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Incoherent excitation and switching of spin states in exciton-polariton condensates

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We investigate, theoretically and numerically, the spin dynamics of a two-component exciton-polariton condensate created and sustained by nonresonant spin-polarized optical pumping in a semiconductor microcavity. Using the open-dissipative mean-field model, we show that the existence of well defined phase-locked steady states of the condensate may lead to efficient switching and control of spin (polarization) states with a nonresonant excitation. Spatially inhomogeneous pulsed excitations can cause symmetry breaking in the pseudospin structure of the condensate and lead to formation of nontrivial spin textures. Our model is universally applicable to two linearly coupled polariton condensates, and therefore can also describe the behavior of condensate populations and phases in “double-well” type potentials.

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I. INTRODUCTION

Created in semiconductor microcavities, two-component spinor exciton-polariton condensates [1,2] provide a rich resource for exploring the effects of internal (spin) degrees of freedom on the dynamics of nonequilibrium superfluids. A remarkable range of spin-related phenomena have been explored in systems with a coherent resonant optical pump, which offers the possibility to directly inject the spin state of exciton-polaritons and explore ultrafast spin-switching dynamics [3–6]. In contrast, the nonresonant incoherent excitation responsible for the spontaneous formation of coherent exciton-polariton condensates is only beginning to be experimentally explored in the context of spin-dependent effects. The reason for this disparity is a long-held belief that the incoherent reservoir (uncondensed hot exciton-polaritons), inevitably created by the incoherent excitations far above the energy of the condensed exciton-polariton quasiparticles, lacks spin-selectivity; i.e., owing to the fast spin-relaxation processes within the reservoir, the pump polarization cannot affect the spin of the condensed polaritons. Thus, many theoretical models of an incoherently excited exciton-polariton system with spin degrees of freedom treat the reservoir as a nonpolarized scalar density state [3,7].

Recent experimental studies of the optical spin-Hall effect [8] employed microcavities with relatively short exciton lifetimes and succeeded in creating a *spin-polarized* reservoir of high-energy “excitonic” polaritons. Such a reservoir can be controllably replenished by a pump laser of a selected polarization, thus opening up the way to create and manipulate spinor condensates of exciton-polaritons in the regime of incoherent far off-resonant optical excitation. The experiments confirming the spin-polarized nature of the reservoir call for a general analysis of spin dynamics under off-resonant excitation conditions, which will assist in understanding stationary spin structures and spin waves [9–11], as well as in the design of optically controlled transistors and switching devices based on the dynamical selection of polariton spin states [12–14].

In this work, we consider the exciton-polariton condensate in the regime of cw far off-resonant incoherent excitation, in

the absence of an external magnetic field. We analyze the non-stationary dynamics of condensate pseudospin (polarization vector) and the relaxation dynamics towards possible steady states for both spatially homogeneous and inhomogeneous cw pumps. We show that the spin-polarized incoherent reservoir allows efficient switching between different spin states of the polariton condensate, both in the spatially homogeneous and some inhomogeneous pumping regimes. Furthermore, we analyze perturbations of steady states by a spatially inhomogeneous optical pulse and demonstrate that heavily damped spin waves excited by the perturbations can lead to the formation of stable nontrivial spin textures. It should be noted that our general analysis, with a suitable modification of physical parameters, can be applied to a quasispin model [15] of weakly coupled, spatially separated polariton condensates, to determine the steady state populations and relative phases, as well as possible dynamical regimes (see, e.g., Ref. [16]).

II. THE MODEL

We model the spontaneously created polariton condensate by the open-dissipative Gross-Pitaevskii (GP) equations [17] derived from the Hamiltonian for a quantum fluid with pseudospin degrees of freedom obtained, in different contexts, in Refs. [3,4,7,18]. The dynamical equations for the condensate wave functions corresponding to the left ψ_- and right ψ_+ handed circular polarization states are written as follows ($\sigma = \pm$) [8]:

$$i\hbar \frac{\partial \psi_\sigma}{\partial t} = \left(-\frac{\hbar^2}{2m} \nabla^2 + u_a |\psi_\sigma|^2 + u_b |\psi_{-\sigma}|^2 + g R n_\sigma + \frac{i\hbar}{2} [R n_\sigma(r) - \gamma_c] \right) \psi_\sigma + J \psi_{-\sigma}. \quad (1)$$

Here m is the effective mass of the lower polariton, γ_c is the loss rate of polaritons which is the inverse of the polariton life time, and R characterizes the stimulated scattering rate from the reservoir into the condensate. The interaction constants u_a and u_b characterize the scattering between the polaritons of the

same and different polarization (spin) states, respectively. It is well established that $|u_b| < |u_a|$ [19]. Also, g_R characterizes interactions between the reservoir and condensate, where for simplicity we assume that interactions between oppositely spin polarized reservoir and condensed polaritons are negligible. The linear spin-mixing term proportional to J (Josephson coupling) accounts for the sample-specific linear polarization splitting observed in experiments [3,4] and usually arising due to stress at the semiconductor interfaces.

The equations for polariton wave functions are coupled to the equations for the “hot,” exciton-like reservoir created by the incoherent cw pump. It is usually assumed that polarization of the incoherent pump cannot be transferred to the polariton condensate [3,18]. However, as mentioned in the Introduction, the creation of a spin-polarized reservoir is possible with a polarized optical pump for sufficiently short-lived reservoir polaritons. To begin with, we will consider the simplest model of a spin-dependent reservoir, which was introduced in [8], and discuss the possible effects of cross-spin stimulated scattering later:

$$\frac{\partial n_\sigma}{\partial t} = P_\sigma(r, t) - (\gamma_R + R|\psi_\sigma|^2)n_\sigma, \quad (2)$$

where γ_R is the decay rate of the reservoir polaritons and $P_\sigma(r, t)$ is the pumping rate. The time scale of the short-lifetime reservoir dynamics is assumed to be comparable to that of the condensate, so that the model of a static reservoir [7,20] is no longer applicable. We note that the coherent resonant injection of spin [4] can be readily introduced into these equations.

By introducing the characteristic time $T = \gamma_c^{-1}$, length $L = \sqrt{\hbar/(m\gamma_c)}$, and energy $E_u = \hbar\gamma_c$ scales [21], we can rewrite Eqs. (1) and (2) in dimensionless form:

$$\begin{aligned} i\partial_t \psi_\sigma &= \left\{ -\frac{1}{2}\nabla^2 + u_a|\psi_\sigma|^2 + u_b|\psi_{-\sigma}|^2 + g_R n_\sigma \right. \\ &\quad \left. + \frac{i}{2}[Rn_\sigma - \gamma_c] \right\} \psi_\sigma + J\psi_{-\sigma}, \\ \partial_t n_\sigma &= P_\sigma(r) - (\gamma_R + R|\psi_\sigma|^2)n_\sigma, \end{aligned} \quad (3)$$

where $u_a \rightarrow u_a/(E_u L^2)$, $u_b \rightarrow u_b/(E_u L^2)$, $g_R \rightarrow g_R/(E_u L^2)$, and $J \rightarrow J/E_u$. In these scaling units, $\gamma_c = 1$; however, we will retain it in the subsequent formulas for clarity. In what follows, unless specified otherwise, all parameters take the physical values given in Ref. [21].

III. HOMOGENEOUS STEADY STATES

We start with the analysis of the spin-mixing dynamics for a spatially homogeneous condensate. In experiments, the quasihomogeneous distribution of condensate density can be achieved, e.g., by a flat-top “super-Gaussian” excitation. In this case, it is convenient to separate the amplitude and phase dynamics of the order parameter by using the transformation $\psi_\sigma(t) = \sqrt{\rho_\sigma(t)} \exp[i\phi_\sigma(t)]$, where $\rho_\sigma(t)$ is the density of the spinor component. The relative phase difference is defined as $\theta = \phi_- - \phi_+$. Substituting these spatially homogeneous states into Eq. (3), we obtain the dynamical equations for the condensate density ρ_σ , the relative phase θ , and the reservoir

density n_σ :

$$\begin{aligned} \dot{\rho}_\sigma &= (Rn_\sigma - \gamma_c)\rho_\sigma + \sigma 2J\sqrt{\rho_\sigma \rho_{-\sigma}} \sin \theta, \\ \dot{\theta} &= \left(U - \frac{J \cos \theta}{\sqrt{\rho_+ \rho_-}} \right) (\rho_+ - \rho_-) + g_R \delta n_R, \\ \dot{n}_\sigma &= P_\sigma - (\gamma_R + R\rho_\sigma)n_\sigma, \end{aligned} \quad (4)$$

where $U = u_a - u_b$ and $\delta n_R = n_+ - n_-$. The dynamics described by Eq. (4) is rather universal, and would apply to the description of any two-mode polariton system with linear (Josephson-type) coupling such as condensates in a double-well potential [22] or a two-level system [23], provided that their spatial variation can be ignored.

The amplitude and phase of each component in the condensate determine the polarization of the coherent photoluminescence emitted from the cavity. It is therefore helpful to use the pseudospin representation (see, e.g., [18]) which allows us to visualize the polarization of the condensate as a point on the Poincaré (Bloch) sphere. The coordinates of the point are given by the components of the Stokes vector defined in the standard manner:

$$\begin{aligned} S_x &= \sqrt{\rho_\sigma \rho_{-\sigma}} \cos \theta, \quad S_y = \sqrt{\rho_\sigma \rho_{-\sigma}} \sin \theta, \\ S_z &= \frac{1}{2}(\rho_+ - \rho_-), \quad S_0 = \frac{1}{2}(\rho_+ + \rho_-), \end{aligned} \quad (5)$$

with the total condensate density $n_c = \rho_+ + \rho_-$ defining the length of the Stokes vector: $S_0^2 = S_x^2 + S_y^2 + S_z^2 = n_c^2/4$. Note that both the pseudospin and the total condensate density, $S_0 = n_c/2$, are functions of time, therefore it makes sense to consider the evolution of this vector in the normalized pseudospin space, $s_i(t) = S_i(t)/S_0(t)$, i.e., on the Poincaré sphere of unit radius. Although we will use the normalized Poincaré sphere representation throughout the text, one should keep it in mind that S_0 can vary with time.

The stationary states of the system (4) define fixed points on the Poincaré sphere, and the stable ones correspond to phase-locked stationary solutions of Eq. (3) with a well-defined polarization state of the polaritons. Desynchronized states with oscillating densities and relative phase correspond to limit cycles of the dynamical system Eq. (4). From the early studies of the bosonic Josephson junction in an ultracold atomic system [24], it is known that its dynamics can be mapped to that of a superconductor Josephson junction driven by a constant bias current. For the system of Eqs. (4), in the limit of a stationary reservoir, this analogy is valid in a narrow parameter regime identified in [7,25], namely when (a) scattering into the polariton condensate dominates the relaxation dynamics of the reservoir $R\rho_\pm \gg \gamma_R$, (b) the polariton interaction energy is greater than the tunneling energy $Un_c \gg J$, and (c) populations of the two spin components are similar (quasilinear polarization), $n_c \gg S_z$. Under these conditions, following the procedure outlined in [7], one can derive the approximate self-consistent equation for the relative phase:

$$\ddot{\theta} + \gamma_c \dot{\theta} - 2UJn_c^0 \sin \theta = (U + g_R)\delta P; \quad (6)$$

here $\delta P = P_+ - P_-$, $n_c^0 = P/\gamma_c$, and $P = P_+ + P_-$. Analysis of the phase space of this dynamical system [26] shows that, for equal pumping of the two spin components, i.e., linearly polarized pump, $\delta P = 0$, only phase-locked steady states exist. For an elliptically polarized pump when

FIG. 1. (Color online) (a)–(c) Stationary solutions of Eq. (4) for a linearly polarized pump. Branches 1 and 3 are elliptically polarized states; branches 2 and 4 are bonding (s_1) and antibonding (s_2) linearly polarized states, respectively. Only the solid branch is stable (see text). (d) Dynamical evolution of a spatially inhomogeneous state governed by Eq. (3) represented by integrated Stokes parameters (see Sec. IV). The pump is linearly polarized ($\bar{P}_0 = 5$, $P_0^-/P_0 = 0.5$) and has a Gaussian profile with the widths $a_x = a_y = 5$. Red dot: initial state of evolution; green rectangle: final state of evolution corresponding to the marked point in (a)–(c). Inset: evolution of the energy functional (7). Here $J = 0.5$ and other parameters are given in Ref. [21].

$|I_c| = [|\delta P|/P][\gamma_c(U + g_R)/(2UJ)] > 1$, the dynamics is dominated by limit cycles. For $|I_c| < 1$ the system admits steady states determined by the condition $\sin \theta = I_c$ and corresponding to the distinct steady states of the homogeneous system for each value of the pump rate.

Although Eq. (6) is obtained under a very restrictive assumption of a nearly linear polarization of the polaritonic state, it can describe some experimental results such as in Ref. [22]. Notably, when the pump rate is large, conditions (a) and (b) can be automatically fulfilled once condition (c) is satisfied (detailed discussion can be found in Ref. [7]). However, in general, Eq. (6) does not apply, especially when the condensate densities are small and the interaction energy is comparable to Josephson coupling. Therefore, in what follows, we will consider the variety of the steady states and dynamical behavior governed by Eq. (4) under various pumping conditions, which fall outside the regime captured by the Josephson junction model (6).

A. Linearly polarized pump

Most commonly, the optical pump producing a polariton condensate in non-spin-resolved experiments is linearly polarized. In our model, this means that the pump is balanced: $P_+ = P_-$. Under this condition, from the first of Eqs. (4), it follows that the synchronized state can be reached when $Rn_\sigma = \gamma_c$ and the internal Josephson current

FIG. 2. (Color online) Trajectories of the pseudospin vector obtained from numerical solutions of Eq. (4) for a linearly polarized, spatially homogeneous pump, $\bar{P}_+ = \bar{P}_- = 1.2$ and $J = 0.5$; the interaction strength g_R between the condensate and the reservoir takes the same value as in Fig. 1 for (a) and (b); for (c) and (d) an artificially enlarged or reduced g_R is used respectively. Red dots: initial states of evolution; green rectangles: final states of evolution. Initial conditions are arbitrary. Parameters are (a) $g_R = 1.5 \times 10^{-2}$, $\theta(0) = 0.5$; (b) $g_R = 1.5 \times 10^{-2}$, $\theta(0) = 2$; (c) artificially enlarged $g_R = 1.5 \times 10^{-1}$; (d) artificially reduced $g_R = 1.5 \times 10^{-3}$. Other parameters are given in Ref. [21].

$I_J = 2J\sqrt{\rho_+\rho_-}\sin\theta$ becomes zero. Consequently, the phases are locked to $\theta = 0$ or $\theta = \pi \pmod{2\pi}$. The two fixed points on the Poincaré sphere corresponding to these phase-locked linearly polarized states are $s_1 = (1, 0, 0)$ (corresponding to $\theta = 0$), and $s_2 = (-1, 0, 0)$ (corresponding to $\theta = \pi$). The density of the steady state grows linearly with the pump rate: $\rho_\sigma = (P_{th}/\gamma_c)(P_\sigma/P_{th} - 1)$; see Figs. 1(a)–1(c), where the normalized pump rate is defined as $\bar{P}_\sigma = P_\sigma/P_{th}$ and for the linearly polarized pump particularly $\bar{P}_+ = \bar{P}_- = \bar{P}$.

Following the standard Lyapunov stability analysis [26], we deduce that, for our choice of the sign of J , the π out-of-phase state s_2 (the so called *anti-bonding state*) is stable and the in-phase state s_1 is unstable. Direct numerical integration of Eq. (4) with different initial conditions also confirms that s_2 is stable, Figs. 2(a)–2(b), where the interaction constant g_R takes the same value as that in Fig. 1. We note that this selection of the antibonding steady state in a weakly (linearly) coupled two-state system has also been recently confirmed in [16,27]. This effect persists even for the inhomogeneous (Gaussian) pumping [16], as shown in Fig. 1(d) and discussed in Sec. IV below.

The other steady states arise when $\rho_\sigma > (P_{th}/\gamma_c)(P_\sigma/P_{th} - 1)$, where the sign of the inequality is the opposite for the other polarization component; see Figs. 1(a)–1(c), dashed branches 1 and 3. They are sustained by a nonzero internal Josephson current which will grow with the increasing pump rate. Lyapunov analysis reveals that they are unstable.

Introduced by the interaction between the condensate and the reservoir, the blue-shift term which is proportional to g_R plays an important role in the spin dynamics. It can modify the properties of steady states when its energy is

FIG. 3. (Color online) (a)–(c) Selected stationary solutions of Eq. (4) with an elliptically polarized pump ($\bar{P}_+ = 2.4$). Only the solid branches are stable. Red dot (green rectangle): the steady states before (during) addition of a pulsed perturbation. (d) Dynamics of spatially inhomogeneous states under pulsed perturbation represented by the integrated Stokes parameters (see Sec. IV). Steady states are supported by a Gaussian cw pump with $\bar{P}_0 = 4$, $P_0^-/P_0 = 0.4$, $a_x = a_y = 5$, and perturbed by a pulse with $\bar{P}_0 = 0.12$, $P_0^-/P_0 = 1$, $a_{px} = a_{py} = 5$. Red dot: initial state of evolution; green rectangle: final state of evolution. Inset: evolution of the energy functional (7). Here $J = 0.5$ and other parameters are given in Ref. [21].

compatible with the Josephson tunneling energy. Figures 2(c)–2(d) illustrates how different values of g_R can affect the evolution of Stokes parameters in phase space. Here the initial conditions are chosen arbitrarily, and the trajectories are obtained by numerical integration of Eqs. (4). Indeed, both the positions of fixed points and their stability can change, e.g., when $g_R = 1.5 \times 10^{-3}$, both s_2 and s_1 become stable; see Fig. 2(d). Higher values of g_R will lead to the system being in an elliptically polarized state; see Fig. 2(c). The elliptically polarized steady states formed under linearly polarized pumping conditions correspond to the so called *self-trapped states* with strongly imbalanced population of the polarization components. It is a consequence of nonlinear interactions [24,28] and has been studied in the context of a double-well polariton system both experimentally [29] and theoretically [27].

B. Elliptically polarized pump

When the imbalanced polarization-selective pumping is introduced, the dependence of the density of the orthogonal spin components on the pump rate displays a hysteresis-like profile (see Fig. 3). However, it does not correspond to the typical bistability behavior, since most of the blue curve corresponds to oscillatory unstable fixed points. Figures 3(a)–3(c) show a typical fixed point distribution against small pumping imbalance, where all unstable solution branches with

at least one vanishing ρ_σ are not shown for clarity. Stable fixed points exist in the semicircular polarized pumping region (blue solid), which will be discussed in the next subsection, and the semilinearly polarized pumping region (red solid), where at the point $\bar{P}_- = \bar{P}_+$ the fixed point coincides with the single antibonding stable state $\rho_+ = \rho_-$ shown in Fig. 1.

One should note that the above discussion is based on the Lyapunov stability analysis of the steady states, which describes the linear stability to a long-wavelength (spatially homogeneous) perturbation with the wave vector $\mathbf{k} = 0$. The steady state may become unstable to a spatially modulated perturbation with $\mathbf{k} \neq 0$; see, e.g., discussions about the *modulational instability* in 2D [7] and 1D [30] cases.

As a verification of the existence of the stable branches, we performed a pulsed excitation of a steady state in Figs. 3(a)–3(c) where the red dots and green rectangles represent stable steady states before and during the pulse. They fit the inhomogeneous simulation result, Fig. 3(d), qualitatively well (see Sec. IV for details).

C. Circularly polarized pump and spin relaxation

The steady state shown in Figs. 3(a)–3(c) at $\bar{P}_- = 0$ corresponds to the situation when only one spin component is pumped, which is physically akin to the strongly imbalanced dissipative double-well Josephson junction [27] (the opposite case, $\bar{P}_+ = 0$, is physically identical). In this limiting case, the fixed point solutions of Eq. (4) are shown in Fig. 4, where the

FIG. 4. (Color online) (a)–(c) Stationary, spatially homogeneous solutions of Eq. (4) with $\bar{P}_- = 0$. Only the solid branch is stable (see text). (d) Dynamical evolution of a dynamically inhomogeneous state governed by Eq. (3) represented by integrated Stokes parameters (see Sec. IV). The pump is circularly polarized ($\bar{P}_0 = 3$, $P_0^- = 0$), and has a Gaussian profile with the widths $a_x = a_y = 5$. Red dot: initial state of evolution; green rectangle: final state of evolution corresponding to the stable steady state indicated in (a)–(c). Inset: evolution of the energy functional (7). Here $J = 0.5$ and other parameters are given in Ref. [21].

steady state for $\bar{P}_+ = 2.4$ is the same as that represented by the dark blue solid curve in Figs. 3(a)–3(c) at $\bar{P}_- = 0$.

Again, Fig. 4 (d) shows us the inhomogeneous simulation result verifying the position of fixed points predicted by the homogeneous model.

Naturally, the condensate supported by a spin-polarized reservoir created by the incoherent circularly polarized pump retains its nearly circular polarization, with the orthogonal polarization component only weakly populated by the Josephson coupling [Fig. 4(b)]. This dynamical behavior of the mean-field model is consistent with the conclusion that the polariton population transfer between the two spin components is due to the linear polarization splitting [31]. In principle, spin-relaxation mechanisms for the condensed polaritons and reservoir excitons involve indirect scattering processes and momentum relaxation due to phonon scattering, and therefore should be treated within the quantum kinetics formalism [31,32]. Within the framework of our mean-field model, we can account for spin-relaxation processes by introducing phenomenological cross-spin reservoir depletion terms [7] into the model equations (3): $Rn_\sigma \rightarrow R_a n_\sigma + R_b n_{-\sigma}$ and $R|\psi_\sigma|^2 \rightarrow R_a |\psi_\sigma|^2 + R_b |\psi_{-\sigma}|^2$. As long as the natural assumption $R_b \ll R_a$ holds (and in fact, up to $R_b \approx 0.4R_a$ for our parameter values), the number and the type of the steady states under the various pumping conditions considered above remain the same. In particular, the circularly polarized pump will inevitably produce a steady state with nearly circular polarization. However, for larger values of the cross-spin scattering coefficient R_b , the dynamics is dominated by the desynchronised states with oscillating polarization.

IV. INHOMOGENEOUS SPIN-SWITCHING DYNAMICS

In the spatially inhomogeneous situation, it is well known that the in-plane magnetic fields, either being externally applied or caused by the TM-TE splitting [31], can lead to nontrivial spatial structures for both density and spin distributions [7,8,33]. The linear spin coupling can also be regarded as an effective magnetic field pointing along the x axis, which can lead to the appearance of spin patterns. However, here we analyze steady states of the polariton spinor systems sustained by a cw pump with a Gaussian shape and show that the *spatially averaged* evolution of the spin components qualitatively follows the prediction of the homogeneous model discussed in Sec. III. The results obtained within the homogeneous approximation can therefore be used as a guide for spin-switching manipulation with a realistic, spatially inhomogeneous pump.

To determine the evolution of pseudospin under spatially inhomogeneous excitation, we solve directly the open-dissipative GP equations (3) with a radially symmetric Gaussian pump by using the split-step method [34]. The pump rates for each of the polarization components are defined as $P^\sigma(x, y) = P_0^\sigma e^{-(x^2/a_x^2 + y^2/a_y^2)}$, where $a_{x,y}$ are the widths of the common Gaussian pump profile along the x and the y directions, and the total peak rate is $P_0 = P_0^+ + P_0^-$. As before, the peak value of the pump rate is conveniently normalized by the homogeneous threshold value $\bar{P}_0 = P_0/P_{th}$, and δP denotes $P_0^+ - P_0^-$.

By adjusting the pump rate and the degree of its circular polarization $\delta P/P_0$, the system can be driven to a synchronized steady state. In numerical simulations, the condensate mean-field energy functional, $E = E_+ + E_-$, serving as the measure of kinetic and interaction energy of the condensate [35], can be used to detect whether the steady state is reached:

$$E_\sigma = \int d^2\mathbf{r} \left[\frac{1}{2} |\nabla \psi_\sigma|^2 + \frac{u_a}{2} |\psi_\sigma|^4 + g_R n_\sigma |\psi_\sigma|^2 + \frac{u_b}{2} |\psi_\sigma|^2 |\psi_{-\sigma}|^2 + J \text{Re}(\psi_\sigma \psi_{-\sigma}^*) \right], \quad (7)$$

where the integration is over the area where the condensate density is non-negligible.

To map the evolution of the inhomogeneous condensate to that of the Stokes parameters, we trace the evolution of the spatially integrated quantities as $s_i^{\text{int}}(t) = \int S_i(x, y) d^2\mathbf{r} / S_0^{\text{int}}$, where S_0^{int} is given by the spatially integrated length of the Stokes vector. The integration area is the same as that used to determine the steady state energy (7). In the following, unless it is indicated explicitly, we shall refer to the normalized $s_i^{\text{int}}(t)$ simply as integrated Stokes parameters. This normalization procedure, although eliminating information on the spatial distribution of the polarization, allows us to recover the spatially averaged polarization state. This averaged spin dynamics could be observable experimentally even without performing the spatially resolved polarimetry of the cavity photoluminescence.

Polarization dynamics for the linearly and circularly polarized pumps are shown in Figs. 1(d) and 4(d), respectively. As we can see, the spatially averaged dynamics in the case of inhomogeneous pumping fits the homogeneous predictions well.

For a typical steady state under elliptically polarized pumping, the values of the cross section densities and spatial distribution of the pseudospin are shown in Fig. 5(a). Its buildup process is similar to that of the linearly-polarized pumping case; see Figs. 5(b) and 5(c). The linear coupling J causes switching of the dominant density between the two components as polaritons spread out from the pump source, which leads to the appearance of radially symmetric domains of polarization density s_z . The larger J leads to the denser spin pattern and loss of correspondence between the averaged spin dynamics in inhomogeneous and homogeneous cases. If J is sufficiently small, which is the case in the planar semiconductor microcavities, the polarization state is almost homogeneous across most of the integration area. Importantly, and similarly to the case of the miscible atomic condensates [36], the two spin components always coexist in space, so that spatial spin domains do not form.

The degree of control of the averaged polarization state of the condensate attainable by an incoherent, far off-resonant spin-polarized excitation according to our model is presented in Fig. 6 and could be tested in an experiment. Due to our particular choice of the spatial pump profile, in Fig. 6, within the range $|\delta P/P_0| < 0.16$ one needs to increase $\bar{P}_0$ in order to reach steady states. With a larger pump rate, however, the degree of control of polarization decreases because the system will easily fall into a desynchronized state even for a small value of imbalance.

FIG. 5. (Color online) (a)–(c) Evolution towards a synchronized steady state with $\bar{P}_0 = 4$, $P_0^-/P_0 = 0.4$, $a_x = a_y = 5$, and $J = 0.5$: (a) density profiles (along $y = 0$) and spatial distribution of s_z for the steady state, (b) evolution of integrated Stokes parameters and energy functional (7) starting from an arbitrary small initial condition, and (c) evolution of integrated Stokes parameters on the Poincaré sphere from the initial (red dot) to the final (green rectangle) steady state. (d) Evolution of integrated Stokes parameters towards a desynchronized state with $\bar{P}_0 = 8$, $P_0^-/P_0 = 0.4$, $a_x = a_y = 5$, and $J = 0.5$. The initial state is marked by the red dot and the final state is a closed orbit. Inset: evolution of the mean-field energy functional (7) with the integration performed over the area where $n_c(x, y) > 10^{-3}$. Other parameters are given in Ref. [21].

For example, for an elliptically polarized pump, within the region marked with D in Fig. 3, none of the stationary state branches are stable, which does not rule out the existence of the closed orbits (limit cycles) on the Poincaré sphere corresponding to desynchronized states. Evolution of the Stokes vector in a desynchronized state is similar to that shown in Fig. 5(d), where the angle of the trajectory precession can vary depending on the system parameters. This regime is not captured by Eq. (6) because the latter can only describe limit cycles in the vicinity of the $s_z = 0$ plane. Indeed, it

FIG. 6. (Color online) Steady state integrated quantities against the ratio P_0^-/P_0 for $\bar{P}_0 = 4$, $a_x = a_y = 5$, $J = 0.5$. Within the gap $|\delta P/P_0| < 0.16$, one needs to increase $\bar{P}_0$ in order to reach steady states. Other parameters are given in Ref. [21].

can be revealed by inhomogeneous simulations of the full model equations that, in such a state, the spatial distribution of the spin pattern keeps breathing. The corresponding Stokes parameters circle around the Poincaré sphere [37], reducing the degree of (time-averaged) polarization of its luminescence. A similar pinning and depinning effect has been observed experimentally [3]. Here this effect is associated with moving in and out of the phase-locked synchronized regime of the condensate dynamics caused by intrinsic interactions between the two spin components, so that it can arise even without taking into account the structural disorder in the sample.

The coexistence of multiple stable branches described in Sec. III indicates that an inhomogeneous pump might be used to switch the system between different steady states. Indeed, we find that the efficient switching between different spin and density states can be achieved by applying a radially symmetric incoherent *pulsed excitation* of the form $P^\sigma(\mathbf{r}, t) = (1/4)P_0^\sigma \{1 + \tanh[\tau_p(t - t_0)]\} \{1 + \tanh[\tau_p(t_1 - t)]\} e^{-(x^2/a_{px}^2 + y^2/a_{py}^2)}$, where $a_{px, py}$ determines the pulse width, t_0 and t_1 ($t_1 > t_0$) are the pulse switch on and switch off times respectively, and τ_p is a coefficient controlling adiabaticity of the excitation (large τ_p corresponds to a nonadiabatic pulse). We assume that the combined pump intensity is given by an interference of the in-phase cw and pulsed laser beams.

A radially symmetric nonadiabatic pulse, $a_{px} = a_{py}$, causes an abrupt change in the condensate energy and a strong outward density flows combined with the internal Josephson currents leading to spin mixing dynamics [Figs. 7(d)–7(f)]. As a result, after an initial transient dynamics, the system can enter a different steady state [Figs. 7(a) and 7(c)], or a phase desynchronized state [Fig. 7(b)]. If the original steady

FIG. 7. (Color online) Spin switching dynamics for the case of a Gaussian pump. (a)–(c) Energy of the initial and final states after (or during) the switching pulse. Dashed line: pulse switched on. Dash-dotted line: pulse switched off. (d)–(f) Integrated total density and integrated Stokes parameters during the switching process for switching between the steady states shown in (c). Parameters of the cw pump and pulsed perturbation in (c) are the same as in Fig. 3(d); parameters of the pulsed perturbation in (a) are $\bar{P}_0 = 0.4$, $P_0^-/P_0 = 0.65$ and in (b) $\bar{P}_0 = 0.9$, $P_0^-/P_0 = 0.65$. Other parameters are given in Ref. [21].

state is stable, the perturbed system is restored to its initial density and polarization state after the pulse is switched off. We note that the transient dynamics shown in Fig. 7 represents time-dependence of a spatially averaged Stokes vector, and therefore does not necessarily represent a spin wave excited by the pulse [11].

Regardless of the fact that with strong spatial variation the homogeneous results would not be generally applicable, one can still compare the spatially averaged inhomogeneous results of Figs. 7(d)–7(f) with the stationary solutions of Eq. (4), as long as the former begins and ends with steady states. Figure 3 shows two of the relevant fixed point solution branches that are qualitatively comparable to the simulation data. The initial (red dot) and final (green square) steady states correspond to the flipping of s_x shown in Fig. 7(d) before and after the pulse was added. These results confirm that the spin switching dynamics that may be observed in experiments under spatially inhomogeneous excitation conditions can be qualitatively inferred from the structure and stability of the steady states of the corresponding homogeneous model.

V. EXCITATION OF NONTRIVIAL SPIN TEXTURES

While the density and spin wave excitation in the regime of a radially symmetric cw pump and pulse leads to the efficient switching between different spin and density states, the spatial dynamics can be captured by the homogeneous or spatial averaging approximation, as long as the radial symmetry of the spatial density distributions for both spin components is preserved. However, the situation changes dramatically when the radial symmetry of the system is broken. This can be achieved by combining a radially symmetric normal incidence pump with elongated excitation, e.g., at a steep angle to the sample surface, or vice versa. Figure 8 demonstrates the consequence of pulsed perturbation of a steady state established with an elongated ($a_x = 16$, $a_y = 5$) pump. The pulse is a strong, off-resonant, tightly focused Gaussian beam with $a_{px} = a_{py} = 4$ [(a)–(d)] or $a_{px} = a_{py} = 3$ [(e)–(h)] and peak power comparable to that of the cw pump. It was set to having a $\pi/2$ phase difference with the pump and thus no interference pattern appeared. Remarkably, the symmetry of the density flows caused by the pulse is now broken, leading to different flow speeds along the two symmetry axes. As a result, the larger area pulse produces a phase fold associated with a pronounced dip in the condensate density; see Figs. 8(a)–8(c). An even tighter pulse will result in the two phase folds terminating on vortices, leading to formation of a *stable and stationary* configuration of two vortex-antivortex pairs.

Although nontrivial phase structures appear only in one of the spin components [see Figs. 8(d) and 8(h)], they are maintained and stabilized by the internal Josephson current between the spin components and do not exist in the absence of the linear coupling ($J = 0$). Therefore, as has been established for spinor systems of ultracold atoms [38,39], further study of these structures in polariton condensates should be performed in the context of heavily damped spin wave excitation [11]. We also note that no phase singularities appear in the linearly polarized basis, which sets these structures apart from the half-charge and spin-vortices in polaritonic [40–42] and atomic [43,44] condensates.

FIG. 8. (Color online) Excitation of vortex pairs for the case of an elongated elliptically polarized pump with $P_0^-/P_0 = 0.4$ and a tightly focused radially symmetric pulse with $P_0^-/P_0 = 0.65$ and different widths (see text). (a) Density profile of the ψ_+ component showing a density dip associated with a phase fold (b). (c) Cross section of density and phase across the fold. (d) Density profile of the ψ_- component corresponding to (a). (e) Density profile of the ψ_+ component showing stationary configuration of vortex pairs connected by a phase fold (f). (g) Cross-section of density and phase across the phase fold. (h) Density profile of the ψ_- component corresponding to (e). All density plots are logarithmic. Other parameters are given in Ref. [21].

VI. CONCLUSIONS

In conclusion, we have investigated in detail the formation and polarization structure of the steady states of an exciton-polariton condensate created by a far off-resonant *spin-polarized* pump in a semiconductor microcavity with linear polarization splitting. We model the dynamics of the system under the assumption that the incoherent pump creates a fast-responding and rapidly decaying reservoir, which significantly affects the condensate dynamics. Our analysis includes effectively spatially homogeneous pumping conditions, as well as spatially inhomogeneous pumping with a Gaussian profile. The polarization dynamics in the latter case can be mapped to the homogeneous dynamics through

a spatial averaging procedure. We have demonstrated that the phase-locking conditions for the existence of stationary states naturally leads to the self-trapped states with a nontrivial phase relationship between the condensate components, including π out-of-phase (“antibonding”) states. The existence of both phase-locked and spin-beating states can enable efficient control and switching of the polarization states in the system. Experimentally, these findings can be tested with spatially averaged polarization measurements under incoherent spin-selective excitation conditions.

In addition, we have analysed the formation of non-trivial spin-textures under pulsed, spatially inhomogeneous excitation, and demonstrated numerically the reliable creation of stable vortex-antivortex pairs.

Our analysis, with suitable modification of parameters, is widely applicable to a general two-state polariton system with linear coupling, including the weakly linked spatially separated condensates or the multimode condensates in shallow potential traps.

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