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Bistability in microcavities with incoherent optical or electrical excitation

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I. INTRODUCTION

Bistability and hysteresis are fundamental properties of resonant nonlinear optical systems with applications in optical memory elements and transistors. Optical nonlinearity ultimately relies on the coupling of light to electronic degrees of freedom, which mediate effective interactions between photons. With an ever-growing effort devoted to the enhancement of light-matter coupling, high nonlinearity allows bistability to be routinely observed in a variety of systems.

One example is the strong light-matter coupling in high-quality factor semiconductor microcavities containing quantum well excitons [1]. The resulting exciton-polaritons exhibit bistability, when driven resonantly and coherently [2–4], i.e., with a pumping laser tuned to the energy of the exciton-polariton quasiparticle. This allows a variety of related effects to occur, including driven superfluidity [5], the suppression of disorder [6], and the formation of various structures: spin patterns [7–9], solitons [10–12], and vortex lattices [13]. Controlled switching has been observed in multimode systems, based on the overlapping of states with different energies and wave vectors [14] or the spin degree of freedom [15–17].

The advantages of exciton-polaritons include their ultrafast velocity and strength of nonlinear interactions, which was reported to result in ultrafast switching typically in the picosecond range using low (milliwatt) power in a micrometer-sized area [17]. These characteristics are very competitive with optical switching in other systems, such as microring resonators [18], organic waveguides [19], or photonic crystals [20]. Exciton-polariton bistability also underpins schemes for photonic circuits [21], where a complete logical architecture can be constructed [22]. In addition, the sensitivity of excitons to electric fields could be exploited in hybrid electro-optic devices [23] combining both photonics and electronic interfaces. The further development of hybrid polaritonic devices requires electrically injected microcavities [24–26], recently shown to enable polariton lasing [27]. However, all of the polaritonic devices based on bistability demonstrated to date require a coherent excitation, implying that they must be coupled with a laser light that is resonant or near-resonant with the exciton-polaritons.

In this paper, we propose a mechanism of optical bistability in a semiconductor microcavity compatible with incoherent nonresonant pumping. We consider the recently developed subwavelength grating type microcavity, reported by Zhang et al. [28] and illustrated in Fig. 1(a). In such a cavity, light is only confined if it has a specific linear polarization (x) such that linearly polarized exciton-polaritons are formed, while a cross-linearly polarized (y) exciton mode remains uncoupled to light. We develop evolution equations for occupation numbers and correlators in this system, using a master equation approach [29] accounting for incoherent pumping and dissipation. First, we consider an optical incoherent excitation of an exciton mode, which could be polarized in the y direction so as to avoid direct excitation of exciton-polariton modes. Second, we consider the electrical injection of excitons, leading to asymmetric pumping of all three modes. In both cases, excitons undergo elastic pair scattering into upper and lower polariton states, in analogy to various intrabranch [30,31] and interbranch [32–34] processes typically studied in resonantly excited microcavities. This leads to bistable behavior of the system under certain pumping conditions, which we expect to be useful for future polaritonic devices able to work under optical and electrical incoherent pumping.

II. THE MODEL

Introducing the field operators of cavity photons $\hat{\phi}$ and excitons $\hat{\chi}$, the Hamiltonian of the system can be written as

$$\hat{H} = E_C \hat{\phi}^{\dagger} \hat{\phi} + E_X (\hat{\chi}^{\dagger}_- \hat{\chi}_- + \hat{\chi}^{\dagger}_+ \hat{\chi}_+) + V (\hat{\chi}^{\dagger}_- \hat{\phi} + \hat{\phi}^{\dagger} \hat{\chi}_+)$$
$$+ \alpha_1 (\hat{\chi}^{\dagger}_- \hat{\chi}^{\dagger}_+ \hat{\chi}_- + \hat{\chi}^{\dagger}_- \hat{\chi}^{\dagger}_+ \hat{\chi}_+) + 2\alpha_2 \hat{\chi}^{\dagger}_+ \hat{\chi}^{\dagger}_+ \hat{\chi}^{\dagger}_- \hat{\chi}^{-}.$$

(1)

where $x, y, +, -$ and $-$ subscripts denote linear, cross-linear, circular, and cross-circular polarizations, respectively. We consider spin degenerate exciton states with energy $E_X$ and a linearly polarized cavity mode with energy $E_C$. The exciton-photon coupling constant $V$ couples only $x$ polarized states. Nonlinear interactions between excitons with parallel and antiparallel spins are described by the parameters $\alpha_1$ and $\alpha_2$, respectively. The circularly polarized states can be rewritten in terms of linearly polarized states, $\hat{\chi}_\pm = (\hat{\chi}_x \pm i \hat{\chi}_y)/\sqrt{2}$. Introducing lower and upper $x$ polarized polariton states, $\hat{\phi}_x = C \hat{\psi}_L - X \hat{\psi}_U$ and $\hat{\chi}_x = X \hat{\psi}_L + C \hat{\psi}_U$, we can diagonalize the
The parametric scattering condition subwavelength grating at the top. This allows confinement of the linear part of the Hamiltonian:

\[ \hat{\mathcal{H}} = E_L \hat{\psi}_L^\dagger \hat{\psi}_L + E_U \hat{\psi}_U^\dagger \hat{\psi}_U + \frac{(\alpha_1 + \alpha_2)}{2}(X^4 \hat{\psi}_L^\dagger \hat{\psi}_L \hat{\psi}_U + C^4 \hat{\psi}_L^\dagger \hat{\psi}_L \hat{\psi}_U) \]

\[ + 4X^2 C^2 \hat{\psi}_L^\dagger \hat{\psi}_L \hat{\psi}_U + \chi \hat{\chi} \hat{\chi} \hat{\chi} + 2 \alpha_1(X^2 \hat{\psi}_L^\dagger \hat{\psi}_L + C^2 \hat{\psi}_L^\dagger \hat{\psi}_L) \hat{\chi} \hat{\chi} \hat{\chi} + E_X \hat{\chi} \hat{\chi} \hat{\chi}. \]

The polariton energies are given by \( E_{U,L} = (E_C + E_X)/2 \pm \sqrt{\delta^2 + 4V^2}/2 \), where \( \delta = E_C - E_X \) denotes a detuning between photon and exciton modes. The corresponding Hopfield coefficients \( X \) and \( C \) can be defined using relations \( X^2, C^2 = \pm (\delta/\sqrt{\delta^2 + 4V^2})/2 \).

Given the Hamiltonian, the system evolution follows the master equation for the density matrix \( \rho \):

\[ \frac{d\rho}{dt} = -\frac{i}{\hbar}[\hat{\mathcal{H}}, \rho] + \sum_{a=x,y} P_a (\hat{a} \rho \hat{a}^\dagger - \hat{a}^\dagger \hat{a} \rho - \rho \hat{a}^\dagger \hat{a}) + \sum_{a=x,y} \Gamma_a (2 \hat{a} \rho \hat{a}^\dagger - \hat{a}^\dagger \hat{a} \rho - \rho \hat{a}^\dagger \hat{a}), \]

where we defined \( \hat{\chi} \equiv \hat{\chi}_y \), \( P_i \) \( (i = X,L,U) \) are the rates of incoherent pumping [35,36] for different modes, and we accounted for different decay rates \( \Gamma_x, \Gamma_U = X^2 \Gamma_C + \Gamma^3 X, \) and \( \Gamma_L = C^2 \Gamma_C + X^2 \Gamma_X, \) corresponding to \( \chi, \psi_U, \) and \( \psi_L \) modes. \( \Gamma_C \) denotes the cavity photon linewidth. We will consider two cases: (1) incoherent optical polarized pump \( (P_X > 0, P_{L,U} = 0) \), where only \( \chi_y \) excitons are pumped, and (2) electrical pump \( (P_{X,L,U} > 0) \), where all three modes are asymmetrically populated. The density matrix allows the calculation of observable quantities such as \( \langle \hat{O} \rangle = \text{tr}(\hat{O} \rho) \).

From Eq. (3), we can thus derive a set of equations for the population numbers of polaritons \( N_{L,U} = \langle \hat{\psi}_L^\dagger \hat{\psi}_L \rangle \) and excitons \( N_X = \langle \hat{\chi} \hat{\chi} \rangle \) (see Ref. [37]):

\[ \frac{dN_L}{dt} = -2\alpha_+ X \Gamma_C m \langle A \rangle - \frac{\Gamma_L}{\hbar} N_L + P_L, \]

\[ \frac{dN_U}{dt} = -2\alpha_- X \Gamma_C m \langle A \rangle - \frac{\Gamma_U}{\hbar} N_U + P_U, \]

\[ \frac{dN_X}{dt} = 4\alpha_+ X \Gamma_C m \langle A \rangle - \frac{\Gamma_X}{\hbar} N_X + P_X, \]

where \( \alpha_\pm = \alpha_1 \pm \alpha_2 \), the correlator \( A = \langle \hat{\psi}_L^\dagger \hat{\psi}_L \hat{\psi}_L \hat{\psi}_U \rangle \), and we used bosonic commutation rules to evaluate the commutators in Eq. (3). By truncating third-order correlators as products of lower order correlators [29] (e.g., \( \langle \hat{\psi}_L^\dagger \hat{\psi}_L \hat{\psi}_L \hat{\psi}_U \hat{\psi}_U \hat{\psi}_U \rangle \approx \langle \hat{\psi}_L^\dagger \hat{\psi}_L \rangle \langle \hat{\psi}_L^\dagger \hat{\psi}_L \hat{\psi}_L \rangle \rangle = A_{NL} \)), the evolution equation for \( A \) is

\[ \frac{dA}{dt} = \frac{1}{\hbar}(\beta_1 A + \beta_2) - \frac{\Gamma}{\hbar} A, \]

where \( \Gamma = (\Gamma_L + \Gamma_U)/2 + \Gamma_X \), and we defined the auxiliary functions \( \beta_1(N_{U,L,X}) \) and \( \beta_2(N_{L,U,L,X}) \), which read

\[ \beta_1 = \delta + \alpha_+ (2X^2 C^2 - 1) \]

\[ - 2(\alpha_+ + \alpha_-) X^2 N_L + C^2 N_U - (\alpha_+ - \alpha_-) N_X \]

\[ + \alpha_+ [N_L (X^4 + 2X^2 C^2) + N_U (C^4 + 2X^2 C^2)], \]

\[ \beta_2 = C^2 (N_{2 N_L U} / 2N_X + 1) - (N_L + N_U + 1) N_X^2. \]

The chosen truncation of correlators, which we use to derive equations of motion, is well justified when the parametric scattering processes are accounted for [38]. In particular, it was shown that higher order correlators play a minor role for description of a system with parametric interaction [39].

Steady-state solutions can be found by setting \( dN_{U,L,X}/dt = 0 \) and \( dA/dt = 0 \). Separating the real and imaginary parts of \( A \), gives \( \Im m \langle A \rangle = \beta_1 \Gamma_C (\beta_1^2 + \Gamma^2) \). The remaining equations can be reduced to a third-order equation for \( N_L \):

\[ N_L \Gamma_L (\beta_1^2 + \Gamma^2) + 2\alpha_+ X C \beta_2 \Gamma - \hbar P_L (\beta_1^2 + \Gamma^2) = 0, \]

where \( \beta_1 \) and \( \beta_2 \) can be rewritten in terms of \( N_L \) using

\[ N_U = \frac{N_L \Gamma_L}{\Gamma_U} + \frac{\hbar (P_U - P_L)}{\Gamma_U}, \]

\[ N_X = -2N_L \frac{\Gamma_L}{\Gamma_X} + \frac{\hbar (P_X + 2P_L)}{\Gamma_X}. \]

We note that the characteristic energy scale of the system is defined by the Rabi splitting \( \Omega_R = 2V \). In particular, the important parameter governing the system behavior is the ratio of the detuning to the Rabi energy, \( \delta/2V \).

Since \( \beta_2 \) has a cubic dependence on occupation numbers, there are in general three solutions to Eq. (9). Here, we shall be interested only in real and non-negative solutions, which for certain parameters can represent bistable behavior of the system. While an analytical solution of Eq. (9) is possible, the resulting expression is too bulky to be presented here. In what follows, we present a numerical treatment of the problem only.
III. INCOHERENT BISTABILITY WITH POLARIZED OPTICAL PUMP

First, we consider excitation with a broad LED beam, representing an incoherent source that is linearly polarized in the y direction using a polarizer. This excitation populates only the weakly coupled excitonic mode $\chi_y$ ($P_{L,U} = 0$). The system depends strongly on the mode detuning $\delta$. For positive detuning and weak pump, the condition of parametric scattering between exciton and polariton modes is not satisfied [Fig. 1(b)]. However, the increase of exciton concentration with $P_X$ leads to a blueshift of the mode, $\mu_X$, and at a threshold intensity (estimated as $P_{th} = 8\Gamma_X/2\alpha_x$) the energy difference between modes becomes equal and parametric scattering is possible [Fig. 1(c)]. This conversion of two $y$-polarized excitons into upper and lower polaritons starts to populate otherwise empty polariton modes.

The results of calculations are shown in Fig. 2 for experimentally reported parameters [28]: $V = 6$ meV, $\Gamma_c = h/\tau_c$ (where $\tau_c = 5$ ps is the cavity photon lifetime), $\tau_X = 100$ ps [42], and $S = 25 \times 10^{-3}$ m$^2$ is the pump area. The exciton-exciton scattering strength in the triplet channel was calculated as $\alpha_1 = 6E_Ba_B^3/S$ [43], where $E_B = 4.8$ meV and $a_B = 11.6$ nm are the exciton binding energy and Bohr radius in GaAs, respectively. $\alpha_x$ was shown to vary in a broad range with varying detuning [44], or can be tuned by other means. Here, we use the value $\alpha_2 = 0.4\alpha_1$ reported in Ref. [45].

For a fixed detuning, Fig. 2(a) shows that at very low pump intensity parametric scattering does not take place, and the polariton population is zero. With increasing intensity parametric scattering is turned on, increasing the polariton population at the expense of limiting the exciton population [Fig. 2(c)]. At high pump powers, the unequal blueshifts of the modes in the system switch the parametric scattering off, with the polariton population returning to small values close to zero. A similar effect was seen in Ref. [4] in resonantly excited parametric oscillators. Here, the lower density branch has a vanishingly small occupation compared to the higher branch or excitonic mode occupancies. While being typically less than unity for small detuning, it can reach an order of one in the higher detuning region. This may not be sufficient for parametric scattering as the occupation of $N_U$ may be suppressed at the same detuning. However, one can still expect energy relaxation mechanisms involving higher energy states [46], which can potentially lead to polariton lasing in the micropillar structure [47]. For a clear demonstration of bistability, we choose parameters where parametric scattering is dominant or the occupancy of $N_L$ is too small for polariton lasing.

Due to the concurrent bistability shown by the vertical gray line in Fig. 2, the system can exist in a high or low polariton density state. Note that while the excitonic mode shows conventional S-shape bistability, the shape of lower polariton bistability is different. This can be explained keeping in mind the relations between modes given by Eq. (11). Additionally, the switching between high- and low-density states can be controlled with detuning between modes, as shown in Figs. 2(b) and 2(d).

The system is further characterized by the phase diagram in the $P_X/\delta$ plane shown in Fig. 2(e). The dashed boundary marks the region in which multiple solutions are present. Taking...
parameters at the red dot in Fig. 2(e), switching between stable states can be demonstrated by numerical solution of Eqs. (4)–(7), as shown in Fig. 2(f). Assuming that the system is initially unoccupied, by switching on the continuous wave (cw) pump, the system is quickly stabilized into the higher polariton density state shown in Figs. 2(a)–2(d). A short pulse applied at $t = 550 \text{ ps}$ is able to switch the system into the lower density state. During the intervals without pulses, the system is stable, remaining in the state set by its history. We note that the appearance of the bistability window also depends on the singlet interaction constant $\alpha_2$, favoring small $\alpha_1 > \alpha_2$.

In order to give an intuitive explanation of the observed bistability effect, we can rewrite Eq. (7) in the form

$$i\hbar \frac{dA}{dt} = [-\delta + \tilde{\mu}(A) - i\Gamma] A + \tilde{P},$$

(12)

where we defined the effective energy shift $\tilde{\mu}(A) = -\beta_1 + \delta$ and effective pumping amplitude at zero frequency $\tilde{P} = -\beta_2$. We expect the bistable window to coincide with the resonance point in Eq. (12) where the amplitude of $A$ can become strongly enhanced. This corresponds to a situation where the nonlinear term (blueshift $\tilde{\mu}$) brings the system into resonance with the effective “coherent” pump of the correlator $A$. This occurs when $\beta_1 \approx 0$ and was verified numerically by tracking the sign of $\beta_1$.

Previously, we restricted our consideration to the purely zero-dimensional case, where both polaritonic and excitonic modes can be described by a three-level system. However, while the confinement of polaritonic modes is fully justified in the case of a microcavity with subwavelength grating [28], the free excitons can still possess spatial dynamics. We have also considered the exciton dispersion (Ref. [37], second section), and have shown that it does not qualitatively change the described incoherent bistability.

While we did not consider thermal/quantum noise, calculations for resonant excitation suggest only a slight narrowing of the hysteresis cycle [48]. Despite fluctuations, resonant experiments show very long memory times [17].

IV. INCOHERENT BISTABILITY WITH ELECTRICAL PUMP

Next, we consider the case of electrical injection of excitons, where both $x$- and $y$-polarized excitonic modes are pumped ($P_{L,U} \neq 0$). The dependence of the lower polariton occupation number $N_L$ on pumping strength $P_X$ is shown in Fig. 3(a) and reveals electrically pumped bistability. The key difference between electrical pumping and the optical polarized pump considered above is additional excitation of polariton modes with rates $P_L$ for lower polaritons and $P_U$ for upper polaritons, which are different in the case of nonzero detuning and due to presence of thermalization processes. Particularly, the upper polariton mode occupation is typically small, while relaxation mechanisms lead to effective incoherent pumping of lower polariton modes. This leads to an asymmetric blue shift of both exciton and polariton energies, and consequent shift of $P_{th}$ to higher pumping rates. Additionally, an incoherent bistability with electrical injection requires larger positive detuning than that with optical pumping and negative $\alpha_2$ [see Fig. 3(b)]. Large, positive values of detuning are accessible in current experiments [49,50].

The phase diagram in the $P_X/\delta$ plane is shown in Fig. 3(c). One can see that the range of the pumping strengths corresponding to the bistability window diminishes compared to the optical case, although is still fully accessible. Finally, the switching behavior of the modes, calculated similarly to the optical pump case, confirms stability of the polariton modes with the high and low population [Fig. 3(d)]. For the calculations in Fig. 3, we considered the system with a single GaAs quantum well (QW) in the cavity, with light-matter interaction constant being $V = 2 \text{ meV}$, $\alpha_2 = -0.4\alpha_1$, and effective pumping amplitudes of polariton states $P_U = 0.01C^2P_X$ and $P_L = 0.1C^2P_X$.

Finally, in order to show the possible window of parameters hosting the bistability effects, we perform calculations for systems with different nonlinear interaction constant ratios $\alpha_2/\alpha_1$ and cavity photon lifetime $\tau_C$. The results are shown in Fig. 4, where we varied parameters for the cases of optical [Figs. 4(a) and 4(c)] and electrical [Figs. 4(b) and 4(d)] pump. In the case of optical incoherent drive, the detuning was fixed to $\delta = 4 \text{ meV}$, and for electrical case, it is chosen as $\delta = 8 \text{ meV}$.

In both cases presented in plots (a) and (b), the increase of the bistable window favors large absolute value of the singlet interaction constant, comparing to the triplet one, $|\alpha_2/\alpha_1| >$
The dependence of incoherent bistability on the cavity photon lifetime is shown in Figs. 4(c) and 4(d) for optical and electrical pumps, respectively. In the case of optical pump, the behavior is rather unexpected, showing that appearance of the bistable window is favored in the low-$Q$ cavity limit (while keeping the strong coupling regime). For the electrical pumping case, the dependence is conventional, meaning that the bistable window grows with the photon lifetime.

V. CONCLUSIONS

We considered parametric scattering between $y$ polarized excitons and $x$ polarized polaritons in a subwavelength grating microcavity. We demonstrated bistable behavior in the cases of both incoherent optical polarized pumping and electrical exciton excitation. This theoretical evidence for fast-response off-resonantly driven bistability lays a foundation for hybrid electro-optic polaritonic devices.

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