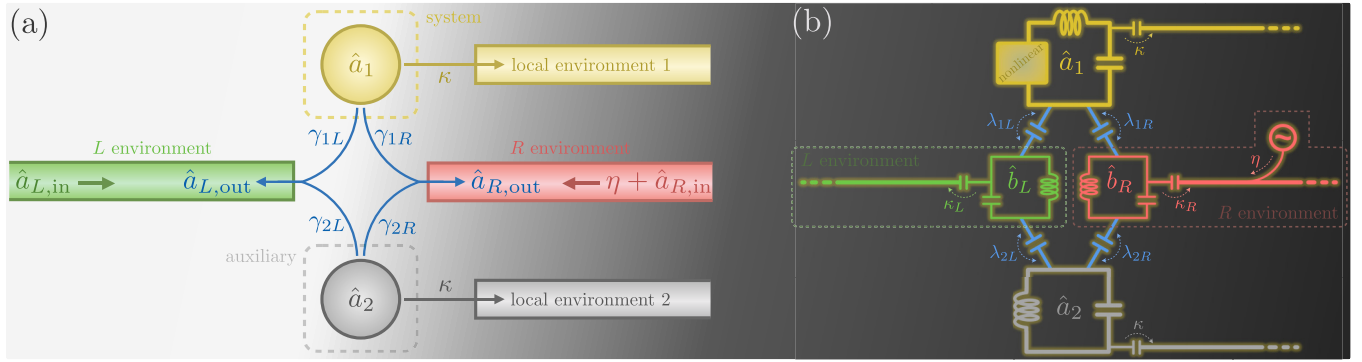


FIG. 1. (a) Sketch of our system, originally formed by one bosonic mode $\hat{a}_1$ presenting a nonlinearity that is strong with respect to its intrinsic decay κ , but weak with respect to the coupling $\gamma_{j\alpha}$ to the environment that we use as input/output channels. We add an auxiliary mode $\hat{a}_2$ and let the excitations of both modes decay collectively to two environments R and L through outputs $\{\hat{a}_{\alpha,out} \propto \sqrt{\gamma_{1\alpha}}\hat{a}_1 + \sqrt{\gamma_{2\alpha}}\hat{a}_2\}_{\alpha=R,L}$, while the environments are driven by input vacuum fluctuations $\{\hat{a}_{\alpha,in}\}_{\alpha=R,L}$ plus a coherent drive of irradiance η^2 for one environment. As we explain in the text, an asymmetric decay of the modes into the environments ($\gamma_{2R} \approx \gamma_{2L} \gg \gamma_{1R} = \gamma_{1L}$) allows the system to reach an effective strong-coupling regime, as evidenced by strong and robust antibunching of the output. (b) Sketch of the experimental proposal. Two LC superconducting circuits (one including some nonlinear element) form the main system, and exchange excitations at rates $\lambda_{j\alpha}$ with two auxiliary LC circuits with corresponding operators $\{\hat{b}_{\alpha}\}_{\alpha=R,L}$, which in turn decay to two transmission lines at rates $\kappa_{\alpha} \gg \lambda_{j\alpha}$. The same idea can be implemented in the optical domain, for example by replacing the LC circuits with whispering gallery mode resonators or photonic crystal cavities, and the transmission lines by photonic fibers.

systems. To demonstrate the strong-coupling regime, we analyze whether the photon-blockade effect is present in the system: It turns a weak input coherent drive into an output field with strong photon antibunching, all while keeping large reflection probabilities. We propose a specific architecture based on superconducting circuits where our ideas can be explored.

Note that our mechanism is completely different from the so-called unconventional photon-blockade effect [15–18], which allows generating antibunched photon statistics even within the weak-coupling regime by exploiting the coherent tunneling of quanta within two resonators, but requires fine tuning of the nonlinearity.

Model. Our starting point is a system in the strong-coupling regime, so that quantum coherent dynamics can compete with incoherent processes such as loss. As a prototype, we consider a bosonic mode with annihilation operator $\hat{a}_1$, subjected to a nonlinear Kerr Hamiltonian characterized by a rate g (see below for details), and with an *intrinsic* loss rate $\kappa < g$ induced by environments that are not accessible to experiments. In order to access the system, one needs to couple it to an experimentally controllable environment that will induce loss at some rate γ . A problem arises when $\gamma > g$, which is unavoidable in some systems where g cannot be very large, so the accessing process would destroy the strong-coupling regime. In the following, we propose a way to solve this issue.

We introduce an auxiliary system which is a copy of the original one except for the Kerr nonlinearity that we set to zero in this one for simplicity. Denoting by $\hat{a}_2$ the annihilation operator of the auxiliary system, we couple the two modes to two experimentally accessible environments denoted by L and R via the collective jump operators $\{\hat{J}_{\alpha} = \sqrt{\gamma_{1\alpha}}\hat{a}_1 + \sqrt{\gamma_{2\alpha}}\hat{a}_2\}_{\alpha=R,L}$, where $\gamma_{j,\alpha} > g$. Hence, rather than considering only independent decay channels for each bosonic mode, we consider a collective coherent decay that will allow us to exploit quantum interference effects. We later put

forward concrete architectures based on currently available superconducting-circuit technologies where our ideas should be readily implementable [see Fig. 1(b)].

Under standard Born-Markov conditions [19–21] and working in a picture rotating at the driving frequency, detuned by Δ with respect to the modes' intrinsic frequencies, the evolution of the state of the system $\hat{\rho}$ is governed by the master equation

$$\partial_t \hat{\rho} = -i[\hat{H}, \hat{\rho}] + \sum_{\alpha=L,R} \mathcal{D}_{J_{\alpha}}[\hat{\rho}] + \sum_{j=1}^2 \kappa \mathcal{D}_{a_j}[\hat{\rho}], \quad (1)$$

where we have defined the Lindblad dissipator

$$\mathcal{D}_J[\hat{\rho}] = 2\hat{J}\hat{\rho}\hat{J}^{\dagger} - \hat{J}^{\dagger}\hat{J}\hat{\rho} - \hat{\rho}\hat{J}^{\dagger}\hat{J}, \quad (2)$$

and the Hamiltonian (divided by $\hbar$) reads

$$\hat{H} = \Delta(\hat{a}_1^{\dagger}\hat{a}_1 + \hat{a}_2^{\dagger}\hat{a}_2) + g\hat{a}_1^{\dagger}\hat{a}_1^2 + \sqrt{2}\eta(\hat{J}_R + \hat{J}_R^{\dagger}), \quad (3)$$

where we have assumed that the system is driven from the R environment with flux η^2 (quanta per unit time). In addition, the fields coming out of the system into the collective environments are characterized by output operators $\hat{a}_{\alpha,out}(t) = \eta\delta_{\alpha,R} + \hat{a}_{\alpha,in}(t) - \sqrt{2}i\hat{J}_{\alpha}(t)$, where $\hat{a}_{\alpha,in}(t)$ annihilates the vacuum state of the environmental modes.

As explained above, we consider the situation in which strong coupling has been achieved with respect to the intrinsic dissipative channels, while the coupling remains weak with respect to the engineered environments used as input/output channels, that is, $\kappa < g \ll \gamma_{j\alpha}$. In this Letter we show how the nonlinear effects that are characteristic of the strong-coupling regime can still be observed by exploiting quantum interference effects coming from the collective character of the jump operators $\hat{J}_{\alpha}$, allowing us to access the system without compromising its strong-coupling regime.

Intuitive picture. In order to understand the essence of the mechanism, it is best to consider the $\kappa = 0$ case, and

perform a canonical transformation to the basis of modes that diagonalize the dissipative part of the master equation,

$$\hat{A}_+ = \hat{a}_1 \cos \theta + \hat{a}_2 \sin \theta, \quad (4a)$$

$$\hat{A}_- = \hat{a}_1 \sin \theta - \hat{a}_2 \cos \theta, \quad (4b)$$

with $\theta = \arctan[2\gamma_{12}/(\gamma_- + \sqrt{\gamma_-^2 + 4\gamma_{12}^2})] \in [0, \pi]$, where $\gamma_j = \gamma_{jR} + \gamma_{jL}$, $\gamma_- = \gamma_1 - \gamma_2$, and $\gamma_{12} = \sqrt{\gamma_{1R}\gamma_{2R} + \gamma_{1L}\gamma_{2L}}$. In terms of these new “normal” modes, which obey canonical commutation relations unlike the jump operators $\hat{J}_\alpha$, we show [22] that the dissipative part of (1) takes the form

$$\sum_{\alpha=L,R} \mathcal{D}_{J_\alpha}[\hat{\rho}] = \sum_{n=\pm} \Gamma_n \mathcal{D}_{A_n}[\hat{\rho}], \quad (5)$$

with normal damping rates $2\Gamma_\pm = \gamma_1 + \gamma_2 \pm \sqrt{\gamma_-^2 + 4\gamma_{12}^2}$.

Interestingly, when $\gamma_{jR} = \gamma_{jL}$ (that is, when the modes couple in the same way to both environments), mode $\hat{A}_+$ gathers all the damping, $\Gamma_+ = \gamma_1 + \gamma_2$, while mode $\hat{A}_-$ effectively becomes dark or undamped, $\Gamma_- = 0$. If, in addition, we work in the $\gamma_2 \gg \gamma_1$ regime, $\theta \approx \pi/2$, and the nonlinear term $g\hat{a}_1^\dagger \hat{a}_1^2$ is dominated by the $g_- \hat{A}_+^\dagger \hat{A}_-^2$ contribution, with $g_- = g \sin^4 \theta$. Hence, while mode $\hat{A}_+$ is still in the weak-coupling regime, $\Gamma_+ \gg g$, this is not the case for mode $\hat{A}_-$, which in this extreme limit behaves as a dissipativeless anharmonic or nonlinear oscillator.

This provides the intuition that by working close to this regime, specifically keeping $y \equiv \gamma_1/\gamma_2 \ll 1$, but allowing for a slightly asymmetric coupling of the modes to the environments, mode $\hat{A}_-$ will effectively enter the strong-coupling regime with $g_- \gg \Gamma_-$. Specifically, in the following we keep mode-1's couplings symmetric for simplicity, $\gamma_{1R} = \gamma_{1L} = \gamma_1/2$, and define the asymmetry parameter $x \equiv (\gamma_{2R} - \gamma_{2L})/\gamma_2$ for the second mode, with $|x| \ll 1$. To the leading orders in x and y [22], the $g_- \gg \Gamma_-$ condition can be recast as $|x| \ll 2\sqrt{g/\gamma_1}$, which sets an upper bound to the asymmetry x that we can introduce.

Interestingly, we show in the Supplemental Material [22] that in this regime an excitation in mode $\hat{a}_1$ is stored during a time $(\Gamma_- + \kappa)^{-1}$, which is much longer than the one characteristic of its interaction with the environment, γ_1^{-1} . This type of protection can have applications in protocols where quantum information storage is needed.

In the following, it is convenient to define $\sqrt{\Gamma_{+\alpha}} = \sqrt{\gamma_{1\alpha}} \cos \theta + \sqrt{\gamma_{2\alpha}} \sin \theta$ and $\sqrt{\Gamma_{-\alpha}} = \sqrt{\gamma_{1\alpha}} \sin \theta - \sqrt{\gamma_{2\alpha}} \cos \theta$, which provide the relation between the collective jump operators and the normal modes, $\hat{J}_\alpha = \sqrt{\Gamma_{+\alpha}} \hat{A}_+ + \sqrt{\Gamma_{-\alpha}} \hat{A}_-$. Note that $\Gamma_\pm = \Gamma_{\pm R} + \Gamma_{\pm L}$.

Effective strong-coupling phenomena. In order to go beyond the qualitative picture described above based on the normal modes $\hat{A}_\pm$, we now show that effects of the nonlinearity are observable through the output field of the R environment, which shows antibunched quanta statistics, while allowing for near-perfect reflection of the incident power (so-called photon-blockade effect). We gain analytical insight by considering the weak-driving regime, performing an expansion in powers of η [22], equivalent to scattering theory formalism [23–25]. The leading order of this expansion should be accurate as long as the number of excitations in the system is small. Neglecting nonlinear effects, the excitation

number is dominated by $N_\eta \equiv 2\Gamma_- \eta^2 / \Gamma_-^2$, that is, the ratio between the driving and damping of the nearly dark mode $\hat{A}_-$, so the leading order of scattering theory is expected to work as long as $N_\eta \ll 1$. In the following we then use N_η instead of η for our analysis.

Consider first the reflection amplitude, which to the lowest order in η we show in the Supplemental Material [22] to read

$$\mathcal{R} = \lim_{t \rightarrow \infty} \frac{\langle \hat{a}_{R,\text{out}}(t) \rangle}{\eta} \approx 1 + 2i \sum_{n=\pm} \frac{\Gamma_{nR}}{\Delta - i(\Gamma_n + \kappa)}. \quad (6)$$

Using the definitions of the different rates, it is easy to show that for a resonant drive ($\Delta = 0$), we obtain $\mathcal{R} = (\Lambda x - 1)/(\Lambda + 1)$, with $\Lambda = \kappa/\Gamma_- \approx 4\kappa/\gamma_2 x^2 y$ [22]. Hence, in the absence of local damping ($\Lambda = 0$), all the energy fed into the system in the form of an input coherent field is transferred back to the R environment, since $\mathcal{R} = -1$. However, the quantum statistics of this reflected field are dramatically changed by the interaction with the system. In particular, we show in the Supplemental Material [22] that, again to the lowest order in η , its normalized second-order correlation function takes the form

$$g^{(2)}(\tau) \equiv \lim_{t \rightarrow \infty} \frac{\langle \hat{a}_{R,\text{out}}^\dagger(t) \hat{n}_{R,\text{out}}(t + \tau) \hat{a}_{R,\text{out}}(t) \rangle}{\langle \hat{n}_{R,\text{out}}(t) \rangle^2} \approx \left[\left(\frac{1 - \Lambda x}{1 + \Lambda} \right)^2 - \frac{y}{1 + \Lambda} e^{-\tilde{\Gamma}_+ \tau} - \frac{1 - 2x}{(1 + \Lambda)^2} e^{-\tilde{\Gamma}_- \tau} \right]^2, \quad (7)$$

where $\hat{n}_{R,\text{out}}(t) = \hat{a}_{R,\text{out}}^\dagger(t) \hat{a}_{R,\text{out}}(t)$, $\tilde{\Gamma}_n = \Gamma_n + \kappa_0$, and we have expanded the final expression to first order in x and y , assuming $\Gamma_-/g_- \ll 1$ as well, which is granted in our setup. For $\kappa = 0$, this expression shows that $g^{(2)}(0) = (2x - y)^2$, signaling strong antibunching over a timescale determined by Γ_-^{-1} . Note that this expression is independent of the coupling g , and hence, such a nonlinear effect only requires the broad conditions $(x, y, \Gamma_-/g_-) \ll 1$, which are naturally satisfied in our construction, unlike previously proposed mechanisms [15–18] that require fine tuning of the parameters. In addition, it is worth remarking that our proposal generates a photon-blockade effect in the reflected light rather than in the transmitted one, which is the most usual case.

In Fig. 2 we confront this analytical approximation with the exact result found numerically [22,26] from the master equation (1). We set $\kappa = 0$ temporarily to remove the limitations set by the intrinsic loss, so that we can estimate well the limitations imposed by our engineered environment; below we discuss the $\kappa \neq 0$ case. $g^{(2)}(\tau)$ is shown as a function of $\Gamma_- \tau$ in Fig. 2(a) for different values of N_η . As expected, we find good agreement when N_η is small. Remarkably, even for large N_η antibunching is still present. In particular, we characterize the antibunching as a function of N_η in Fig. 2(b), where we plot $g^{(2)}(0)$, the smallest time delay $\tilde{\tau}$ for which $g^{(2)}(\tilde{\tau}) = 1/e$ (characterizing the antibunching timescale), and the output photon flux $\bar{n}_{\text{out}} = \lim_{t \rightarrow \infty} \langle \hat{n}_{R,\text{out}}(t) \rangle$. If used as a single-photon source [27,28], we see that our system is optimized for $N_\eta \approx 1$, with $\bar{n}_{\text{out}} \approx \Gamma_-/2$, $g^{(2)}(0) \approx 10^{-5}$, and $\tilde{\tau} \approx \Gamma_-^{-1}$, while the triggering pulses used as input need to

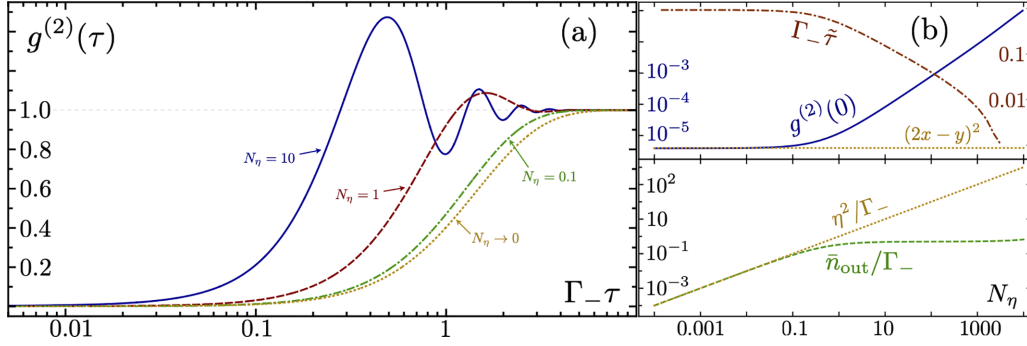


FIG. 2. (a) Second-order correlation function $g^{(2)}(\tau)$ as a function of the normalized time delay $\Gamma_- \tau$ for different values of the normalized input photon flux $N_\eta = 2\Gamma_{-R}\eta^2/\Gamma_-^2$. Strong antibunching is appreciated, even when $N_\eta > 1$, where our analytical results based on scattering theory (dotted yellow line) cease to apply. (b) characterizes the antibunching as a function of N_η , showing $g^{(2)}(0)$ in solid blue, the normalized output photon flux $\tilde{n}_{\text{out}}/\Gamma_-$ in dashed green, and in dashed-dotted green the normalized time delay $\Gamma_- \tilde{\tau}$ that it takes for the second-order correlation function to grow up to $g^{(2)}(\tilde{\tau}) = 1/e$, providing a measure for the antibunching timescale. The dotted yellow curves are the analytic results obtained from scattering theory. Note that on the upper panel the left ticks refer to $g^{(2)}(0)$ while the right ones to $\Gamma_- \tilde{\tau}$. In all plots we have chosen $y = 0.01$, $g = 4 \times 10^{-4}\gamma_2$, $x = 0.02\sqrt{g/\gamma_1} = 0.004$, and $\kappa = 0$.

have a width shorter than Γ_-^{-1} , which is the width of both the reflection peak and the $g^{(2)}(\tau)$ dip.

Note that in the case of nonvanishing intrinsic damping $\kappa \neq 0$, it is best to set the parameters such that Λ is not very large ($\Gamma_- \sim \kappa$), so that the reflection $\mathcal{R}$ remains reasonably large, while still achieving strong antibunching $g^{(2)}(0) \ll 1$. Below we show that this and all the previous conditions are feasible with superconducting-circuit architectures.

Universality of the mechanism. Importantly, we have checked [22] that the photon blockade found in our system is not specific to the Kerr nonlinearity. In particular, we have considered more intricate nonlinearities such as coupling mode $\hat{a}_1$ to a two-level system (Jaynes-Cummings model [29–31]) or to a mechanical oscillator [32], finding that the output through our collective decay channel presents the photon-blockade effects expected for such systems [22].

Implementation. We propose now a generic way in which the collective decay required in our setup can be implemented, which we support with specific parameters of current experimental platforms. The basic idea is sketched in Fig. 1(b) through a superconducting-circuit architecture. The original modes of our model, $\hat{a}_j$, correspond to *LC* circuits, where the first one has an additional additional nonlinear element (e.g., a Josephson junction for Kerr nonlinearity [9], but Jaynes-Cummings [9] or optomechanical [10,33] nonlinearities are also available in such platform). These circuits are coupled to two identical auxiliary *LC* circuits with annihilation operators $\{\hat{b}_\alpha\}_{\alpha=R,L}$, which in turn decay each to a transmission line at rate $\{\kappa_\alpha\}_{\alpha=R,L}$. In addition, the auxiliary circuit *R* is resonantly driven with an external generator. In a picture where all modes rotate at the driving frequency, the master equation describing the dynamics of the state of the whole system, denoted by $\hat{\rho}$, reads

$$\partial_t \hat{\rho} = -i[\hat{H}, \hat{\rho}] + \sum_{\alpha=R,L} \kappa_\alpha \mathcal{D}_{b_\alpha}[\hat{\rho}] + \sum_{j=1}^2 \kappa \mathcal{D}_{a_j}[\hat{\rho}], \quad (8)$$

with

$$\begin{aligned} \hat{H} = & \sum_{j=1}^2 \Delta \hat{a}_j^\dagger \hat{a}_j + \hat{H}_{\text{NL}} + i\sqrt{2\kappa_R}\eta(\hat{b}_R^\dagger - \hat{b}_R) \\ & + \sum_{j=1}^2 \sum_{\alpha=R,L} \lambda_{j\alpha}(\hat{a}_j \hat{b}_\alpha^\dagger + \hat{a}_j^\dagger \hat{b}_\alpha), \end{aligned} \quad (9)$$

where $\hat{H}_{\text{NL}}$ includes the nonlinear processes in mode $\hat{a}_1$. When $\kappa_\alpha \gg \lambda_{j\alpha}$, the auxiliary modes can be adiabatically eliminated as we show in the Supplemental Material [22], leading precisely to our original master equation (1), with $\gamma_{j\alpha} = \lambda_{j\alpha}^2/\kappa_\alpha$. Hence, we see that the lossy auxiliary modes act as intermediate sinks where the main modes interfere before decaying to the transmission lines.

In order to prove the feasibility of our proposal with current experimental platforms, let us consider some specific parameters, but keeping in mind that our ideas do not require fine tuning of these. We consider *LC* circuits with 4 GHz resonance frequency (all frequencies are in 2π units in the following), and consider low-*Q* auxiliary circuits with $\kappa_\alpha = 200$ MHz decay rates, and couplings $\lambda_{1\alpha} = 2$ MHz, $\lambda_{2R} = 20.88$ MHz, and $\lambda_{2L} = 19.08$ MHz. These values are common in this platform [9], and lead to $y \approx 0.01$, $x \approx 0.05$, and $\gamma_2 \approx 4$ MHz. Assuming then a state-of-the-art 10^8 quality factor for the *LC* circuits, we obtain $\kappa = 40$ Hz, while the nonlinearity g can easily reach kHz in these platforms, so that we stay in the $\kappa < g \ll \gamma_{j\alpha}$ regime. For these parameters, our theory predicts strong antibunching $g^{(2)}(0) \approx 10^{-3}$ with a large reflection $|\mathcal{R}| \approx 0.64$.

Conclusions. In this Letter we have presented a way to externally access systems that are in the strong-coupling regime (with respect to their intrinsic loss), without compromising them. The idea relies on the addition of an auxiliary system, which together with the main system is coupled asymmetrically to two environments in an unconventional way: They decay collectively, rather than independently. Using as an example the photon-blockade effect of a Kerr

resonator, we have shown that the mechanism works, and argued that it is universal, in the sense that it works for any type of nonlinearity. We have finally offered a generic way in which the unconventional decays can be engineered, showing that superconducting-circuit technologies are specially suited for the observation of the phenomena predicted in this work. We hope that our proposal will allow more experimental platforms to reach the strong-coupling needed not only for the generation of single photons, but also for other applications such as low-threshold lasing or switching [34].

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- [1] J. M. Raimond, M. Brune, and S. Haroche, Manipulating quantum entanglement with atoms and photons in a cavity, *Rev. Mod. Phys.* **73**, 565 (2001).
 - [2] R. Miller, T. E. Northup, K. M. Birnbaum, A. Boca, A. D. Boozer, and H. J. Kimble, Trapped atoms in cavity QED: coupling quantized light and matter, *J. Phys. B: At., Mol. Opt. Phys.* **38**, S551 (2005).
 - [3] H. Walther, B. T. H. Varcoe, B.-G. Englert, and T. Becker, Cavity quantum electrodynamics, *Rep. Prog. Phys.* **69**, 1325 (2006).
 - [4] D. Leibfried, R. Blatt, C. Monroe, and D. Wineland, Quantum dynamics of single trapped ions, *Rev. Mod. Phys.* **75**, 281 (2003).
 - [5] C. Schneider, D. Porras, and T. Schaetz, Experimental quantum simulations of many-body physics with trapped ions, *Rep. Prog. Phys.* **75**, 024401 (2012).
 - [6] D. Jaksch and P. Zoller, The cold atom Hubbard toolbox, *Ann. Phys. (NY)* **315**, 52 (2005).
 - [7] I. Bloch, J. Dalibard, and W. Zwerger, Many-body physics with ultracold gases, *Rev. Mod. Phys.* **80**, 885 (2008).
 - [8] M. H. Devoret and R. J. Schoelkopf, Superconducting circuits for quantum information: An outlook, *Science* **339**, 1169 (2013).
 - [9] A. Blais, A. L. Grimsmo, S. M. Girvin, and A. Wallraff, Circuit quantum electrodynamics, *Rev. Mod. Phys.* **93**, 025005 (2021).
 - [10] M. Aspelmeyer, T. J. Kippenberg, and F. Marquardt, Cavity optomechanics, *Rev. Mod. Phys.* **86**, 1391 (2014).
 - [11] S. Hu and S. M. Weiss, Design of photonic crystal cavities for extreme light concentration, *ACS Photonics* **3**, 1647 (2016).
 - [12] H. Choi, M. Heuck, and D. Englund, Self-similar nanocavity design with ultrasmall mode volume for single-photon nonlinearities, *Phys. Rev. Lett.* **118**, 223605 (2017).
 - [13] K. Lalumière, B. C. Sanders, A. F. van Loo, A. Fedorov, A. Wallraff, and A. Blais, Input-output theory for waveguide QED with an ensemble of inhomogeneous atoms, *Phys. Rev. A* **88**, 043806 (2013).
 - [14] A. Imamoglu, H. Schmidt, G. Woods, and M. Deutsch, Strongly interacting photons in a nonlinear cavity, *Phys. Rev. Lett.* **79**, 1467 (1997).
 - [15] T. C. H. Liew and V. Savona, Single photons from coupled quantum modes, *Phys. Rev. Lett.* **104**, 183601 (2010).
 - [16] M. Bamba, A. Imamoglu, I. Carusotto, and C. Ciuti, Origin of strong photon antibunching in weakly nonlinear photonic molecules, *Phys. Rev. A* **83**, 021802(R) (2011).
 - [17] H. Flayac, D. Gerace, and V. Savona, An all-silicon single-photon source by unconventional photon blockade, *Sci. Rep.* **5**, 11223 (2015).
 - [18] H. Flayac and V. Savona, Single photons from dissipation in coupled cavities, *Phys. Rev. A* **94**, 013815 (2016).
 - [19] C. W. Gardiner and P. Zoller, *Quantum Noise* (Springer, New York, 2004).
 - [20] I. de Vega and D. Alonso, Dynamics of non-Markovian open quantum systems, *Rev. Mod. Phys.* **89**, 015001 (2017).
 - [21] C. Navarrete-Benlloch, Introduction to quantum optics, [arXiv:2203.13206](https://arxiv.org/abs/2203.13206).
 - [22] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/PhysRevResearch.7.L022024> for details about the diagonalization of the dissipative terms, the lifetime of single excitations, the analytical calculations in the weak-driving regime, the numerical simulations, the results for the Jaynes-Cummings and optomechanical couplings, and the elimination of the auxiliary modes in order to generate the model we seek.
 - [23] Y. Chang, A. González-Tudela, C. Sánchez Muñoz, C. Navarrete-Benlloch, and T. Shi, Deterministic down-converter and continuous photon-pair source within the bad-cavity limit, *Phys. Rev. Lett.* **117**, 203602 (2016).
 - [24] T. Caneva, M. T. Manzoni, T. Shi, J. S. Douglas, J. I. Cirac, and D. E. Chang, Quantum dynamics of propagating photons with strong interactions: A generalized input-output formalism, *New J. Phys.* **17**, 113001 (2015).
 - [25] T. Shi, D. E. Chang, and J. I. Cirac, Multiphoton-scattering theory and generalized master equations, *Phys. Rev. A* **92**, 053834 (2015).
 - [26] C. Navarrete-Benlloch, Open systems dynamics: Simulating master equations in the computer, [arXiv:1504.05266](https://arxiv.org/abs/1504.05266).
 - [27] H. Mabuchi, Qubit limit of cavity nonlinear optics, *Phys. Rev. A* **85**, 015806 (2012).
 - [28] M. E. Reimer and C. Cher, The quest for a perfect single-photon source, *Nat. Photon.* **13**, 734 (2019).
 - [29] L. Tian and H. J. Carmichael, Quantum trajectory simulations of two-state behavior in an optical cavity containing one atom, *Phys. Rev. A* **46**, R6801 (1992).

- [30] K. M. Birnbaum, A. Boca, R. Miller, A. D. Boozer, T. E. Northup, and H. J. Kimble, Photon blockade in an optical cavity with one trapped atom, [Nature \(London\) **436**, 87 \(2005\)](#).
- [31] T. Shi, S. Fan, and C. P. Sun, Two-photon transport in a waveguide coupled to a cavity in a two-level system, [Phys. Rev. A **84**, 063803 \(2011\)](#).
- [32] P. Rabl, Photon blockade effect in optomechanical systems, [Phys. Rev. Lett. **107**, 063601 \(2011\)](#).
- [33] J. D. Teufel, T. Donner, D. Li, J. W. Harlow, M. S. Allman, K. Cicak, A. J. Sirois, J. D. Whittaker, K. W. Lehnert, and R. W. Simmonds, Sideband cooling of micromechanical motion to the quantum ground state, [Nature \(London\) **475**, 359 \(2011\)](#).
- [34] D. S. Dovzhenko, S. V. Ryabchuk, Y. P. Rakovich, and I. R. Nabiev, Light-matter interaction in the strong coupling regime: Configurations, conditions, and applications, [Nanoscale **10**, 3589 \(2018\)](#).