

Temperature-dependent modeling of exciton-polariton bistability in semiconductor microcavities

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We develop a model to investigate the collective dynamics and coherent interaction of exciton-polaritons and longitudinal acoustic phonons in semiconductor microcavities beyond Fermi's golden rule. Our model predicts temperature-induced dephasing and an energy redshift of exciton-polaritons together with the formation of a reservoir of incoherent excitons with very high in-plane momenta owing to scattering with acoustic phonons. We validate the proposed model by reproducing the temperature-dependent changes of the polariton bistability in coherently driven semiconductor microcavities. The bistability threshold and the extent of the bistable domain are considerably modified by the blueshift induced by the incoherent reservoir. Furthermore, we describe the effect of stimulated backscattering from the reservoir into the coherent polariton state, which occurs at higher reservoir densities.

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I. INTRODUCTION

Strong coupling between cavity photons and quantum-well (QW) excitons in semiconductor planar microcavities leads to the formation of hybrid quasiparticles known as intra-cavity exciton-polaritons [1–4]. This strong exciton–photon hybridization gives rise to pronounced cubic nonlinearities originating from the Coulomb interaction between the excitonic components [5–11]. Therefore, exciton-polaritons in optical microresonators constitute a unique platform for studying nonlinear collective phenomena, including optical bistability [12–16], parametric scattering [14,17,18], superfluidity [19–21], and the formation of various structures such as quantized vortices [22–24] and solitons [25–32]. In the context of the spontaneous coherence building, exciton-polariton systems have enabled nonequilibrium analogues of Bose–Einstein condensation (BEC) [33] and low-threshold polariton lasing [34–36]. Furthermore, the exciton-polaritons are attracting increasing interest as building blocks for neuromorphic and information-processing architectures [37,38].

Beyond microcavities with conventional III–V semiconductors considered above, the nonequilibrium condensation and polariton-lasing phenomena have been demonstrated across a variety of material platforms, including organic microcavities [39], perovskite-based structures [40], and transition-metal dichalcogenide (TMD) microcavities [41–43], in some cases even at elevated or room temperatures.

Recent studies of microcavities with monolayer TMDs [42,43] have highlighted the crucial role of phonon-assisted scattering into dark exciton states and incoherent reservoirs

for polariton relaxation and thermalization. Furthermore, an appropriate excitation of the dark exciton states can be utilized to control optical bistability in a coherently driven microcavity with GaAs QWs [16].

These findings underline the general importance of phonon-mediated relaxation mechanisms across different material platforms, and motivate the development of theoretical approaches capable of capturing such processes.

A major challenge in the study of exciton-polaritons is the description of energy relaxation caused by their interaction with phonons. The resulting relaxation kinetics and thermalization processes have been investigated within the framework of Boltzmann rate equations [10,44,45] derived from Fermi's golden rule. However, while addressing scattering between states in energy–momentum space, this model neglects the coherent behavior of phonons in real space.

In a related context, polariton relaxation in lower-dimensional systems such as micropillars and photonic molecules has been simulated using semiclassical Boltzmann equations that are self-consistently coupled to a Schrödinger-type equation [46]. This equation can be solved at any given time to obtain the eigenstates in the presence of potentials created by the exciton reservoir and by polaritons populating these states. However, this approach is not well suited for simulating nonlinear polariton dynamics within an arbitrary potential landscape.

Another approach is based on a stochastic Gross–Pitaevskii-type equation that also accounts for phonon-mediated relaxation [47]. It can formally be solved, by analogy with the Markov approximation within the Langevin framework, under the assumption that acoustic phonons possess randomly varying phases. This approach has shown good agreement with experimental measurements of energy relaxation in propagating condensates within semiconductor microwires [48].

In this paper, we introduce a theoretical framework for simulating the collective nonlinear dynamics of exciton-polaritons interacting with acoustic phonons beyond the

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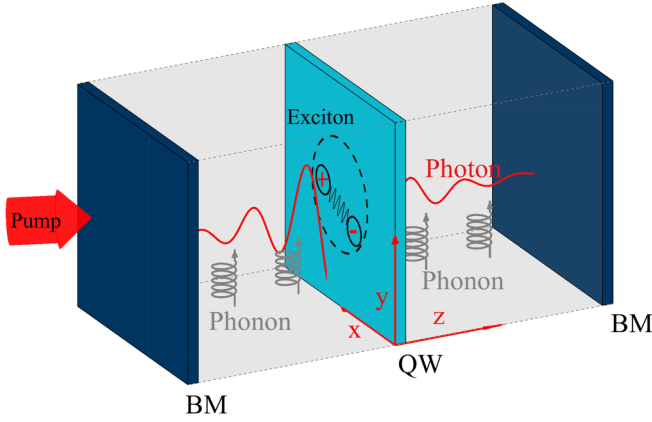


FIG. 1. Schematic illustration of the GaAs-based planar microcavity containing a single InGaAs semiconductor quantum well (QW) embedded at the antinode of the optical field between two distributed Bragg mirrors (BM). The intracavity photons are strongly coupled to the 2D excitons in the QW. Longitudinal acoustic phonons interact with the excitons via the deformation potential.

applicability of Fermi's golden rule. Starting from the well-established approach [19,20] describing the evolution of QW excitons coupled to classical photons in semiconductor microcavities, we extend the model to include a wave equation describing the evolution of acoustic phonons. Primarily, the proposed model is intended for the simulation of exciton-polariton dynamics in structures based on III-V inorganic semiconductor materials, such as GaAs, hosting Wannier-Mott excitons. Their interaction with acoustic phonons can be described via electron and hole deformation potentials [44,53]. The excitons remain well defined at the sample temperatures below $T \approx 116$ K, where the typical energy scale of thermal fluctuations ($k_B T$) is below the binding energy of excitons, which is roughly 10 meV in a typical QW [10]. This condition sets an approximate upper temperature limit for the quantitative applicability of the model. Nevertheless, for a qualitative understanding of the underlying physics behind the nonlinear dynamics of exciton-polaritons, we will also consider higher temperatures.

In particular, we examine the influence of lattice temperature on optical bistability within a coherently pumped semiconductor microcavity. We note that the proposed model, which incorporates microscopic scattering processes, is suitable for describing polaritons within an arbitrary two-dimensional potential landscape (for a review, see [49]) and is also applicable under incoherent pumping conditions.

This paper is organized as follows. In Sec. II, we present our theoretical approach for a semiconductor microcavity operating in the strong-coupling regime (Fig. 1) and derive a set of coupled differential equations to simulate the nonlinear dynamics of exciton-polaritons. In Sec. III, we analyze the interaction between exciton-polaritons and longitudinal acoustic phonons and discuss the formation of a reservoir of incoherent excitons with high momenta. Based on our model, we also determine the temperature-dependent dephasing times of exciton-polaritons. In Sec. IV, we examine the optical bistability response of a microcavity coherently pumped near the lower polariton branch. We develop an analytical model

to investigate the bistability domain for different lattice temperatures and compare our analytical predictions with direct numerical simulations, finding very good agreement.

II. THEORETICAL MODEL

In this section, we discuss the physical processes and the corresponding theoretical approaches required to simulate the nonlinear dynamics of exciton-polaritons in the presence of acoustic phonons. The Hamiltonian of the system can be conveniently divided into two parts,

$$H_0 = \sum_{\mathbf{k}} H_{\text{ph-ex},\mathbf{k}} + H_{\text{phonon}}, \quad (2.1)$$

where the wave vector $\mathbf{k} = (k_x, k_y)$ is defined in the plane of the microcavity (see Fig. 1). In Hamiltonian (2.1), the photon-exciton term $H_{\text{ph-ex},\mathbf{k}}$ describes the intracavity photons that are strongly coupled to the QW excitons, while H_{phonon} accounts for the interaction between excitons and acoustic phonons. In the following, we describe both subsystems in detail and derive the corresponding master equations.

A. Photon-exciton subsystem

The dynamics of intracavity photons strongly coupled to QW excitons in a semiconductor microcavity (see Fig. 1) are described in momentum space by the Hamiltonian [2,3,17,19,20],

$$\begin{aligned} H_{\text{ph-ex},\mathbf{k}} = & \left(\hbar\omega_c + \frac{\hbar^2|\mathbf{k}|^2}{2m_c} \right) a_{\mathbf{k}}^\dagger a_{\mathbf{k}} + \hbar\Omega_R (a_{\mathbf{k}}^\dagger b_{\mathbf{k}} + b_{\mathbf{k}}^\dagger a_{\mathbf{k}}) \\ & + \left(\hbar\omega_e + \frac{\hbar^2|\mathbf{k}|^2}{2m_e} \right) b_{\mathbf{k}}^\dagger b_{\mathbf{k}} + \hbar g \sum_{\mathbf{k}',\mathbf{q}} \frac{1}{A} b_{\mathbf{k}+\mathbf{q}}^\dagger b_{\mathbf{k}'-\mathbf{q}}^\dagger b_{\mathbf{k}} b_{\mathbf{k}'} \\ & + i\hbar(E_p(\mathbf{k})e^{-i\omega_p t} a_{\mathbf{k}}^\dagger - E_p^*(\mathbf{k})e^{i\omega_p t} a_{\mathbf{k}}). \end{aligned} \quad (2.2)$$

Here, $a_{\mathbf{k}}^\dagger$ ($a_{\mathbf{k}}$) and $b_{\mathbf{k}}^\dagger$ ($b_{\mathbf{k}}$) denote the creation (annihilation) operators for cavity photons and excitons with in-plane wave vector $\mathbf{k}$. The prefactors of the photon and exciton number operators, $a_{\mathbf{k}}^\dagger a_{\mathbf{k}}$ and $b_{\mathbf{k}}^\dagger b_{\mathbf{k}}$, represent their respective dispersions, and Ω_R is the half Rabi frequency describing exciton-photon coupling. m_c and m_e are the effective masses of cavity photons and excitons, respectively, while ω_c and ω_e denote their resonance frequencies.

The last term in Eq. (2.2) represents an external driving of the cavity by a coherent field of frequency ω_p and amplitude $E_p(\mathbf{r})$, with $E_p(\mathbf{k})$ denoting its Fourier transform. This term injects photons into the cavity at the pump frequency. Simulation parameters correspond to a microcavity containing a single 10-nm InGaAs QW with zero photon-exciton detuning ($\omega_c = \omega_e$) and $\hbar\Omega_R = 3.291$ meV [20]. Typical effective masses are $m_c = 2 \times 10^{-5} m_0$ and $m_e = 0.5 m_0$, where m_0 is the free electron mass ($m_0 = 5678$ meV ps² μm⁻²). In the low-density limit and neglecting the external pumping term ($|E_p| = 0$), the Hamiltonian (2.2) can be diagonalized in a basis of exciton-polariton quasiparticles. The resulting dispersion relation, i.e., the dependence of the polariton energy

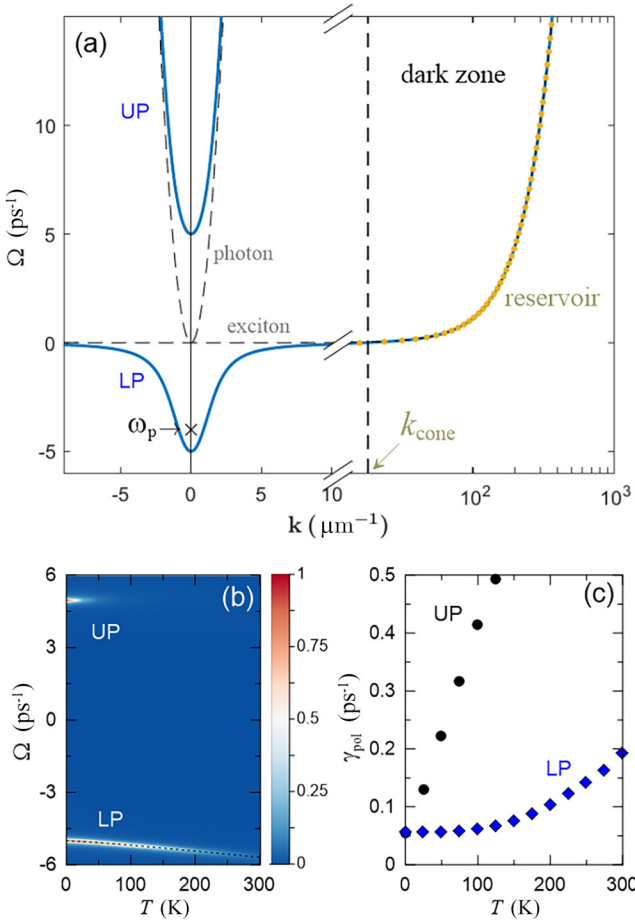


FIG. 2. (a) Polariton dispersion of the semiconductor microcavity calculated from Eq. (2.3) as a function of the in-plane wave vector $\mathbf{k} = (k_x, k_y)$. LP and UP denote the lower- and upper-polariton branches, respectively. Intracavity photons couple to excitons within the light cone ($|\mathbf{k}| < k_{\text{cone}}$, indicated by the vertical dashed line), whereas in the “dark zone” ($|\mathbf{k}| > k_{\text{cone}}$) excitons form an incoherent reservoir. (b) Normalized spectral response of photonic part and (c) damping constants of the LP and UP branches, calculated within the theoretical framework [see Eqs. (2.6)–(2.8)] and plotted as functions of sample temperature T . The LP energy shift [dashed line in (b)] is fitted by the redshift of the exciton resonance given by the empirical function $\omega_e(T) = -0.044 \text{ ps}^{-1} \text{ K}^{-1} T / (1 + 0.0004 \text{ K}^{-1} T)$, with T in Kelvin. Parameters: $\hbar\Omega_R = 3.291 \text{ meV}$, $\omega_c = \omega_e = 0$, $m_c = 2 \times 10^{-5} m_0$, $m_e = 0.5 m_0$, where m_0 is the free electron mass.

$E(\mathbf{k}) \equiv \hbar\Omega(\mathbf{k})$ on the in-plane wave vector $\mathbf{k}$, reads

$$\hbar\Omega_{\mp}(\mathbf{k}) = \frac{\hbar}{2}(\omega_c + \omega_e) + \frac{\hbar^2(m_e + m_c)}{4m_cm_e}|\mathbf{k}|^2 \mp \sqrt{\left[\frac{\hbar}{2}(\omega_c - \omega_e) + \frac{\hbar^2(m_e - m_c)}{4m_cm_e}|\mathbf{k}|^2\right]^2 + \hbar^2\Omega_R^2}. \quad (2.3)$$

The + and – signs correspond to the upper- (UP) and lower-polariton (LP) branches, respectively [see Fig. 2(a)].

Since the photon effective mass is much smaller than the exciton effective mass, the LP branch becomes almost purely

excitonic for in-plane momenta $|\mathbf{k}| > k_{\text{cone}}$, i.e., beyond the maximum photon wave vector determined by the light cone, $k_{\text{cone}} = n_{\text{eff}}\omega_e/c$, where c is the speed of light in vacuum and n_{eff} the effective refractive index of the microcavity material. Although photons do not couple directly to high-momentum excitons, these excitons play a crucial role in forming an incoherent reservoir through phonon-mediated scattering. Here, we retain the parabolic approximation for the excitonic dispersion, even though the exciton effective mass is typically four to five orders of magnitude larger than the photon effective mass.

The nonlinear term in the second line of Hamiltonian (2.2), where A represents the macroscopic in-plane quantization area, accounts for Coulomb scattering between excitons within the Born approximation [7,10,11]. It gives rise to an effective repulsive interaction characterized by a positive coupling constant ($g > 0$), which dominates for excitons with parallel spins [11,50–52]. For the simulations, we use $\hbar g = 0.0033 \text{ meV } \mu\text{m}^2$ [11,51].

B. Phonon subsystem

Within this model, we focus exclusively on longitudinal acoustic phonons, which interact with excitons through the deformation potential. This potential accounts for the local change in the excitonic energy induced by lattice compression owing to phonons [44,45,53].

We are particularly interested in thermal effects, for which the exciton–phonon interaction is dominated by low-energy acoustic phonons whose wavelengths are much larger than both the QW width and the exciton Bohr radius. Consequently, the phonons considered here cannot resolve the individual electron and hole constituents of the exciton; instead, they interact with the exciton as a single quasiparticle. Under these conditions, the interaction Hamiltonian of two-dimensional (2D) QW excitons with three-dimensional (3D) acoustic phonons can be simplified [44,45,53] to the form

$$H_{\text{phonon}} = \sum_{q_z, \mathbf{q}} \hbar v_{\text{ph}} (|\mathbf{q}|^2 + q_z^2)^{1/2} u_{\mathbf{q}, q_z}^\dagger u_{\mathbf{q}, q_z} + \sum_{q_z, \mathbf{q}, \mathbf{k}} G(\mathbf{q}, q_z) (u_{\mathbf{q}, q_z} - u_{-\mathbf{q}, -q_z}^\dagger) b_{\mathbf{k}}^\dagger b_{\mathbf{k}-\mathbf{q}}. \quad (2.4)$$

Here, $u_{\mathbf{q}, q_z}^\dagger$ ($u_{\mathbf{q}, q_z}$) denotes the creation (annihilation) operator of a 3D phonon with in-plane wave vector $\mathbf{q}$ and out-of-plane component q_z , where the z axis corresponds to the growth direction. Since the z component of the phonon momentum is not conserved, the summation over q_z is required in Eq. (2.4).

All parameters characterizing the exciton–phonon interaction are contained in the coupling coefficient

$$G(\mathbf{q}, q_z) = -i\alpha_0 \sqrt{\frac{\hbar(|\mathbf{q}|^2 + q_z^2)^{1/2}}{2\rho v_{\text{ph}} V}}. \quad (2.5)$$

Here, $\rho = 5.32 \text{ g cm}^{-3}$ and $v_{\text{ph}} = 4.73 \times 10^{-3} \text{ } \mu\text{m}/\text{ps}$ denote the mass density and the longitudinal sound velocity of the semiconductor, respectively, while V is the quantization volume. The effective deformation potential for InGaAs QW excitons is $\alpha_0 = 10.4 \text{ eV}$ (see [54,55] and Appendix A).

Because the overlap integral between the exciton envelope function and the phonon wave function vanishes for very large momenta [44,45], cutoffs are introduced in Eq. (2.5) for $|\mathbf{q}| > 1/a_B$ and $q_z > 2\pi/d$, where a_B denotes the Bohr radius of an exciton confined in the QW of width d . For our model, we assume an exciton radius $a_B = 10$ nm, which is typical for a GaAs QW with a width $d \approx a_B$ [45].

C. Mean-field master equations

The subsequent analysis is conducted in real space; therefore, we perform a transformation from momentum space to a finite-dimensional 2D system of quantized area A , given by

$$a_c(\mathbf{r}) = \frac{1}{\sqrt{A}} \sum_{\mathbf{k}} a_{\mathbf{k}} e^{i\mathbf{k}\cdot\mathbf{r}}, \quad b_e(\mathbf{r}) = \frac{1}{\sqrt{A}} \sum_{\mathbf{k}} b_{\mathbf{k}} e^{i\mathbf{k}\cdot\mathbf{r}}.$$

Using the Heisenberg equation of motion based on the Hamiltonian (2.2) and (2.4), and performing the inverse Fourier transformation, we derive the evolution equations for operators describing intracavity photons strongly coupled to QW excitons (see [2,3,20] and Appendix A).

Following a mean-field approximation, we introduce the classical cavity field ψ_c and the exciton polarization ψ_e as expectation values $\psi_c = \langle a_c(\mathbf{r}) \rangle$ and $\psi_e = \langle b_e(\mathbf{r}) \rangle$. The evolution equations of these complex fields are given by

$$\begin{aligned} \partial_t \psi_c = & -i \left(\delta_c - i\gamma_c - \frac{\hbar}{2m_c} \Delta_{\parallel} \right) \psi_c \\ & + i\Omega_R \psi_e + E_p, \end{aligned} \quad (2.6)$$

$$\begin{aligned} \partial_t \psi_e = & -i \left(\delta_e - i\gamma_e + g|\psi_e|^2 - \frac{\hbar}{2m_e} \Delta_{\parallel} \right) \psi_e \\ & + i\Omega_R \psi_c + \frac{i\alpha_0}{\hbar} P_{\text{ph}}(\mathbf{r}, z=0) \psi_e. \end{aligned} \quad (2.7)$$

All fields are driven by the incident field E_p , now acting as source terms in the equations. $\Delta_{\parallel} = \partial_{x,y}^2$ is the Laplace operator in the plane of the cavity, and $\delta_{e,c} = \omega_{e,c} - \omega_p$ denote the detunings of the pump frequency ω_p from the resonances of excitons (ω_e) and cavity photons (ω_c), respectively. It is noteworthy that the model takes into account excitons with very high momenta, far beyond the range allowed for photons. In addition, it accounts for unavoidable losses by including damping constants associated with the lifetimes of intracavity photons γ_c and QW excitons γ_e , with typical values of $\hbar\gamma_c = 0.0658$ meV and $\hbar\gamma_e = 0.0066$ meV [13].

The space and time-dependent lattice compression P_{ph} caused by longitudinal acoustic phonons changes the exciton energy via the deformation potential α_0 , thus inducing a time-varying energy landscape as represented in Eq. (2.7). In contrast to previous approaches [47], we explicitly consider the dynamics of these longitudinal acoustic phonons by solving the wave equation for the lattice compression P_{ph} (for the derivation of this equation from the Hamiltonian, see Appendix A1)

$$\partial_t^2 P_{\text{ph}} = v_{\text{ph}}^2 (\Delta_{\parallel} + \partial_z^2) P_{\text{ph}} - \frac{\alpha_0}{\rho d} \Delta_{\parallel} |\psi_e|^2 f(z), \quad (2.8)$$

with the mass density of the material ρ and the thickness d of the QW. The QW localization function is defined as $f(z) \equiv 1$ within the QW layer $-d/2 < z < d/2$ and $f(z) \equiv 0$

otherwise. Thus, each variation of the exciton density drives phonon waves by inducing a local lattice deformation, which quickly spreads with the velocity of sound v_{ph} into the surrounding material. The escape of the acoustic phonons from the QW into the volume of the microcavity is compensated by the phonons entering from the environment, which is included in the model by taking into account the third dimension z orthogonal to the QW. In all simulations based on Eq. (2.8), we start with a phonon population in thermodynamic equilibrium with the lattice temperature T (see Appendix B for further details).

To simplify the numerical treatment (see Appendix C for numerical details), but still obtain a qualitative understanding of the underlying physics, we limit our analysis to a small-sized system with dimensions of 800 nm in both transverse directions (x and y). This choice of geometry allows us to focus on the influence of phonons while suppressing previously investigated optically driven modulational instabilities of the exciton-polariton system [19,26,56]. The physical rationale for using a reduced in-plane simulation domain is provided in Appendix C. The finite speed of sound permits us to restrict the z direction orthogonal to the QW to 1000 nm while still avoiding any influence of the boundaries within our simulation time.

In an initial test of the proposed model [Eqs. (2.6)–(2.8)], we calculated the spectral response as a function of lattice temperature [see Fig. 2(b)]. For this analysis, we assumed excitation by a sequence of short input pulses (~ 1 ps) and determined the spectra of the emitted optical field. We found the energy of the LP branch to experience a gradual redshift with increasing temperature. A plausible explanation for this behavior is that excitons with lower momenta $\mathbf{k}$ tend to regions of lower energy within the phonon-created potential. Another interesting aspect is a strong imbalance between phonon-mediated damping of the UP and LP branches [see Fig. 2(c)]. This enhanced scattering of the UP branch is explained by the large number of available excitonic states within the same energy interval.

III. PHONON-MEDIATED DECOHERENCE

Conventional models for simulating exciton-polariton dynamics [12–14,17–21,26–28,56–58] typically consider only the coherent excitons with relatively small in-plane momenta $\mathbf{k}$. Therefore, very small excitonic damping constants are usually included phenomenologically to describe dephasing induced by exciton scattering. In our model [Eqs. (2.6)–(2.8)], phonon-induced dephasing of excitons is naturally incorporated via scattering from the deformation potential.

In this section, we estimate the phonon-mediated decoherence time of exciton-polaritons using the theoretical approach described above [see Eqs. (2.6)–(2.8)]. As shown below, this temperature-dependent decoherence mechanism is associated with the formation of a reservoir of incoherent excitons and plays a crucial role in the occurrence of exciton-polariton bistability.

We assume that the microcavity is coherently pumped by an external optical beam at a frequency near the bottom of the LP branch [see Fig. 2(a)]. Thus, coherently created exciton-polaritons scatter at the deformation potential caused

by acoustic phonons [see the term with α_0 in Eq. (2.7)] and populate excitonic states with high in-plane momenta $\mathbf{k}$, generating a reservoir of incoherent excitons. Excitons with very high momenta ($|\mathbf{k}| > k_{\text{cone}}$) cannot couple to photons [see “dark zone” in Fig. 2(a)] and lose any strict phase relation owing to further scattering on the acoustic phonons. This decoherence effect becomes stronger with increasing temperature T , resulting in a shorter lifetime of the coherent exciton-polariton state at the bottom of the band [see Figs. 2(b) and 2(c)].

To illustrate the mechanism of phonon-driven dephasing, we first split the exciton population into a coherent component $\psi_{\text{coh}}(\mathbf{r})$, which can interact with light, and an incoherent component $\psi_{\text{inc}}(\mathbf{r})$, whose momentum distribution lies outside the light cone, as follows:

$$\psi_{\text{coh}}(\mathbf{r}) = \int \int_{|\mathbf{k}| < k_{\text{cone}}} dk_x dk_y \psi_e(\mathbf{k}) e^{i\mathbf{k} \cdot \mathbf{r}},$$

$$\psi_{\text{inc}}(\mathbf{r}) = \int \int_{|\mathbf{k}| \geq k_{\text{cone}}} dk_x dk_y \psi_e(\mathbf{k}) e^{i\mathbf{k} \cdot \mathbf{r}},$$

where the entire population in momentum space is given by

$$\psi_e(\mathbf{k}) = \frac{1}{(2\pi)^2} \int \int_A dx dy \psi_e(\mathbf{r}) e^{-i\mathbf{k} \cdot \mathbf{r}}.$$

To quantify the strength of phonon-induced dephasing, we simplify Eq. (2.7) by concentrating on coherent excitons and representing the influence of phonons by a temperature-dependent dephasing rate Γ_e as

$$\begin{aligned} \partial_t \psi_{\text{coh}}(\mathbf{r}) = & -i[\delta_e - i(\gamma_e + \Gamma_e - RN_{\text{inc}}(\mathbf{r}))]\psi_{\text{coh}}(\mathbf{r}) \\ & + i\Omega_R \psi_c(\mathbf{r}) - ig(|\psi_{\text{coh}}(\mathbf{r})|^2 + N_{\text{inc}}(\mathbf{r}))\psi_{\text{coh}}(\mathbf{r}). \end{aligned} \quad (3.1)$$

For a vanishing density of incoherent excitons $N_{\text{inc}}(\mathbf{r}) \equiv |\psi_{\text{inc}}(\mathbf{r})|^2 = 0$, Eq. (3.1) coincides with the widely used model for exciton-polariton dynamics [12–14, 17–21, 26–28, 56–58], where the dephasing rate Γ_e is introduced phenomenologically.

A nonvanishing reservoir density N_{inc} gives rise to frequency blueshift and the gain of the coherent state in Eq. (3.1). This gain is induced by phonon-mediated backscattering, which we phenomenologically introduce via a term $RN_{\text{inc}}(\mathbf{r})$. Such stimulated emission into the condensate mode depletes the incoherent reservoir and is well known in Gross–Pitaevskii models [20, 59].

Neglecting the diffusion of excitons, the reservoir of incoherent excitons is described by the rate equation,

$$\partial_t N_{\text{inc}}(\mathbf{r}) = -2(\gamma_e + R|\psi_{\text{coh}}(\mathbf{r})|^2)N_{\text{inc}}(\mathbf{r}) + 2\Gamma_e|\psi_{\text{coh}}(\mathbf{r})|^2. \quad (3.2)$$

Since the rate equation (3.2) describes the evolution of a real-valued quantity, namely, the density of excitons N_{inc} , the frequency blueshift and, in general, all phase-shift terms with prefactor g in Eq. (3.1) must vanish. Therefore, the influence of the exciton-exciton interaction on the reservoir dynamics via parametric scattering [2, 14, 20] is phenomenologically included in the model [Eqs. (3.1) and (3.2)] via a gain term with coefficient R .

By comparing the full and simplified models, we now estimate the dephasing constant Γ_e and the scattering rate R for different lattice temperatures T . According to Eq. (3.2), the decoherence rate Γ_e is given in the steady state by

$$\begin{aligned} \Gamma_e(T) &= (\gamma_e + R(T)|\psi_{\text{coh}}(\mathbf{r})|^2) \frac{N_{\text{inc}}}{|\psi_{\text{coh}}|^2} \\ &\equiv (\gamma_e + R(T)|\psi_{\text{coh}}(\mathbf{r})|^2)\beta(T), \end{aligned} \quad (3.3)$$

and is directly related to the density ratio of incoherent to coherent excitons β .

In order to determine this ratio, we perform a series of numerical simulations within the original model [Eqs. (2.6)–(2.8)] for a homogeneous, time-independent coherent pump of low intensity. As soon as the system has reached a steady state, we calculate $\beta(T)$ by separating the respective components in Fourier space. Now, the dephasing rate $\Gamma_e(T)$ is estimated based on Eq. (3.3) in the limit of very small densities of coherent excitons ($|\psi_{\text{coh}}(\mathbf{r})|^2 \approx 0$) [see Fig. 3(a)]. The resulting dependence is well described by a smooth, monotonically growing analytical function [see the solid line in Fig. 3(a)], which will be used in our later analysis of bistability in Sec. IV.

In order to determine the scattering rate R , we now analyze the stationary state of the full model for increasing pumping and consequently higher densities of coherent excitons. As displayed in Fig. 3(b), the system is dominated by the incoherent reservoir for low excitations, but the ratio β decreases monotonically with increasing coherent-exciton density. Such behavior is caused by stimulated backscattering from the reservoir [see the term with the backscattering rate R in Eq. (3.2)]. Having previously determined the dephasing rate Γ_e , the scattering coefficient R is now determined based on Eq. (3.3). By repeating this procedure for different temperatures, we find the backscattering rate R to be well approximated by a linear function of T [see the inset in Fig. 3(b)].

As the dephasing rate Γ_e was determined solely on the basis of the stationary solution of the reservoir rate equation (3.2), we now test our findings by simulating the dynamics of the full system. To this end, we model the decay of low-density exciton-polaritons at different temperatures. Indeed, exponential decay curves based on the previously determined decoherence rates $\Gamma_e(T)$ match the numerical findings very well [compare the black and green lines in Fig. 3(c)].

IV. BISTABILITY OF EXCITON-POLARITONS

In this section, we investigate the bistable response of exciton-polaritons by launching a coherent optical beam with frequency ω_p slightly above the resonance frequency of the LP branch [see Fig. 2(a)]. For stronger pumping, the intracavity exciton-polaritons experience a frequency blueshift owing to the nonlinear interaction between excitonic components. Once this frequency blueshift exceeds a certain threshold, it can compensate for the frequency detuning of the pump, resulting in an abrupt increase in the intracavity field. This mechanism is responsible for the formation of bistability and hysteresis, as demonstrated for two different detunings in Fig. 4(a).

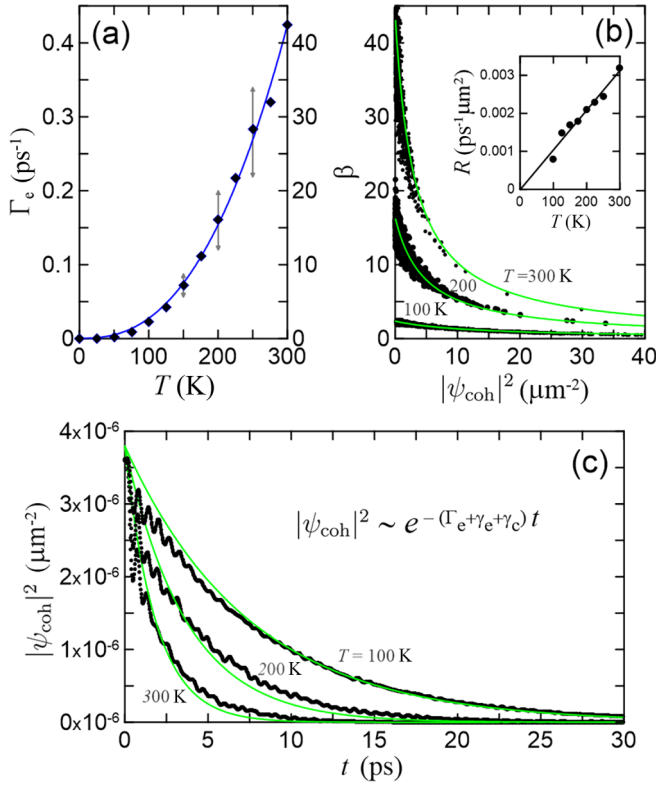


FIG. 3. (a) Effective dephasing rate Γ_e near the LP branch estimated by Eq. (3.3) as a function of the lattice temperature T . The vertical lines represent intervals of statistical tolerance. The solid line shows a fit to the function $f(T) = \gamma_e \times 2.72 \times 10^{-5} T^{2.5}$ with T given in Kelvin and γ_e in ps^{-1} . (b) The ratio of incoherent to coherent excitons β calculated for varying densities of coherent excitons (filled circles). The solid lines represent fits based on Eq. (3.3) including a temperature-dependent fitting parameter R . The inset shows the extracted coefficient R vs T , which is approximated by the function $f(T) = 1.045 \times 10^{-5} \text{ ps}^{-1} \mu\text{m}^2 \text{ K}^{-1} T$. (c) Numerically observed decay of the coherently generated exciton-polaritons calculated within the full model [Eqs. (2.6)–(2.8)] for different temperatures (100 K, 200 K, and 300 K). The solid lines represent exponential fits with damping rates $\Gamma = \Gamma_e + \gamma_e + \gamma_c$ with Γ_e taken from (a). Parameters are $\delta_{c,e} = 4.75 \text{ ps}^{-1}$, $\gamma_c = 0.1 \text{ ps}^{-1}$, $\gamma_e = 0.01 \text{ ps}^{-1}$, $\Omega_R = 5 \text{ ps}^{-1}$.

A. Semi-analytic approach neglecting the reservoir

We use Eq. (3.1), which includes the previously determined dephasing rate Γ_e and assume a negligible population of the incoherent reservoir $N_{\text{inc}} \approx 0$. This assumption is justified for relatively fast switch-on timescales shorter than the typical build-up time of the incoherent reservoir described by Eq. (3.2).

We then determine the steady-state solution of Eqs. (2.6) and (3.1) by setting the time derivatives to zero ($\partial_t = 0$). For a spatially uniform coherent pump E_p , the system exhibits translational invariance in the QW plane, allowing for homogeneous (steady) solutions (HSs), thus eliminating the in-plane Laplacian $\nabla_{\parallel}^2$ in Eq. (2.6). The corresponding densities are $n_c = |\psi_c|^2$ and $n_e = |\psi_{\text{coh}}|^2$. We note that while the semi-analytic treatment assumes strict translational invariance, the full numerical simulations include a thermal phonon

bath for $T > 0$, which modulates the exciton and photon densities in space even under spatially uniform coherent driving. However, the photon modulation is weak, and the corresponding solutions are therefore still referred to as homogeneous, consistent with previous theoretical treatments where phonon effects enter via an effective damping only.

It was shown in [56–58] that $n_c(E_p)$ and $n_e(E_p)$ are multi-valued functions of the pump field amplitude E_p provided that $f(\delta_c, \delta_e) > 0$ holds for

$$f(\delta_c, \delta_e) \equiv \delta_c \Omega_R^2 - \delta_e (\delta_c^2 + \gamma_c^2) - \sqrt{3} \gamma_c \Omega_R^2 - \sqrt{3} (\gamma_e + \Gamma_e) (\delta_c^2 + \gamma_c^2). \quad (4.1)$$

Figure 4(a) shows hysteresis loops of HSs for two different frequency detunings. Within the bistability range, two states are stable whereas the middle branch (depicted by dashed lines) is unstable (for details, see [56–58]). Two turning points (TPs) define the boundaries of the bistability region. They fulfill the condition $\frac{\partial |E_p|^2}{\partial n_c} = 0$, which yields

$$n_c^{\pm} = \frac{n_e^{\pm}}{\Omega_R^2} [(\delta_e + g n_e^{\pm})^2 + (\gamma_e + \Gamma_e)^2], \quad (4.2)$$

with

$$n_e^{\pm} = \frac{1}{3g} \left[2 \left(\frac{\delta_c \Omega_R^2}{\delta_c^2 + \gamma_c^2} - \delta_e \right) \pm \sqrt{\left(\frac{\delta_c \Omega_R^2}{\delta_c^2 + \gamma_c^2} - \delta_e \right)^2 - 3 \left(\gamma_e + \Gamma_e + \frac{\gamma_c \Omega_R^2}{\delta_c^2 + \gamma_c^2} \right)^2} \right].$$

In contrast to the previous theoretical studies of exciton-polariton bistability [56–58], our results, based on Eqs. (4.1) and (4.2), account for temperature dependence via the dephasing rate $\Gamma_e(T)$. Using the estimates determined in the previous section [see Fig. 3(a)] and Eq. (4.2), we plot the turning points of the bistability region as a function of the phonon temperature [see Fig. 4(b)]. The extent of the bistable response of the system is influenced by the phonon-mediated dephasing of the exciton-polaritons $\Gamma_e(T)$. With increasing temperature, its effect becomes stronger, resulting in a shorter lifetime of the coherent exciton-polariton state. Thus, the bistable range shrinks and eventually disappears at higher temperatures above a certain threshold ($T > T_{\text{th}}$). Notably, this cutoff temperature rises with the frequency detuning of the pump field (ω_p) from the resonance frequency of the LP branch. Higher lattice temperatures also require stronger pumping for bistability [see Fig. 4(c)].

The buildup of the incoherent reservoir substantially affects the bistability on timescales longer than the lifetime of coherent excitons ($\gamma_e^{-1} \gtrsim \Gamma_e^{-1}$). The additional blueshift induced by the incoherent reservoir shifts the resonance, thereby reducing the bistability threshold at higher temperatures. On the other hand, the recovery of coherent excitons driven by stimulated scattering from the reservoir [see the term with R in Eq. (3.1)] can support the formation of bistability, which we will describe in the following subsection.

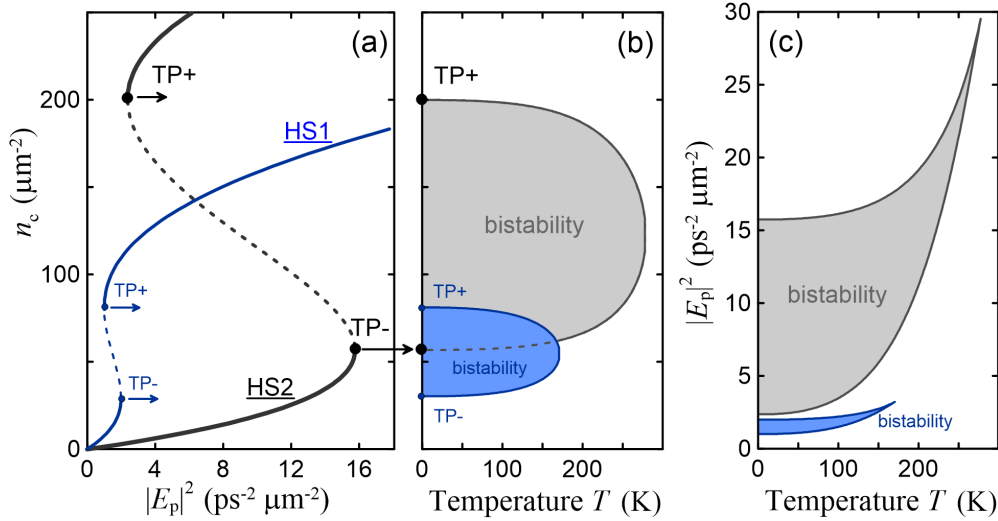


FIG. 4. (a) The intracavity photon density $n_c = |\psi_c|^2$ of the homogeneous solution (HS) plotted versus the pump intensity $|E_p|^2$ at $T = 0$ K and two different detunings: $\delta_{c,e} = 4.81 \text{ ps}^{-1}$ (the curve “HS1”) and $\delta_{c,e} = 4.6 \text{ ps}^{-1}$ (the curve “HS2”). The solid/dashed lines indicate stable/unstable HSs. The bistability domain is bounded by two turning points (TPs) where stable and unstable solutions merge. (b) The intracavity fields of TPs calculated from Eq. (4.2) vs lattice temperature T for the frequency detunings provided in (a). (c) The pump intensities of the TPs for the frequency detunings given in (a) and (b). The shaded regions in (b) and (c) indicate the temperature-dependent bistable domain. Other parameters are given in the caption to Fig. 3.

B. Effect of incoherent reservoir: Direct numerical simulations

Finally, we perform numerical simulations of the bistability domain using the full model [Eqs. (2.6)–(2.8)] to conduct a detailed study of the influence of the incoherent reservoir. During each simulation, we adiabatically sweep the pumping amplitude E_p within the estimated bistability range slowly enough to ensure steady-state conditions for the incoherent reservoir. In order to ensure steady-state conditions, we require the sweeping time over the entire bistability range to be much larger than the longest time constant of the model, here given by the exciton lifetime ($t_{\text{sweep}} \gg \gamma_e^{-1}$). Here, we used $t_{\text{sweep}} = 40 \text{ ns}$ for all calculations and also verified that further increasing the sweeping time leads to negligible changes in the hysteresis loop, confirming temporal convergence.

At the turning points of bistability (TP), the intracavity field undergoes a sudden jump either up or down, depending on the sweeping direction. The respective pump and intracavity field intensities are plotted versus the phonon temperature [marked by rhombuses and circles in Figs. 5(a) and 5(b)].

In contrast to the case without the reservoir discussed above, the bistable range of intracavity intensities does not monotonically shrink with temperature [compare Fig. 4(b) and Fig. 5(a)]. Figures 5(c) and 5(d) illustrate two exemplary hysteresis loops for $T = 25 \text{ K}$ and $T = 100 \text{ K}$, for which the bistable domain widens with increasing temperature. This behavior is directly related to the additional frequency blueshift induced by the incoherent reservoir, together with the frequency redshift—shown by the dashed line in Fig. 2(b)—associated with the phonon-created potential. Nevertheless, bistability finally disappears for higher temperatures [see Fig. 5(b)].

We now estimate the extent of the bistability domain in the presence of the incoherent reservoir. We adapt the semi-analytic approach [Eq. (4.2)] and compare its results with direct numerical simulations. To account for the influence of the incoherent reservoir, we express its density as $N_{\text{inc}} \equiv \beta n_{\text{coh}}$ and substitute it in Eq. (3.1). As discussed in Sec. III, the steady-state ratio of incoherent and coherent excitons β is given by Eq. (3.3) [see also Fig. 3(b)]. It can be shown that the inclusion of the reservoir can be addressed by substituting the parameters: $g \rightarrow (1 + \beta(T, n_{\text{coh}}))g$ and $\Gamma_e \rightarrow \gamma_e \beta(T, n_{\text{coh}})$ in the semi-analytic approach [Eq. (4.2)]. In addition, we also incorporate a temperature-dependent exciton energy $\delta_e(T)$ to account for the phonon-induced frequency redshift of the LP branch [see the dashed line along the LP resonance in Fig. 2(b)]. Including the reservoir in the semi-analytic approach [Eq. (4.2)] causes the TPs to depend not only on temperature but also on the density of coherent excitons n_{coh} mediated by the factor $\beta(T, n_{\text{coh}})$. To address this, we estimated the lower turning point (TP[−]) in the low-density limit $n_{\text{coh}} = 0$, while the upper turning point (TP⁺) was calculated for a density $n_{\text{coh}} = 115 \mu\text{m}^{-2}$ (the density at the turning point in the low-temperature limit $T \rightarrow 0$). The solid lines in Figs. 5(a) and 5(b) show the analytically estimated turning points (TP[±]) versus lattice temperature (T). Both the adapted semi-analytic approach [Eq. (4.2)] and the direct numerical simulations [Eqs. (2.6)–(2.8)] agree very well.

V. DISCUSSION AND CONCLUSIONS

We developed a dynamical theory that accounts for the coherent interaction between longitudinal acoustic phonons and strongly coupled microcavity exciton–polaritons. It allows us to predict and quantify the influence of a thermal bath on the

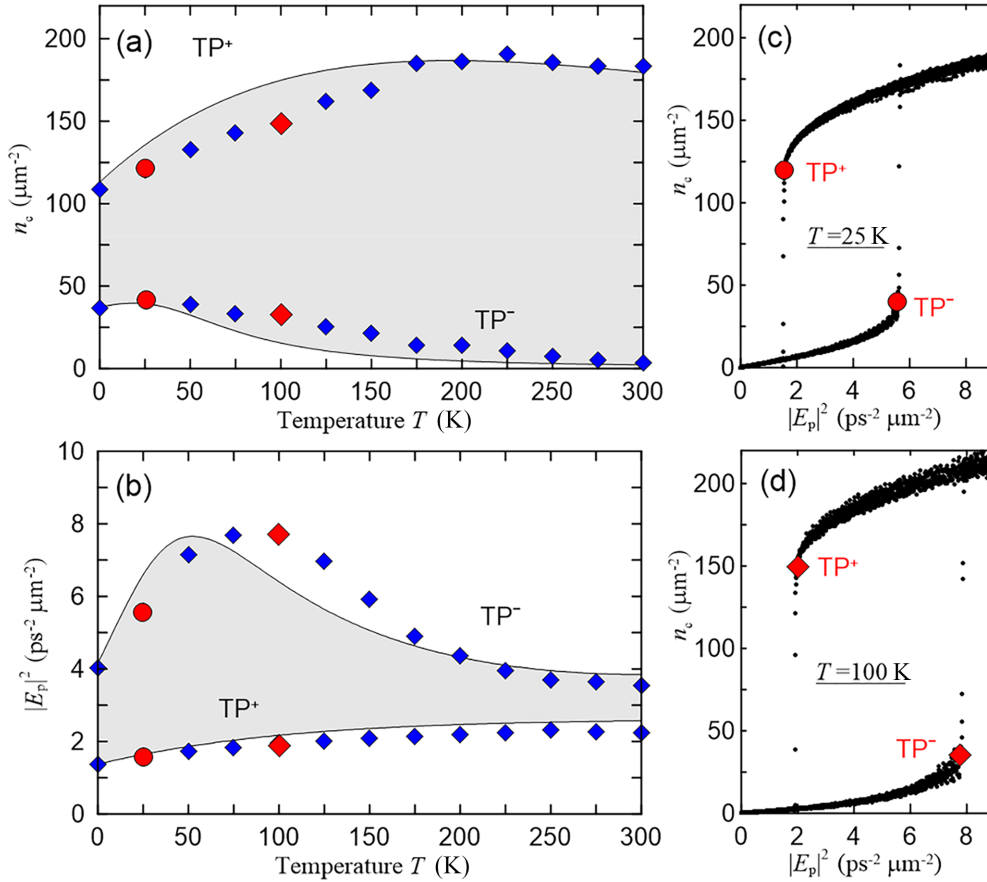


FIG. 5. Comparison of turning points $TP^\pm$ of HS solutions calculated with the full model [Eqs. (2.6)–(2.8)] (marked by rhombuses and circles) and the analytical approach [Eq. (4.2)] adapted to the presence of the incoherent reservoir (shaded area between the two solid lines). Panels (a) and (b) show the intracavity photon density n_c and the pump intensity $|E_p|^2$ as a function of temperature T , respectively. Panels (c) and (d) show two examples of hysteresis loops obtained by an adiabatic sweeping of the pump intensity for the temperatures $T = 25$ K and $T = 100$ K, respectively. Respective TPs are highlighted by larger symbols (circles and rhombuses) in the panels (a) and (b). Parameters are $\delta_{c,e} = 4.75$ ps $^{-1}$, $\gamma_c = 0.1$ ps $^{-1}$, $\gamma_e = 0.01$ ps $^{-1}$, $\Omega_R = 5$ ps $^{-1}$.

nonlinear optical response and, in particular, on the extent of the bistability domain of the resonator.

To our knowledge, direct experiments isolating the effects of a thermal acoustic phonon bath and detecting polariton hysteresis loops as a function of temperature are not yet available. However, several experimental studies support our findings. First, Rozas *et al.* [16] report the sensitivity of hysteresis cycles to additional scattering channels and reservoirs in GaAs microcavities. For instance, the occupation of a dark exciton reservoir narrows or collapses the hysteresis loop.

Second, Borri *et al.* [60] performed temperature-dependent linewidth measurements for GaAs microcavities, showing stronger line broadening for UP-like states than for LP-like states, which is in very good agreement with our findings displayed in Fig. 2(c). These observations align with our modeling of phonon-induced dephasing and energy renormalization.

To put our results in a broader context, we list the homogeneous linewidth broadening obtained with our model for different temperatures in Table I and compare it with the results obtained by Savona *et al.* [61] and with experimental results [16,60]. The homogeneous broadening, calculated

with our model beyond Fermi's golden rule, agrees very well with the theoretical results obtained by Savona *et al.* [61]

TABLE I. The exciton-induced linewidth broadening of LP and UP branches (i) obtained beyond Fermi's golden rule (this paper), (ii) calculated within Fermi's golden rule by Savona *et al.* [61], and (iii) experimental measurements by Borri *et al.* [60].

T (K)	This paper LP (meV)	Fermi's Golden LP (meV)	Experiment LP (meV)
10	0.0032	0.0004	0.00
25	0.0035	0.002	0.04
50	0.005	0.005	0.12
75	0.008	0.008	0.14
100	0.012	0.010	0.18
	UP (meV)	UP (meV)	UP (meV)
10	0.035	0.028	0.00
25	0.067	0.06	0.08
50	0.117	0.12	0.30
75	0.187	0.17	0.35
100	0.247	0.23	0.55

for a wide range of temperatures (from 10 K to 100 K) for both the lower and upper polariton branches. Nevertheless, compared with the experimental measurements [60], both theoretical approaches underestimate the exciton-induced linewidth broadening, especially for the lower polariton branch.

The enhanced linewidth broadening observed in the experiment could be explained by the coexistence of both exciton species, i.e., heavy-hole and light-hole excitons. An excess of free carriers in the microcavity owing to, e.g., unintentional doping of the Bragg reflectors might also have caused additional broadening as argued in [60]. Furthermore, the Rabi splitting measured in the experiment (3.6 meV) is almost half the Rabi splitting used in our calculations (6.582 meV). This could have further enhanced the dephasing of the lower polariton branch owing to the altered shape of the polariton dispersion.

In this context, we also note that the recent experimental study by Rozas *et al.* [16] assumed the damping rate of the lower-branch polaritons to be 0.283 meV for a GaAs microcavity measured at 10 K. Even after subtracting the photonic contribution (0.19 meV), this value is substantially larger than those shown in Table I for 10 K. This can be attributed to additional inhomogeneous broadening mechanisms, such as structural disorder [61], determined by the fabrication quality of the respective GaAs/AlGaAs microcavity.

Our model aims to provide a qualitative understanding of the underlying physics. A more realistic quantitative modeling would require the inclusion of additional degrees of freedom that enhance the dephasing dynamics by increasing the number of decay channels. For instance, considering spinor effects, including dark excitonic states (with spin projection of ± 2), will further improve the modeling of thermalization dynamics, in agreement with recent observations in TMD- [42] and GaAs- [16] based polariton systems, where the activation of additional dark exciton scattering channels leads to increased polariton linewidths and energy shifts.

While the spin of bright excitons (with spin projection of ± 1) is optically accessible, spinor dynamics introduces additional scattering and relaxation channels between the two circularly polarized components of microcavity polaritons [20,62]. We expect any imbalance between these polarization components to modify the thermalization dynamics of the reservoir and to introduce polarization-dependent dephasing or energy shifts of polaritons.

More generally, the role of phonon-induced dephasing and incoherent reservoir formation is expected to be relevant to a broad class of driven-dissipative polariton systems. From this perspective, a microscopic treatment of phonon-induced dephasing and reservoir formation, as developed in the present work, is relevant for assessing the robustness and thermal stability of polariton-based devices [38]. Finally, direct phonon-polariton signatures and significant couplings have been reported in the TMD platform [63], underlining the importance of phonons in the strong-coupling regime.

In conclusion, we scrutinized the effect of phonon dynamics on the bistable response of exciton-polaritons in semiconductor microcavities operating in the strong-coupling regime

and driven by a coherent optical pump field. We demonstrated that scattering from longitudinal acoustic phonons results not only in dephasing of coherent exciton-polaritons but also in the formation of an incoherent exciton reservoir. In particular, we found that coherent exciton-polaritons also interact nonlinearly with the reservoir, resulting in an additional frequency blueshift. This additional reservoir-induced frequency shift can reduce the bistability threshold and even increase the bistable domain.

We note that the proposed approach is appropriate for the description of the collective coherent dynamics of polaritons within an arbitrary two-dimensional potential landscape [49], and also for incoherent pumping conditions [33–36].

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DATA AVAILABILITY

There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

APPENDIX A: DERIVATION OF MASTER EQUATIONS

The model describes the dynamics of longitudinal acoustic phonons interacting with excitons via the deformation potential they induce. This potential accounts for the local change in the excitonic energy caused by phonon-induced lattice compression [44,45,53]. The Hamiltonian for 3D phonons interacting with 2D QW excitons can be simplified to

$$H_{\text{phonon}} = \sum_{q_z, \mathbf{q}} \hbar v_{\text{ph}} (|\mathbf{q}|^2 + q_z^2)^{1/2} u_{\mathbf{q}, q_z}^\dagger u_{\mathbf{q}, q_z} + \sum_{q_z} \sum_{\mathbf{q}, \mathbf{k}} G(\mathbf{q}, q_z) (u_{\mathbf{q}, q_z} - u_{-\mathbf{q}, -q_z}^\dagger) b_{\mathbf{k}}^\dagger b_{\mathbf{k}-\mathbf{q}}. \quad (\text{A1})$$

Here, $b_{\mathbf{k}}^\dagger (b_{\mathbf{k}})$ is the creation (annihilation) operator of a 2D QW exciton and $u_{\mathbf{q}, q_z}^\dagger (u_{\mathbf{q}, q_z})$ is the creation (annihilation) of a 3D phonon. The summations run over the in-plane wave vector components $\mathbf{q} = (m_x \mathbf{e}_x + m_y \mathbf{e}_y)^{2\pi/\sqrt{A}}$ and over the out-of-plane wave vector of the microcavity $q_z = m_z^{2\pi/L_z}$, where $m_{x,y,z}$ are integers and A is the quantization area of the assumed square QW. L_z is the thickness of the microcavity (z is the growth direction). All parameters of the interaction are contained in the coefficient,

$$G(\mathbf{q}, q_z) = -i\alpha(|\mathbf{q}|, q_z) \sqrt{\frac{\hbar(|\mathbf{q}|^2 + q_z^2)^{1/2}}{2\rho v_{\text{ph}} V}}. \quad (\text{A2})$$

Here, ρ and v_{ph} are the mass density and the longitudinal sound velocity of the semiconductor, respectively, and

$V = AL_z$ is the quantization volume. The effective deformation potential for the excitons is given by the coefficients of the respective potentials of electrons α_c and holes α_v ,

$$\alpha(|\mathbf{q}|, q_z) = \alpha_v I_h^\parallel(|\mathbf{q}|) I_h^\perp(q_z) - \alpha_c I_e^\parallel(|\mathbf{q}|) I_e^\perp(q_z).$$

The terms $I_{e(h)}^\parallel(|\mathbf{q}|)$ and $I_{e(h)}^\perp(q_z)$ are the overlap integrals of electrons and holes given by the exciton envelope and the phonon wave function (plane wave) in the in-plane direction and the z direction, respectively (for details, see [44,45,53]). These overlap integrals vanish for very high momenta, introducing cutoffs in Eq. (A2) for $|\mathbf{q}| > 1/a_B$ and $q_z > 2\pi/d$, where a_B the Bohr radius of exciton confined in the QW of width d . In our model, we assume an exciton radius $a_B = 10$ nm, which is typical for a GaAs QW with a width $d \approx a_B$ [45].

For simplicity, we neglect the momentum dependence of the exciton deformation potential and substitute the term $\alpha(|\mathbf{q}|, q_z)$ by the respective constant coefficient α_0 [see Eq. (2.5)], which can be estimated by the expression

$$\alpha_0 = \frac{a_B^2 d}{2\pi^2} \int_{|\mathbf{q}| < a_B^{-1}} d\mathbf{q} \int_{q_z < \pi/d} dq_z \alpha(\mathbf{q}, q_z).$$

Here, the prefactor $a_B^2 d / (2\pi^2)$ corresponds to the inverse of the integration volume in momentum space, such that α_0 represents the average value of $\alpha(\mathbf{q}, q_z)$ within the integration domain. For our calculations, we use the deformation potential constants $\alpha_c = -9$ eV and $\alpha_v = 2.7$ eV which are consistent with the material parameters for InGaAs QWs known from the literature [54,55]. From these values, we estimate the effective deformation potential of InGaAs QW excitons to be $\alpha_0 = 10.4$ eV.

1. Equation of motion for acoustic phonons

The aim of this section is to derive the master equation for the propagation of the lattice deformation (compression) waves within the bulk of the microcavity, and their interaction with excitons in the QW. The corresponding equation for the operators in the Heisenberg picture has the following form:

$$\begin{aligned} \partial_t u_{\mathbf{q}, q_z} &= \frac{i}{\hbar} [H_{\text{phonon}}, u_{\mathbf{q}, q_z}] = -iv_{\text{ph}} (|\mathbf{q}|^2 + q_z^2)^{1/2} u_{\mathbf{q}, q_z} \\ &+ \frac{\alpha_0}{\hbar} \sqrt{\frac{\hbar (|\mathbf{q}|^2 + q_z^2)^{1/2}}{2\rho L_z v_{\text{ph}}}} \varrho_{\mathbf{q}}, \end{aligned} \quad (\text{A3})$$

where the term on the right-hand side

$$\varrho_{\mathbf{q}} = \frac{1}{\sqrt{A}} \sum_{\mathbf{k}} b_{\mathbf{k}}^\dagger b_{\mathbf{k}+\mathbf{q}} = \mathcal{F}\{b_e(\mathbf{r})^\dagger b_e(\mathbf{r})\} \quad (\text{A4})$$

is the Fourier transform of the 2D exciton density operator calculated in the plane of the QW of area A . The evolution of phonons as described by Eq. (A3) is associated with the propagation of a longitudinal wave causing a lattice displacement described by the operator

$$\begin{aligned} \hat{\mathbf{w}}(\mathbf{r}, z) &= \sum_{\mathbf{q}, q_z} \sqrt{\frac{\hbar}{2\rho V v_{\text{ph}} (|\mathbf{q}|^2 + q_z^2)^{3/4}}} (\mathbf{q} + q_z \mathbf{e}_z) \\ &\times (u_{\mathbf{q}, q_z} e^{i(\mathbf{q} \cdot \mathbf{r} + q_z z)} + u_{\mathbf{q}, q_z}^\dagger e^{-i(\mathbf{q} \cdot \mathbf{r} + q_z z)}). \end{aligned} \quad (\text{A5})$$

We note that the normalization of Eq. (A5) ensures that the mechanical energy of the standing wave with a wave vector $(\mathbf{q}, q_z)$ is equal to the energy of all phonons carried by the respective mode [for details, see Eq. (B3) in Appendix B]. The longitudinal wave induces the relative deformation of the lattice given by

$$\begin{aligned} \hat{P}(\mathbf{r}, z) &= \nabla \cdot \hat{\mathbf{w}}(\mathbf{r}, z) = i \sum_{\mathbf{q}} \sum_{q_z} \sqrt{\frac{\hbar (|\mathbf{q}|^2 + q_z^2)^{1/2}}{2\rho V v_{\text{ph}}}} \\ &\times (u_{\mathbf{q}, q_z} e^{i(\mathbf{q} \cdot \mathbf{r} + q_z z)} - u_{\mathbf{q}, q_z}^\dagger e^{-i(\mathbf{q} \cdot \mathbf{r} + q_z z)}) \\ &= i \sum_{\mathbf{q}} \sum_{q_z} \sqrt{\frac{\hbar (|\mathbf{q}|^2 + q_z^2)^{1/2}}{2\rho V v_{\text{ph}}}} (u_{\mathbf{q}, q_z} - u_{-\mathbf{q}, -q_z}^\dagger) e^{i(\mathbf{q} \cdot \mathbf{r} + q_z z)}. \end{aligned} \quad (\text{A6})$$

These phonon-induced lattice deformations (compressions) change locally the exciton energy, thus creating a space-dependent potential for the excitons $\alpha_0 \hat{P}(\mathbf{r}, z)$, originating from the phonon–electron–hole interaction [44,53].

We proceed with the derivation of an equation for the evolution of the deformation $\hat{P}(\mathbf{r}, z)$. For this purpose, we apply a time derivative to Eq. (A3) and obtain a second-order differential equation for the annihilation operator $u_{\mathbf{q}, q_z}$,

$$\begin{aligned} \partial_t^2 u_{\mathbf{q}, q_z} &= -v_{\text{ph}}^2 (|\mathbf{q}|^2 + q_z^2) u_{\mathbf{q}, q_z} - i \frac{\alpha_0 v_{\text{ph}}}{\hbar} \sqrt{\frac{\hbar (|\mathbf{q}|^2 + q_z^2)^{1/2}}{2\rho L_z v_{\text{ph}}}} \\ &\times (|\mathbf{q}|^2 + q_z^2)^{1/2} \varrho_{\mathbf{q}} + \frac{\alpha_0}{\hbar} \sqrt{\frac{\hbar (|\mathbf{q}|^2 + q_z^2)^{1/2}}{2\rho L_z v_{\text{ph}}}} \partial_t \varrho_{\mathbf{q}}. \end{aligned}$$

A similar expression can be derived for the creation operator $u_{-\mathbf{q}, -q_z}^\dagger$, resulting in the following differential equation for their difference, which determines the lattice compression:

$$\begin{aligned} \partial_t^2 (u_{\mathbf{q}, q_z} - u_{-\mathbf{q}, -q_z}^\dagger) &= -v_{\text{ph}}^2 (|\mathbf{q}|^2 + q_z^2) (u_{\mathbf{q}, q_z} - u_{-\mathbf{q}, -q_z}^\dagger) \\ &+ i \frac{\alpha_0 v_{\text{ph}}}{\hbar} \sqrt{\frac{2\hbar}{\rho L_z v_{\text{ph}} (|\mathbf{q}|^2 + q_z^2)^{1/2}}} (|\mathbf{q}|^2 + q_z^2) \varrho_{\mathbf{q}}, \end{aligned}$$

where the symmetry of the exciton density operator in Fourier space $\varrho_{\mathbf{q}}^\dagger = \varrho_{-\mathbf{q}}$ [see Eq. (A4)] was used. Multiplying by the coefficients according to Eq. (A6) and performing a 3D Fourier transformation yields a wave equation for the lattice deformation $\hat{P}(\mathbf{r}, z)$ coupled to the excitons of the 2D layer of the QW (for $z = 0$) as

$$\begin{aligned} \partial_t^2 \hat{P}(\mathbf{r}, z) &= v_{\text{ph}}^2 (\Delta_\parallel + \partial_z^2) \hat{P}(\mathbf{r}, z) - \frac{\alpha_0}{\rho} \Delta_\parallel (b_e(\mathbf{r})^\dagger b_e(\mathbf{r})) \delta(z) \\ &+ \frac{\alpha_0}{\rho} (b_e(\mathbf{r})^\dagger b_e(\mathbf{r})) \frac{1}{L_z} \sum_{q_z} (-q_z^2) e^{iq_z z}. \end{aligned} \quad (\text{A7})$$

It is convenient to separate the free-propagating phonons in the bulk of the microcavity from those that interact with the excitons in the QW. To do so, we define the 2D acoustic phonons in the QW layer by integrating the lattice

deformation $\hat{P}(\mathbf{r}, z)$ over the QW with thickness d ,

$$\hat{P}_{\text{ph}}(\mathbf{r}, 0) \equiv \frac{1}{d} \int_{-d/2}^{d/2} \hat{P}(\mathbf{r}, z) dz \approx \hat{P}(\mathbf{r}, 0). \quad (\text{A8})$$

Integrating Eq. (A7) over the QW removes the term on the right-hand side corresponding to a second derivative of the delta distribution. We arrive at an equation for the 2D acoustic phonons (A8) within the QW layer,

$$\begin{aligned} \partial_t^2 \hat{P}_{\text{ph}}(\mathbf{r}, 0) &= v_{\text{ph}}^2 \Delta_{\parallel} \hat{P}_{\text{ph}}(\mathbf{r}, 0) - \frac{\alpha_0}{\rho d} \Delta_{\parallel} |b_e(\mathbf{r})|^2 \\ &+ \frac{v_{\text{ph}}^2}{d} \partial_z \hat{P}(\mathbf{r}, z) \Big|_{-d/2}^{d/2}. \end{aligned} \quad (\text{A9})$$

The last term on the right-hand side of Eq. (A9) describes the influx and outflux of phonons from the QW layer. In our model, we consider the two-dimensional dynamics of excitons interacting with the three-dimensional acoustic phonons propagating outside the QW layer in the microcavity material. Rather than introducing a phenomenological phonon bath at this stage, we proceed by deriving a wave equation for the lattice deformation field in the full three-dimensional microcavity, which naturally accounts for phonon propagation and thermal fluctuations.

Applying the same procedure at any position z_i along the out-of-plane z direction, within a “phonon” bath, we generalize the equation for acoustic phonons to the following form:

$$\begin{aligned} \partial_t^2 \hat{P}_{\text{ph}}(\mathbf{r}, z) &= v_{\text{ph}}^2 (\Delta_{\parallel} + \partial_z^2) \hat{P}_{\text{ph}}(\mathbf{r}, z) \\ &- \frac{\alpha_0}{\rho d} \Delta_{\parallel} (b_e(\mathbf{r})^\dagger b_e(\mathbf{r})) f(z), \end{aligned} \quad (\text{A10})$$

with the QW localization function defined as $f(z) \equiv 1$ within the interval $-d/2 < z < d/2$ and $f(z) \equiv 0$ otherwise. Initially (for $t = 0$), the phonons are assumed to be in thermal equilibrium with a predefined lattice temperature T (see Appendix B for further details).

Following a mean-field approximation, the field operators are treated as classical variables after a standard quantum-mechanical averaging procedure, $\psi_e(\mathbf{r}) = \langle b_e(\mathbf{r}) \rangle$, $|\psi_e(\mathbf{r})|^2 = \langle b_e(\mathbf{r})^\dagger b_e(\mathbf{r}) \rangle$ and $P(\mathbf{r}, z) = \langle \hat{P}(\mathbf{r}, z) \rangle$. In this case, the model described by Eq. (A10) becomes equivalent to Eq. (2.8) of the main text.

2. Phonon-mediated deformation potential in the equation of motion for excitons

In this section, we derive the term responsible for the scattering of QW excitons from the deformation potential created by longitudinal acoustic phonons. This term can be obtained in the Heisenberg representation from the Hamiltonian in Eq. (A1),

$$\begin{aligned} \partial_t b_{\mathbf{k}}|_{\text{exc-phon}} &= \frac{i}{\hbar} [H_{\text{phonon}, \mathbf{k}}, b_{\mathbf{k}}] \\ &= -\frac{i}{\hbar} \sum_{\mathbf{q}, q_z} G(\mathbf{q}, q_z) (u_{\mathbf{q}, q_z} - u_{-\mathbf{q}, -q_z}^\dagger) b_{\mathbf{k}-\mathbf{q}}. \end{aligned} \quad (\text{A11})$$

We switch to the real-space representation of Eq. (A11) by performing a 2D Fourier transformation in the plane of the

microcavity with the quantization area A ,

$$\begin{aligned} \partial_t \frac{1}{\sqrt{A}} \sum_{\mathbf{k}} b_{\mathbf{k}} e^{i\mathbf{k} \cdot \mathbf{r}}|_{\text{exc-phon}} \\ = -\frac{i}{\hbar} \frac{1}{\sqrt{A}} \sum_{\mathbf{k}} \sum_{\mathbf{q}, q_z} G(\mathbf{q}, q_z) (u_{\mathbf{q}, q_z} - u_{-\mathbf{q}, -q_z}^\dagger) b_{\mathbf{k}-\mathbf{q}} e^{i\mathbf{k} \cdot \mathbf{r}}. \end{aligned} \quad (\text{A12})$$

We then change the summation index $\mathbf{k}' = \mathbf{k} - \mathbf{q}$ and obtain

$$\begin{aligned} \partial_t \frac{1}{\sqrt{A}} \sum_{\mathbf{k}} b_{\mathbf{k}} e^{i\mathbf{k} \cdot \mathbf{r}}|_{\text{exc-phon}} \\ = \frac{-\alpha_0}{\hbar} \sum_{\mathbf{q}, q_z} \sqrt{\frac{\hbar(|\mathbf{q}|^2 + q_z^2)^{1/2}}{2\rho V v_{\text{ph}}}} (u_{\mathbf{q}, q_z} - u_{-\mathbf{q}, -q_z}^\dagger) \\ \times e^{i\mathbf{q} \cdot \mathbf{r}} \frac{1}{\sqrt{A}} \sum_{\mathbf{k}'} b_{\mathbf{k}'} e^{i\mathbf{k}' \cdot \mathbf{r}}. \end{aligned} \quad (\text{A13})$$

After introducing the space-dependent exciton operator $b_e(\mathbf{r}) = \frac{1}{\sqrt{A}} \sum_{\mathbf{k}} b_{\mathbf{k}} e^{i\mathbf{k} \cdot \mathbf{r}}$, Eq. (A13) reads as

$$\begin{aligned} \partial_t b_e(\mathbf{r})|_{\text{exc-phon}} &= \frac{i\alpha_0}{\hbar} \left[i \sum_{\mathbf{q}, q_z} \sqrt{\frac{\hbar(|\mathbf{q}|^2 + q_z^2)^{1/2}}{2\rho V v_{\text{ph}}}} \right. \\ &\times (u_{\mathbf{q}, q_z} - u_{-\mathbf{q}, -q_z}^\dagger) e^{i\mathbf{q} \cdot \mathbf{r}} e^{iq_z \cdot 0} \Big] b_e(\mathbf{r}). \end{aligned} \quad (\text{A14})$$

In accordance with Eq. (A6), the term in square brackets in Eq. (A14) is the lattice compression $\hat{P}(\mathbf{r}, z)$ induced by 3D longitudinal acoustic phonons in the plane of the QW ($z = 0$). Then, using the definition of the 2D phonons within the QW (A8), we simplify the equation of motion to

$$\partial_t b_e(\mathbf{r})|_{\text{exc-phon}} = \frac{i\alpha_0}{\hbar} \hat{P}_{\text{ph}}(\mathbf{r}) b_e(\mathbf{r}). \quad (\text{A15})$$

After averaging over the respective quantum field operators $\psi_e(\mathbf{r}) = \langle b_e(\mathbf{r}) \rangle$ and $P_{\text{ph}}(\mathbf{r}) = \langle \hat{P}_{\text{ph}}(\mathbf{r}) \rangle$, we obtain the equation for the classical variables [see Eq. (2.7) in the main text].

Equation (A15) shows that the lattice deformations induced by the longitudinal acoustic phonons modify locally the exciton energy, thus inducing a time-varying energy landscape for excitons in the QW layer.

APPENDIX B: THERMAL POPULATION OF THE PHONON BATH

The model [Eq. (A10)] describes the dynamics of a longitudinal-acoustic phonon bath that is in thermal contact with a single QW positioned at the antinode of a semiconductor microcavity.

Before proceeding with the derivation of the phonon-bath population, it is instructive to relate the quantum field operators to observable quantities. For this purpose, we consider the oscillation energy of the phonon mode $E_{\mathbf{q}, q_z}$ with a fixed momentum in the plane $\mathbf{q}$ and out of the plane q_z of the QW. A formal solution of the homogeneous part of Eq. (A3) for $\alpha_0 = 0$ has the form $u_{\mathbf{q}, q_z}(t) = u_{0, \mathbf{q}, q_z} e^{-i\omega_{\mathbf{q}, q_z} t}$ with the corresponding dispersion relation $\omega_{\mathbf{q}, q_z} = v_{\text{ph}}(|\mathbf{q}|^2 + q_z^2)^{1/2}$.

First, we derive the expression for the kinetic energy by considering the corresponding component of the lattice displacement $\hat{\mathbf{w}}_{\mathbf{q},q_z}(\mathbf{r}, z)$ given by Eq. (A5),

$$E_{\mathbf{q},q_z}^{\text{kin}} = \left\langle \int dV \frac{\rho}{2} \left| \partial_t \left(\sqrt{\frac{\hbar}{2\rho V \omega_{\mathbf{q},q_z}}} (u_{0,\mathbf{q},q_z} e^{i(\mathbf{q}\cdot\mathbf{r}+q_z z - \omega_{\mathbf{q},q_z} t)} + u_{0,\mathbf{q},q_z}^\dagger e^{-i(\mathbf{q}\cdot\mathbf{r}+q_z z - \omega_{\mathbf{q},q_z} t)}) \right) \right|^2 \right\rangle, \quad (\text{B1})$$

where ρ is the mass density of the semiconductor, and the integral is calculated over the quantization volume $V = AL_z$.

Calculating the time derivatives in Eq. (B1) simplifies the expression to

$$\begin{aligned} E_{\mathbf{q},q_z}^{\text{kin}} &= \left\langle \int dV \frac{\hbar \rho \omega_{\mathbf{q},q_z}^2}{4\rho V \omega_{\mathbf{q},q_z}} \left| -iu_{0,\mathbf{q},q_z} e^{i(\mathbf{q}\cdot\mathbf{r}+q_z z - \omega_{\mathbf{q},q_z} t)} + iu_{0,\mathbf{q},q_z}^\dagger e^{-i(\mathbf{q}\cdot\mathbf{r}+q_z z - \omega_{\mathbf{q},q_z} t)} \right|^2 \right\rangle \\ &= \frac{1}{4} \hbar \omega_{\mathbf{q},q_z} (\langle u_{0,\mathbf{q},q_z} u_{0,\mathbf{q},q_z}^\dagger \rangle + \langle u_{0,\mathbf{q},q_z}^\dagger u_{0,\mathbf{q},q_z} \rangle) = \frac{1}{4} \hbar \omega_{\mathbf{q},q_z} (1 + 2\langle u_{0,\mathbf{q},q_z}^\dagger u_{0,\mathbf{q},q_z} \rangle), \end{aligned} \quad (\text{B2})$$

where we implemented the commutation rules for bosonic operators and neglected the fast-oscillating terms, since their contribution vanishes after time averaging. Following the virial theorem, the total averaged energy of the oscillator is distributed equally between the kinetic and potential components. Therefore, we obtain the total energy of the mode as

$$\begin{aligned} E_{\mathbf{q},q_z} &= 2E_{\mathbf{q},q_z}^{\text{kin}} = \hbar \omega_{\mathbf{q},q_z} (1/2 + \langle u_{\mathbf{q},q_z}^\dagger u_{\mathbf{q},q_z} \rangle) \\ &\approx \hbar \omega_{\mathbf{q},q_z} \langle u_{\mathbf{q},q_z}^\dagger u_{\mathbf{q},q_z} \rangle. \end{aligned} \quad (\text{B3})$$

As we see from Eq. (B3), the total energy of the mode is indeed equal to the energy of a single phonon multiplied by the phonon number (obtained after averaging the quantum field operator $u_{\mathbf{q},q_z}^\dagger u_{\mathbf{q},q_z}$).

In thermodynamic equilibrium, the phonon occupation of modes is given by a Bose-Einstein distribution with an average phonon number per mode of

$$N_{\mathbf{q},q_z} = \langle u_{\mathbf{q},q_z}^\dagger u_{\mathbf{q},q_z} \rangle = \frac{1}{[\exp(\hbar \omega_{\mathbf{q},q_z}/k_B T) - 1]}, \quad (\text{B4})$$

where k_B is Boltzmann's constant, and T is the lattice temperature. After substitution of the mode amplitudes $\sqrt{N_{\mathbf{q},q_z}}$ into the expression for the lattice deformation in Eq. (A6), we can write the initial distribution ($t = 0$) as

$$\begin{aligned} P_{\text{ph}}(\mathbf{r}, z) &= \sum_{\mathbf{q}} \sum_{q_z} \sqrt{\frac{2\hbar \omega_{\mathbf{q},q_z}}{\rho V v_{\text{ph}}^2 [\exp(\hbar \omega_{\mathbf{q},q_z}/k_B T) - 1]}} \\ &\times \sin(\mathbf{q} \cdot \mathbf{r} + q_z z + \chi_{\mathbf{q},q_z}), \end{aligned} \quad (\text{B5})$$

where $\chi_{\mathbf{q},q_z}$ are arbitrarily chosen random phases of the phonon waves. Equation (B5) describes the initial distribution of the lattice deformation [see Eq. (2.8)], in a phonon bath with the predefined environmental temperature T .

APPENDIX C: NUMERICAL METHODS AND SYSTEM-SIZE ANALYSIS

The coupled equations of motion for the cavity-photon and QW-exciton fields [Eqs. (2.6) and (2.7)] were solved numerically using a split-step Fourier method under periodic in-plane boundary conditions. It is important to note that the phonon coupling term in the exciton evolution equation [Eq. (2.7)] is nonlinear, and the phonon wave equation [Eq. (2.8)] was solved in Fourier space.

The simulation domain for the exciton and photon fields corresponds to a two-dimensional area of $800 \times 800 \text{ nm}^2$, while the phonon field was modeled in a three-dimensional box of size $800 \times 800 \times 1000 \text{ nm}^3$, accounting for the vertical extent of phonon propagation in the microcavity structure. The thermal phonon field of Eq. (B5) is initialized by assigning to each discrete mode an independent random phase $\chi_{\mathbf{q},q_z}$ (sampled once at $t=0$) and a Bose-Einstein modal amplitude proportional to $\sqrt{N_{\mathbf{q},q_z}}$ [see Eq. (B4)]. This procedure ensures that the initial lattice deformation corresponds to a thermally populated phonon bath with random spatial phase correlations. The random phases define the initial spatial configuration (realization) of the deformation potential in the QW plane, while the subsequent time evolution of the phonon field, governed by the wave equation [Eq. (A10)], produces the expected temporal and spatial fluctuations of the deformation potential.

A spatial discretization of $\Delta x = \Delta y = 4.4 \text{ nm}$ was used in the in-plane direction for the coupled Eqs. (2.6)–(2.8), and $\Delta z = 5 \text{ nm}$ in the out-of-plane (z) direction for phonons. The temporal step size for the time evolution was chosen to be $\Delta t = 10^{-4} \text{ ps}$, which ensures numerical stability and accuracy of the split-step method.

To confirm the reliability of the numerical results, we performed a convergence analysis with respect to both spatial and temporal discretizations. The mesh size and time step were systematically varied, and the deviation in key physical observables (e.g., intracavity field intensity and bistability thresholds) was found to be less than 2% when refining the mesh or reducing Δt by a factor of two. This confirms the local and global convergence of our simulations.

The selected simulation box size allows us to focus on phonon-induced effects while suppressing modulational instabilities of the exciton-polariton system [19,26,56]. Such instabilities correspond to the destabilization of spatially homogeneous steady states owing to the growth of finite-wave-vector perturbations driven by the interplay of dispersion, nonlinearity, and external pumping. In resonantly driven polariton systems, modulational instability can lead to spontaneous pattern formation, filamentation, or fragmentation of the homogeneous field [19,26,56].

Moreover, the reduced geometry substantially lowers the computational cost of solving the fully coupled microcavity

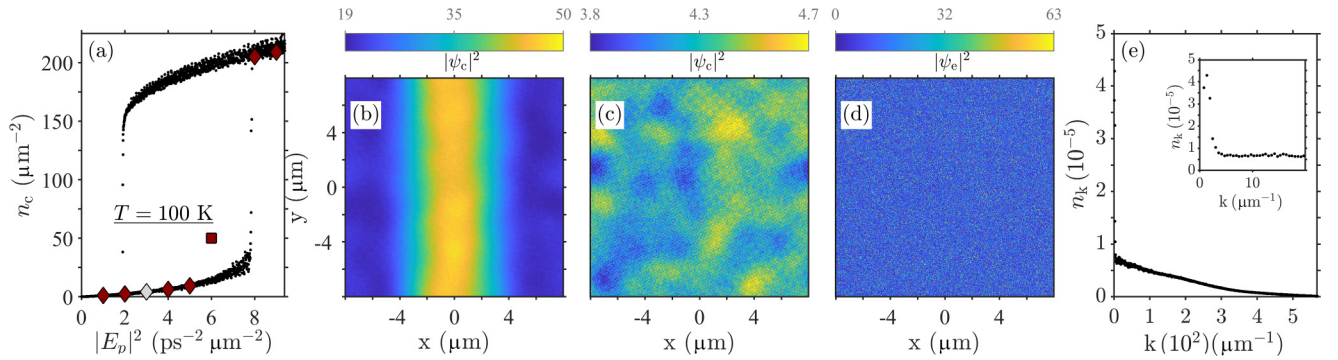


FIG. 6. Bistability characteristics and corresponding field profiles of the coupled exciton–photon–phonon system. (a) Cavity photon density n_c as a function of the pump intensity $|E_p|^2$ at $T = 100$ K. The dotted line shows the bistability loop for the smaller transversal domain ($800 \times 800 \text{ nm}^2$) [see Fig. 5(d)], while filled markers for the larger system ($1.6 \times 10^4 \times 1.6 \times 10^4 \text{ nm}^2$). (b) Photon density $|\psi_c|^2$ at the pump intensity corresponding to the filled square in (a), showing the emergence of spatial modulations. (c) Photon density $|\psi_c|^2$ at the steady-state condition [filled gray marker in (a)], exhibiting a nearly homogeneous distribution. (d) Exciton density $|\psi_e|^2$, including both the coherently driven exciton field and the incoherent reservoir contribution. (e) Momentum-space exciton population n_k corresponding to (d), normalized to its zero-momentum value $n_{k=0}$, as a function of the in-plane wave number k . The inset highlights the coherently driven exciton population at around $k = 0$.

exciton–photon–3D phonon system over long time intervals, allowing for systematic parameter sweeps.

To further verify the validity of our model and assess the effects of increasing the system size, we relax this restriction. Specifically, we (i) compare the bistability characteristics for small and large simulation domains and (ii) present representative two-dimensional spatial profiles of the coupled photon and exciton densities in an extended geometry.

Figure 6(a) compares the bistability characteristics obtained for the small transverse domain ($800 \times 800 \text{ nm}^2$, dotted lines) used in the main text with those calculated for a significantly larger system ($1.6 \times 10^4 \times 1.6 \times 10^4 \text{ nm}^2$, filled markers) at 100 K. Close to the lower bistability branch boundary (TP[−]) [square marker in Fig. 6(a)], the cavity photon density [panel (b)] exhibits pronounced spatial modulations, indicating the emergence of modulational instabilities with increased system size. Panels (c) and (d) show the steady-

state particle density at the pump intensity corresponding to the gray marker in panel (a). Away from the lower bistability branch boundary, the photon density is spatially homogeneous on average, with small persistent fluctuations about the mean density arising from phonon coupling. The exciton density, in contrast, contains both a coherent component that couples directly to the cavity photons and an incoherent reservoir populated by phonon-assisted scattering processes. Panel (e) shows the exciton population corresponding to panel (d) as a function of the in-plane wave number $k = \sqrt{k_x^2 + k_y^2}$, normalized to $n_{k=0}$. The inset shows the pronounced peak around $k = 0$ highlighting the coherent exciton component, whereas for $k > 0$ the distribution decreases smoothly with increasing wave number without additional narrow features, which is compatible with the formation of an incoherent exciton reservoir mediated by phonon-assisted scattering.

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