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Low intensity directional switching of light in semiconductor microcavities

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Recently it was demonstrated that in atomic vapors weak control beams can manipulate (or switch) the propagation direction of strong light beams [Dawes et al., Science **308**, 672 (2005)]. As a semiconductor analog of such all-optical switching, we present a proposal for similar manipulation and switching in planar semiconductor microcavities. Using a microscopic many-particle theory, we investigate the spatio-temporal dynamics of four-wave mixing signals and related instabilities in these systems. Even though the underlying physical processes are different from atomic systems, we find that microcavities allow for *reversible directional manipulation* of light.

1 Introduction Recently, an all-optical switch operating at very low light intensities was demonstrated in an atomic vapor system [1]. For two pump beams counter-propagating through an atomic vapor, four-wave mixing (FWM) instability-induced spontaneous off-pump-axis pattern formation can occur [2]. With a slight breaking of the system symmetry to impose a preferred direction to these off-pump-axis patterns, it was demonstrated that a very weak "control" beam is sufficient to redirect the patterns away from the preferred direction into the control's direction [1]. In this way, an all-optical switch was realized, where the actual switching signal is much stronger than the control.

A semiconductor analog of this effect is without doubt desirable. In this Letter, using a semiconductor microcavity, we present a specific conceptual proposal for a semiconductor implementation of the switching phenomenon reported in [1]. Our analysis is firmly based on a microscopic many-particle theory combined with numerical simulations of the spatio-temporal polariton dynamics. Our approach allows us to identify the physical processes underlying FWM instabilities and the directional switching action. We find that, on a microscopic level, the switching action is related to stimulated polariton scattering, an effect that has been observed and well characterized in the recent past (e.g., [3, 4]). As in the atomic case, we find that in the studied microcavity system, a weak control beam can steer a stronger signal beam into its own (the control's) propagation direction. As a crucial ingredient for possible future switching device applications, in the slightly anisotropic microcavity system under investigation, the switching action is found to be (spontaneously) reversed upon switching off the control.

2 Theory We analyze the nonlinear cavity-polariton dynamics in a typical planar GaAs microcavity [5]. The linear polariton dispersion is shown in Fig. 1(a). Our theory is a $\chi^{(3)}$ Hartree–Fock (HF) theory in the coherent limit, evaluated self-consistently up to arbitrary order in the optical fields [6]. The polariton interactions are strongly spin-dependent [5, 7, 8]. Here we concentrate on the dynamics in one spin subsystem (say spin up) by choosing circularly polarized excitation. We neglect the longitudinal–transverse (TE–TM) polariton-dispersion splitting [9] and use the optical dipole selection rules and matrix elements appropriate in quasi-normal incidence [10] (for a complete vectorial formulation of the theory, see [5]). In order to study the spatio-temporal dynamics of the system, we apply a spatial decomposition of cavity field and exciton polarization into Fourier components $E_{\boldsymbol{k}}$ and $p_{\boldsymbol{k}}$, respectively, with in-plane momentum $\boldsymbol{k}$ [10]. The exciton dynamics is evaluated in the 1s heavy-hole exciton subspace [11]. The nonlinear set of equations of motion for $E_{\boldsymbol{k}}$ and $p_{\boldsymbol{k}}$ reads [5, 10]:

$$i\hbar\dot{E}_{\boldsymbol{k}} = \hbar\omega_{\boldsymbol{k}}^{c} E_{\boldsymbol{k}} - \Omega_{\boldsymbol{k}}\, p_{\boldsymbol{k}} + i\hbar t_{c} E_{\boldsymbol{k},\mathrm{inc}}^{\mathrm{eff}}\,, \tag{1}$$

$$i\hbar\dot{p}_{\boldsymbol{k}} = (\varepsilon_{\boldsymbol{k}}^{x} - i\gamma_{x})\, p_{\boldsymbol{k}} - \Omega_{\boldsymbol{k}} E_{\boldsymbol{k}} + \sum_{\boldsymbol{qk'k''}} (2\tilde{A}\Omega_{\boldsymbol{k''}}\, p_{\boldsymbol{q}}^{*} p_{\boldsymbol{k'}} E_{\boldsymbol{k''}} + V_{\mathrm{HF}}\, p_{\boldsymbol{q}}^{*} p_{\boldsymbol{k'}} p_{\boldsymbol{k''}})\, \delta_{\boldsymbol{q},\boldsymbol{k'}+\boldsymbol{k''}-\boldsymbol{k}}\,. \tag{2}$$

The cavity field in Eq. (1) is treated in quasi-mode approximation. The effective incoming field $E_{\boldsymbol{k},\mathrm{inc}}^{\mathrm{eff}}$ driving the field $E_{\boldsymbol{k}}$ in the cavity is obtained from a simple transfer-matrix formalism that includes the radiative width ($\Gamma = \omega\hbar^{2}t_{c}^{2}/(\varepsilon_{0}cn_{b})$, with the background refractive index n_{b}, the vacuum velocity of light c and dielectric constant ε_{0}) of the cavity mode and yields transmitted and reflected field components: $E_{\boldsymbol{k},\mathrm{inc}}^{\mathrm{eff}} = E_{\boldsymbol{k},\mathrm{trans}} = E_{\boldsymbol{k},\mathrm{inc}} - E_{\boldsymbol{k},\mathrm{refl}}$ with $E_{\boldsymbol{k},\mathrm{refl}} = -[\hbar t_{c}/(2n_{b}c\varepsilon_{0})]\, E_{\boldsymbol{k}}$. The cavity-to-outside coupling constant t_{c} is chosen such that $\Gamma \approx 0.4$ meV for $\hbar\omega = 1.5$ eV. We include excitonic phase-space filling (PSF) and HF exciton–exciton Coulomb interaction in Eq. (2); two-exciton correlations are neglected and are expected to give merely quantitative changes because the pump is tuned far (several meV) below the bare exciton resonance [10, 11] (cf. Fig. 1(a)). The bare exciton and cavity in-plane dispersions are denoted by $\varepsilon_{\boldsymbol{k}}^{x}$ (with $\varepsilon_{0}^{x} = 1.497$ eV) and $\omega_{\boldsymbol{k}}^{c}$, with $\hbar\omega_{\boldsymbol{k}}^{c} = \varepsilon_{0}^{x}/\cos\vartheta$ and $\sin\vartheta = |\boldsymbol{k}|\,c/(\omega n_{b})$. A dephasing $\gamma_{x} = 0.4$ meV is included for the exciton polarization, $\Omega_{\boldsymbol{k}} = 8$ meV is the vacuum Rabi splitting, and $\tilde{A} = A_{\mathrm{PSF}}/\phi_{1s}^{*}(0)$ (with $A_{\mathrm{PSF}} = 4a_{0}^{x}\sqrt{2\pi/7}$, the bulk Bohr radius $a_{0}^{x} \approx 170$ Å, and the two-dimensional 1s wavefunction $\phi_{1s}(\boldsymbol{r})$ at $\boldsymbol{r} = 0$) and $V_{\mathrm{HF}} = 2\pi(1 - 315\pi^{2}/4096)/a_{0}^{x2} E_{b}^{x}$ (with the binding energy $E_{b}^{x} \approx 13$ meV) are the excitonic PSF and HF Coulomb matrix elements, respectively. A spatial anisotropy in the system can be modeled, e.g., by including an anisotropic cavity dispersion $\omega_{\boldsymbol{k}}^{c}$.

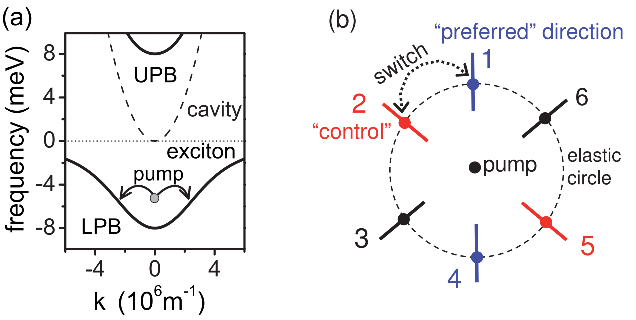

Figure 1 (online colour at: www.pss-rapid.com) (a) Cavity and exciton dispersions, and lower (LPB) and upper (UPB) polariton branches of the coupled cavity-mode exciton system. The fundamental pairwise off-pump-axis scattering of polaritons is indicated. (b) Hexagonal switching geometry in the transverse plane. The basic switching action triggered by the control beam is indicated. The radial bars indicate the variation in the magnitude of momenta k as included in the nonlinear dynamics.

We consider steady-state pump excitation in normal incidence ($\boldsymbol{k}_{\mathrm{pump}} = 0$), spectrally below the bare exciton resonance and above the lower polariton branch (cf. Fig. 1(a)). For this excitation configuration, FWM processes triggered by fluctuations in the cavity photon field can give rise to resonant phase-matched pairwise scattering of pump-excited polaritons into two polaritons with finite and opposite in-plane momentum $\boldsymbol{k}$ and $-\boldsymbol{k}$ (Fig. 1(a)). For a pump-induced exciton density above the instability threshold ($\approx 10^{10}$ cm^{-2} here) these scattering processes can lead to strong off-pump-axis signals (cf. [12]). At moderate pump intensity and in a system with a small symmetry breaking (anisotropy), hexagons and their subsets were the favored instability-induced patterns in Ref. [1]. In Eqs. (1) and (2), we allow for signals in the pump direction ($\boldsymbol{k}_{\mathrm{pump}} = 0$) and six off-pump directions forming hexagons in the transverse plane (Fig. 1(b)). Contributions to $E_{\boldsymbol{k}}$ and $p_{\boldsymbol{k}}$ with finite in-plane momentum $\boldsymbol{k} \neq 0$ are restricted to $k_{\min} < |\boldsymbol{k}| < k_{\max}$. The lower ($k_{\min} \approx 2.0 \times 10^{6}$ m^{-1}) and upper ($k_{\max} \approx 2.7 \times 10^{6}$ m^{-1}) momentum cut-offs are chosen such that the elastic circle (including the dynamical nonlinear renormalizations) lies within the numerical domain. This approximation can be justified as long as the pump excitation takes place spectrally below the "magic angle" [12].

3 Results and discussion In Fig. 2, we show results where we have numerically integrated Eqs. (1) and (2) for quasi steady-state pump excitation in normal incidence. The pump frequency is tuned 5 meV below the bare exciton resonance. The pump (not shown) reaches its peak intensity $I_{\mathrm{pump}} \approx 19.5$ kW cm^{-2} shortly after 0 ps and is then kept constant. We impose a slight anisotropy in the cavity dispersion by shifting $\omega_{\boldsymbol{k}}^{c}$ to lower energies by 0.075 meV in directions 1 and 4. Above a certain pump threshold intensity, stimulated scattering of pump-induced polaritons, driven mainly by the HF term in Eq. (2), leads to spontaneous (fluctuation-triggered) off-pump-axis signal formation (with intensity ≈ 1.5 W cm^{-2}). Initially, signals in all considered off-pump directions grow simultaneously. However, as these signals grow over time, the anisotropy (symmetry breaking) fixes the spontaneous signals at directions 1 and 4 (Fig. 2, times <2 ns). (About 10% of the pump-induced polaritons are scattered off-axis.) After 2 ns, we apply a weak control ($I_{\mathrm{control}} \approx 0.1$ W cm^{-2}) with the same frequency as the pump in direction 2 (Fig. 2(d)). Now, the strong off-pump-axis signal switches to directions 2 and 5 and vanishes in the "preferred" directions 1 and 4. Note that the switching signal in directions 2 and 5 is about *15 times stronger* than the control itself (i.e., part of the pump is redirected from normal incidence into directions 2 and 5). When switching off the weak control at ≈ 5 ns, the strong off-pump-axis signal switches back to the preferred directions 1 and 4. The switching can then be repeated. On/off switching times in our study are ≈ 1 ns, cor-

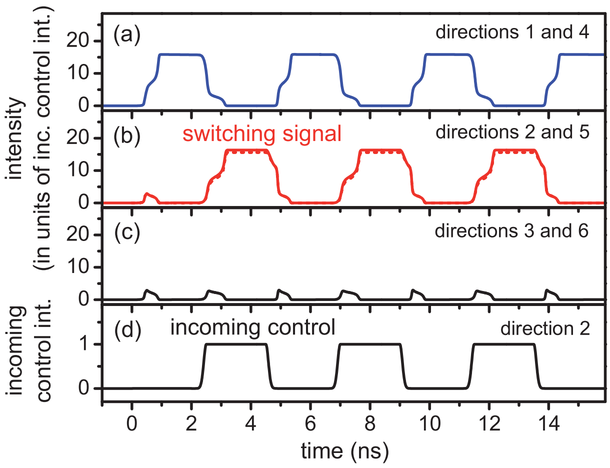

Figure 2 (online colour at: www.pss-rapid.com) (a)–(c) Switching in the output signals in a reflection geometry (the signals with out-of-plane momentum opposing the incident pump's are plotted). The intensities per direction are normalized to the incoming control intensity. The switching signal in (b) is about 15 times stronger than the incoming control in (d) that is triggering this signal (note the different scales on the vertical axes in panels (a)–(c) and (d)). In panel (b), direction 2 is shown as the solid line and direction 5 as the dashed line. Similar switching is observed in a transmission geometry (not shown).

responding to switching brought about by ≈13 photons if a beam waist of 2 μm diameter is assumed.

Figure 2 demonstrates the switching for one set of parameters. However, the general mechanism is robust. By fine-tuning parameters in the relatively large parameter space (e.g., pump and control intensities and frequencies), the switching performance can be optimized further, e.g., to achieve larger gain or lower switching intensities. Since the total strength of the off-pump-axis signals mainly depends on pump parameters, the largest gain is typically achieved for the lowest control intensity that overcomes the built-in anisotropy and thus triggers the switching process. We also note that the system dynamics drastically changes when the pump intensity and/or frequency are chosen such that bistability [13] plays a role. Finally, we briefly discuss possible limitations of the proposed scheme. The pump excitation is off-resonant. Thus, a relatively strong pump is required to reach the instability threshold. In an experimental setup, unintended off-pump-axis scattering of light can reduce the contrast ratio between "on" and "off" states. However, this practical issue might be alleviated using another existing microcavity design [14] with resonant pump excitation. Based on the simplifying assumption that the pump-induced exciton density is approximately $\sim(\gamma^2 + \Delta^2)^{-1}$, the threshold intensity would be reduced by about two orders of magnitude (for $\gamma = 0.4$ meV, and pump detuning $\Delta = 3$ meV from the bottom of the lower polariton branch). Also, fully two-dimensional simulations beyond the hexagonal geometry considered in this work would be desirable for the future.

4 Conclusions and remarks Our theoretical studies show that despite fundamental differences with atomic system, in the studied microcavity systems a strong optical beam can be *reversibly* controlled with a weaker one. Experiments on stimulated polariton scattering in microcavities have shown an impressive trend towards higher operational temperatures [4, 14]. However, one must realize that our analysis, is only a first step towards a "real-life" application. Apart from most desirable experimental investigations of the switching, it would be worthwhile to study vectorial polarization effects [5, 8], and to transfer the analysis to other systems [14, 15].

After arXiv upload and conference submission of our results [16], an independent but closely related study was published [17]. Kheradmand et al. investigate all-optical switching by pattern selection in a semiconductor microresonator. The microresonator response is modelled using an atomic nonlinearity caused by incoherent carriers. Spatial redirection of patterns with beams two orders of magnitude weaker than the actual patterns is achieved in a system with perfect in-plane isotropy. Pattern dynamics in an anisotropic system, including spontaneously reversible switching – the subject of our present paper – was not studied in [17].

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References

[1] A. M. C. Dawes et al., Science **308**, 672 (2005).
[2] A. Yariv and D. M. Pepper, Opt. Lett. **1**, 16 (1977).
G. Grynberg, Opt. Commun. **66**, 321 (1988).
R. Chang et al., Opt. Commun. **88**, 167 (1992).
[3] R. Huang et al., Phys. Rev. B **61**, R7854 (2000).
P. G. Savvidis et al., Phys. Rev. Lett. **84**, 1547 (2000).
C. Ciuti et al., Phys. Rev. B **62**, R4825 (2000).
D. M. Whittaker, Phys. Rev. B **63**, 193305 (2001).
[4] M. Saba et al., Nature **414**, 731 (2001).
[5] S. Schumacher et al., Phys. Rev. B **76**, 245324 (2007).
[6] T. Östreich and L. J. Sham, Phys. Rev. Lett. **83**, 3510 (1999).
[7] A. Kavokin et al., Phys. Rev. B **67**, 195321 (2003).
P. Renucci et al., Phys. Rev. B **72**, 075317 (2005).
[8] S. Schumacher, Phys. Rev. B **77**, 073302 (2008).
[9] V. Savona et al., Phys. Rev. B **53**, 13051 (1996).
[10] S. Savasta et al., Phys. Rev. Lett. **90**, 096403 (2003).
[11] N. H. Kwong et al., Phys. Rev. Lett. **87**, 027402 (2001).
[12] M. Romanelli et al., Phys. Rev. Lett. **98**, 106401 (2007).
[13] M. Wouters and I. Carusotto, Phys. Rev. B **75**, 075332 (2007).
[14] C. Diederichs et al., Nature **440**, 904 (2006).
[15] S. Schumacher et al., Appl. Phys. Lett. **92**, 131109 (2008).
[16] S. Schumacher et al., arXiv.org:cond-mat/0712.2060v1 (2007); CLEO/QELS Conference Proceedings, No. QMC4 (2008).
[17] R. Kheradmand et al., Eur. Phys. J. D **47**, 107 (2008).

