The remarkable progresses in the control of matter-light interaction in semiconductor optical microcavities have made it possible to design a new generation of optoelectronic devices [1-7]. These are based on the peculiar properties of exciton-polaritons, half-light half-matter bosonic quasiparticles arising from the strong coupling between photonic cavity modes and excitons in quantum wells placed inside the cavity. One of the most remarkable properties of polaritons is that they have a spin degree of freedom inherited from the photon chirality and exciton spin angular momentum that shows long coherence time and the possibility to be actively manipulated by external fields through the excitonic component [8-9]. This additional feature of polaritons significantly broadens the range of their possible applications to include what is known under the name of spinoptronics. By now, there are already several implemented concepts in the form of spinoptronics devices such as the “Datta and Das” spin transistor [5, 10], the polaritonic analogue of a Berry-phase interferometer [6] and the exciton-polariton spin switch [7, 11]. In all these works a central aspect is the ability to control the spin of polaritons using internal as well as external factors to affect their polarization properties.

From a fundamental point of view, the dominant effect on the polariton spin dynamics is the optical spin Hall effect (OSHE) [12-14]. Basically, it originates from the longitudinal-transverse (LT) splitting of exciton polariton states. The effect of this splitting can be described by an effective magnetic field, strongly dependent on the direction of the quasi-particle propagation and its velocity, giving rise to the polariton spin precession as it propagates in the cavity plane. One of the main problems, in particular when polarization is a key parameter in polariton devices, is the control of such an effective magnetic field which in turn directly affects the polariton state. Usually, the Faraday geometry with the magnetic field directed along the cavity normal is used. In this configuration the studies of the influence of the magnetic field on the polariton dispersion [15], coherence properties [16] as well as on the spin textures in excitonic [17] and polaritonic [18] systems have been reported. However, such geometry cannot be effective on the control of the OSHE [19-22].

By contrast, the effect on exciton-polariton spin dynamics in the Voigt geometry, with the external magnetic field within the plane of the quantum well, has not been studied so far. This is because the in-plane field does not directly couple with the polariton pseudospin and one may expect its effect to be quite minor. In this Letter we demonstrate that this is not the case by reporting for the first time the direct control of the polariton spin transport both in a confined one-dimensional (1D) geometry and in the whole two-dimensional (2D) plane of the cavity through the application of an external magnetic field directed in the cavity plane. In this context, we show the possibility to control and even totally suppress the OSHE for polaritons propagating in a given direction by properly choosing the magnitude and the direction of the applied field. Moreover the application of the external magnetic field causes a stretching of the circular pseudospin patterns in the axis normal to the magnetic field direction and a contraction in the same direction of the external field. From the point of view of applications, the possibility to completely control the OSHE related intensity oscillations allows one to remove the residual density modulations that appear during polaritons propagation and that would be detrimental for the elaboration of the signal inside the devices [10]. All the experimentally observed phenomena can be described within an analytical model taking into account both the field-induced mixing of bright polariton doublet with dark excitonic states and nonlinear effects originating from the polariton-polariton interactions.
The sample studied in this work is a high finesse 3λ/2 GaAs/AlGaAs planar microcavity grown along z || [001] axis with a state-of-the-art polariton lifetime of about 100 ps. The high quality factor $Q > 10^5$ and the low density of defects allow ballistic propagation of polaritons to cover several hundreds of microns, as recently demonstrated by different groups [23, 24]. The microcavity contains 12 GaAs quantum wells (7 nm-wide) placed at three anti-node positions of the electric field, providing a vacuum Rabi splitting of 16 meV. The front (back) mirror consists of 34 (40) pair of AlAs/Al$_{0.2}$Ga$_{0.8}$As layers and cavity-exciton detuning is slightly negative, about $-2$ meV. Polaritons are injected both resonantly and nonresonantly using a low-noise, narrow-linewidth Ti:sapphire laser with stabilized output frequency in a continuous wave operation mode. The sample is kept at a temperature of around 10 K. The sample emission is collected, filtered in polarization and sent to the entrance slit of a spectrometer coupled with a charge coupled device camera.

In order to study the 1D case, we resonantly inject $\sigma^+$ circularly polarised polaritons with a finite momentum into natural misfit dislocations present along the [110] axis of the sample. The 1D confinement is shallow but allows to observe the polariton propagation up to 400 μm from the laser injection point with a negligible spread in the perpendicular direction (see Fig. 1). The coherent spin oscillations during the propagation of polaritons are measured by selectively detecting the emission intensity co- ($\sigma^+$) and cross-polarised ($\sigma^-$) with respect to the exciting laser, as shown in Fig. 1h and Fig. 1i, respectively.
The circular polarization degree, \( P_c \), is then obtained as \( P_c = (I_x + I_y) / (I_x + I_y - I_z) \). For data shown in Fig. 1, polaritons are injected with a speed of 1.5 \( \mu \text{m/ps} \) and the magnetic field, applied in the plane of polariton propagation normally to the propagation direction, spans a range from 0 to 9 T. The pronounced oscillations of the circular polarization degree as a function of coordinate are observed being a signature of the polariton pseudospin precession in the course of propagation. Interestingly and somewhat unexpectedly, the magnetic field affects the frequency of the spatial oscillations in \( P_c \) that increases quadratically with the intensity of the magnetic field, as shown in Fig. 1(g). Figure 2 shows the results of measurement performed with the sample rotated by 90° with respect to the orientation shown in Fig. 1 in order to have the external magnetic field oriented parallel to the dislocation and, hence, to the polariton propagation direction. A deeper dislocation is considered in this case, showing a non-trivial polarization pattern along propagation. Indeed, increasing the magnetic field intensity, the spatial oscillations frequency first decreases (up to 5 T) and then increases with a constant positive offset in the co-polarised component.

The observed effects can be quantitatively described within a pseudospin model parametrizing the polariton spin density matrix through the vector \( S_k \) whose \( z \) component describes the circular polarization degree of the particles and the in-plane components characterize the linear polarization degree in two sets of axes. The equation of motion for the pseudospin of polaritons in the \( k \) state is written as [12]

\[
\frac{\partial S_k}{\partial t} + S_k \times \Omega_k + \gamma_s S_k = 0, \tag{1}
\]

Hereafter we assume ballistic propagation of polaritons, use the set of axes with \( x \parallel [110], y \parallel [\overline{1}10] \) and \( z \parallel [001], \Omega_k \) is the effective pseudospin precession frequency and the last term accounts for the spin relaxation processes with the rate \( \gamma_s \). The effective precession frequency components read

\[
\Omega_{k,x} = \frac{\Delta_{LT}(k_x^2 - k_y^2)}{k^2} - \beta(B_x^2 - B_y^2), \tag{2a}
\]

\[
\Omega_{k,y} = 2\Delta_{LT}k_x k_y / k^2 - 2\beta B_x B_y, \tag{2b}
\]

\[
\Omega_{k,z} = \alpha S_z, \tag{2c}
\]

and contain contributions from the LT splitting of polariton modes, \( \Delta_{LT} \), in the linear polarization components due to the applied external magnetic field in the cavity plane \( \beta B_x^2, B_y \), and the so-called self-induced Lamb–Morr precession of the polariton pseudospin due to the polariton-polariton interactions (\( \alpha \)) treated here within the mean-field approach [8, 26, 27]. Importantly, the in-plane components of the effective field contain the quadratic magnetic field contributions. The form of these terms follows from the symmetry arguments since the quadratic combinations \( B_i B_j \) and \( k_i k_j \) with \( i, j = x, y \) transform in the same way already in the isotropic approximation and the effects of \( C_{2n} \) point symmetry of the studied structure on the magnetic-field induced terms are disregarded. Microscopically, the parameter \( \beta \) results from magneto-induced mixing of polariton states and dark (spin-forbidden) excitons with the additional contribution from the diamagnetic effect, just like for excitons in quantum wells and quantum dots [28, 29], see Supplementary material [30] for the details.

Equation (1) can be solved analytically for various experimentally relevant configurations. Let us first set \( \alpha = 0 \), i.e., neglect the effect of the polariton-polariton interactions and assume that the polaritons propagate along the \( y \)-axis and \( B \) is either parallel or perpendicular to the polariton velocity. From Eqs. (2) it follows that \( \Omega_y \equiv 0 \) and

\[
S_z(y) = S_{z0} \cos(\kappa y) e^{-\nu/\ell_z}. \tag{3}
\]

Here the propagation constant of the oscillatory distribution of pseudospin is:

\[
\kappa = \frac{|\Delta_{LT} + \beta(B_x^2 - B_y^2)/}v, \tag{4}
\]

and the pseudospin decay length \( \ell_z = v/\gamma_s \), where \( v = k/m \) (with \( m \) being the polariton effective mass) is the polariton velocity. For positive \( \beta \) and \( B \perp k \), Eq. (1) predicts a monotonic increase of the propagation constant \( \kappa \) with the square of the magnetic field magnitude \( B^2 \), in full agreement with the experimental results shown in Fig. 1(g). By contrast, with \( B \parallel k \) the dependence \( \kappa(B) \) is more complicated: from Eq. (4), it follows that \( \kappa \) is a nonmonotonic function of \( B \). First it decreases for \( B < B_c = \sqrt{\Delta_{LT}/\beta} \) and then increases for higher field intensities. Under the condition \( B = B_c \) the complete stop of oscillations is expected due to the suppression of the LT-splitting by the magnetic field similarly to the field-induces suppression of exciton anisotropic splitting in quantum dots [28, 29]. Qualitatively, these results are in agreement with the experimental data presented in Fig. 2 with \( B_c \approx 4 \text{T} \).

While Eqs. (3) and (4) quantitatively fit the data in Fig. 1, see solid lines in the panels (b,c), the linear model is not sufficient to describe all the peculiarities of the polariton polarization dynamics shown in Fig. 2. This is because the nonlinearity due to spin-dependent polariton-polariton interactions becomes of particular importance in the situation of \( B \approx B_c \). Also, in the experimental geometry with \( B \parallel k \), the polariton propagation velocity is relatively small, \( v = 0.3 \mu \text{m/ps} \), that results in the weaker manifestation of the LT-splitting effect. Indeed, one of the peculiarities seen in Fig. 2 is the positive offset in the \( S_z(y) \)-dependence. This effect is the manifestation of the presence of the third component of the effective field \( \Omega_z \propto S_z \), which tilts the pseudospin precession axis towards the \( z \)-axis and suppresses partially the effect of the LT-splitting [18, 31]. The results of calculations for this configuration are shown by blue curves.
FIG. 3. Measured circular polarization degree patterns with $\mathbf{B} \parallel y$ for the magnetic field intensities of (0, 7, 9) T and the initial polarizations $\sigma^+$ (a,b,c) and $\sigma^-$ (d,e,f). The propagation velocity is about 1.8 $\mu$m/ps. (g) Cross section of $S_z$ along the vertical directions at 0 T (red line) and 9 T (blue line) as indicated by the dashed line in (a) and (c), respectively. The fitting function is $Ae^{-b\gamma} \sin(ky + \phi)$. (h) Same as in (g) but for $\sigma^-$ (see (d) and (f)). (i) Propagation constant $\kappa$ extracted as best fit to the data for $\sigma^+$ (red) and $\sigma^-$ (blue) polarization. (l)-(q) The circular polarization degree in real space simulation based on Eq. (1). The magnetic field intensities correspond to ones in (a-f). The magnetic field intensities and the dynamics of $S_z(y)$ return to the harmonic character at the distance of about 300 $\mu$m, Fig. 2.

In order to understand how the complete two dimensional polarization spatial distribution is affected by the application of an in-plane magnetic field, we performed a different experiment using a free, high velocity, radially expanding polariton condensate in a clean region of the sample. Even though obtained with a nonresonant circularly polarized pumping scheme, polaritons, in our case, inherit about 40% of circular polarization from the pump [13]. With the aim of efficiently injecting polaritons, the energy of the pump is tuned at the first minimum of the reflection stop band making it possible to reach the condensation density threshold in a region within the laser spot, blueshifted of about 4 meV from the bottom of the lower polariton branch. From this central region, a polariton flow is ballistically expelled and is free to radially propagate in the plane of the cavity outside the excitation spot region, with an acceleration due to the gradient in the potential resulting from the blueshift under the excitation spot [32]. This configuration enables polaritons to ballistically propagate with a speed of about 2 $\mu$m/ps.

Figure 3(a–f) shows the $P_c$ for different magnetic field intensities. A small asymmetry in the distribution function already present at $B = 0$ T is initially compensated, and then enhanced by the external magnetic field, see Fig. 3(a–c). Indeed, the pattern was at $B = 0$ T elliptical with the horizontal axis greater than the vertical but, by increasing the applied magnetic field, these axes become inverted and the ellipse gets rotated. This is confirmed by the opposite behaviour of the spatial frequency with respect to the vertical and horizontal axes as reported in the inset of Fig. 3. By changing the polarisation degree of the exciting laser from right-circular to left-circular, we change the relative orientation of the polariton spin with respect to the magnetic field. In Fig. 3(d–f), we show that in this case the long eccentricity axis at 9 T is oriented roughly perpendicularly to Fig. 3(c).

The resulting circular polarization degree of polaritons during the expansion in the two-dimensional plane can be qualitatively described by the same equation (1) applied to all $k$ on the elastic circle of the radius $k$. Figure 3(i–q) illustrates the 2D expansion of the $P_c$ theoretically predicted from the model. The values of the effective magnetic field magnitude $B$ in this figure correspond to those in Fig. 3(a–f). In order to distinguish the influence of different effects on the shape of the polarization pattern it is possible to consider the following. At the pure circular initial polarisaton, in the presence of only the LT-splitting, the circular polarization degree patterns in real space are rotationally invariant. A slight tilt and the squeezing of the patterns in $x$ direction even in the absence of the external magnetic field are due to the fact that the initial polarization is not pure circular but it contains a small admixture of the linear components, see caption to Fig. 3. Moreover, the external magnetic field tends to change the spatial frequency $\kappa$ with, as shown above, a different effect depending on the relative direction of the external field and the propagation. According to Eq. (4), the absolute value of the propagation constant $\kappa(k, y, B)$ decreases with increasing $B$ (directed along the $y$-axis) until the magnetic field compensates the effect of the LT splitting while a further increase in $B$ leads to the increase of the spatial frequency. In the propaga-
tion direction orthogonal to B, the propagation constant $k_\perp = k_x + i k_y$ monotonically increases with the increase of $B$ on all extent.

In conclusions, we have experimentally demonstrated the control of the optical spin Hall effect by tuning an external magnetic field applied in the direction of propagation of polaritons. This can be useful to avoid unwanted rotation of the polarization in polariton devices or by controlling the spin degree at a given position. In fact, if the spin precession is instead required [13], the spin beat frequency can be tuned by the external magnetic field applied perpendicularly to the propagation direction of polaritons. In the 2D expansion the in-plane magnetic field induces an additional polarization anisotropy in the structure that manifests itself as a deformation of the pseudospin patterns in real space. We have developed a theoretical model that qualitatively explains the observed effects.

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