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One-Way Reflection-Free Exciton-Polariton Spin-Filtering Channel

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We consider theoretically exciton-polaritons in a strip of honeycomb lattice with zigzag edges and it is shown that the interplay among the spin-orbit coupling, Zeeman splitting, and an on-site detuning between sublattices can give rise to a band structure where all the edge states of the system split in energy. Within an energy interval, one of the spin-polarized edge states resides with the gapless bulk that has opposite spin. Being surrounded by opposite spin, and the absence of the backward-propagating edge state, ensures both reflection-free and feedback-suppressed one-way flow of polaritons with one particular spin in the system. The edge states in this system are more localized than those in the standard topological-polariton systems and are fully spin polarized. This paves the way for feedback-free spin-selective polariton channels for transferring information in polariton networks.

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I. INTRODUCTION

It was recognized long ago that one of the crucial ingredients for optical circuits is the isolation of input and output [1]. Equivalently, circuits should operate in a one-way fashion, where feedback in the backward direction is suppressed. As Maxwell's equations are time-reversal invariant, the realization of such a requirement is not obvious in many optical systems and reflections must be considered carefully. With the emergence of topological photonics [2,3], special techniques have been introduced to artificially break time-reversal symmetry (e.g., helical-waveguide arrays as in Floquet topological insulators [4]) or exotic material solutions employed at particular wavelengths (e.g., magnetically sensitive gyrotropic materials [5]). One of the well-known features of topological photonic lattices is the appearance of chiral edge states. As these propagate in only a certain direction and do not undergo back scattering, it has been natural to expect that such states are relevant for the robust propagation of optical signals.

Instead of breaking time-reversal symmetry directly in a photonic system, an alternative is to couple it to a system in which time-reversal symmetry is broken. In particular, when photons are coupled to excitons in semiconductor micropillars, they hybridize to form exciton-polaritons, which inherit a Zeeman splitting due to their excitonic component when placed in a magnetic field [6]. Since such micropillars can be arranged into lattices, it has

been possible to form topological-polariton systems [7], accounting for the effect of spin-orbit coupling [8–10].

Exciton-polaritons also exhibit a significant nonlinearity, which in theory has been predicted to itself induce topological behavior [11], antichiral behavior [12], and topological solitons [13]. Separate from topological physics, nonlinearity has been shown to provide low-energy polaritonic switches [14–16] and to realize other information-processing elements such as transistors [16–19], routers [20–22], and amplifiers [23,24]. The possible use of topological-polariton states at the edges of a strip to couple these elements suffers, however, from an obvious drawback: such states come in pairs, which propagate in opposite directions on the opposite edges of the strip [9,10]. Consequently, there will be an unwanted counterpropagating signal at the opposite edge of the strip, leading to potentially detrimental feedback (unless the strip is made very wide, which would lead to other problems in scalability).

Furthermore, it has long been expected that polaritonic devices would make use of the spin degree of freedom [25], which can take one of two spin projections on the structure growth axis (denoted $\sigma_{\pm}$). Individual spin switches [26] and gates [27] have encoded information in this spin. Unfortunately, though, topological-polariton states have a specific polarization that mixes the σ_{+} and σ_{-} components [9,10], which would not allow the preservation of individual σ_{+} or σ_{-} spin-polarized wave packets.

In this work, we aim to develop the recent advances in topological-polariton physics to introduce a channel for connecting polaritons of a given spin. We aim for the channel to be reasonably thin and allow only one

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FIG. 1. (a) A schematic diagram of coupled polariton micropillars arranged in a honeycomb strip, where two different sublattices are indicated by two colors. The considered system can act as a spin filter by allowing the polaritons with one of the spins in the system to propagate robustly from right to left, while propagation in the opposite direction is restricted. The red arrow indicates the perpendicularly applied magnetic field (H_z). (b) The potential application of the polariton spin filter (SF) as a channel that can be used to couple two polariton networks without feedback (left). The same cannot be done using a polariton topological insulator (TI), where there exists a possibility of feedback via an opposite edge state (right).

spin to propagate in a specific direction. To meet this aim, we consider a strip of honeycomb polariton lattice, which is composed of two interlocking sublattices, A and B [see Fig. 1(a)]. The sites of the different sublattices are engineered with different on-site energies. Accounting for a magnetic field and spin-orbit coupling, we find that the states at different edges are split in energy and become almost completely spin polarized. Furthermore, an edge state of a given spin can be embedded in an energy range where the only existing states have opposite spin polarization. As the scattering of polaritons with disorder tends to be spin conserving, this maintains a unidirectional propagating edge state for a specific spin, which is not possible with a regular topological-polariton system [see Fig. 1(b)]. A further key difference with existing topological-polariton realizations is that the strip can be made remarkably thin (three lattice constants thick) and still maintain its unidirectional spin-polarized behavior.

II. SCHEME

We start by considering a planar semiconductor microcavity containing exciton-polaritons subjected to a resonant excitation, expressed by the driven dissipative

Schrödinger equation:

$$i\hbar \frac{\partial \psi_{\sigma_{\pm}}(x, y)}{\partial t} = \left[-\frac{\hbar^2 \nabla^2}{2m} + V(x, y) - i\frac{\Gamma}{2} \pm \Delta_z \right] \psi_{\sigma_{\pm}}(x, y) + \Delta_T \left(i\frac{\partial}{\partial x} \pm \frac{\partial}{\partial y} \right)^2 \psi_{\sigma_{\mp}}(x, y) + F_{\sigma_{\pm}}(x, y) e^{i(k_0 x - \omega_p t)}. \quad (1)$$

Here, $\psi_{\sigma_{\pm}}(x, y)$ are the wave functions for $\sigma_{\pm}$ spin-polarized polaritons; $V(x, y)$ is a potential describing a strip of honeycomb lattice having zigzag edges along the y direction; Γ is a uniform decay throughout the sample; and $2\Delta_z$ and Δ_T are the Zeeman and transverse-electric–transverse-magnetic splittings, respectively. $F_{\sigma_{\pm}}(x, y)$ are the circular-polarization components of a resonant excitation with frequency ω_p . Each micropillar is modeled by placing a Gaussian of width σ at each site [28] and the lattice has a periodicity a along the x direction, such that

$$V(x, y) = V(x + a, y) = -(v_o \pm \delta) \sum_{x_n, y_n} \exp \left[-\frac{(x - x_n)^2 + (y - y_n)^2}{\sigma^2} \right]. \quad (2)$$

Here, the summation runs over all sites, $(v_o + \delta)$ is the potential depth corresponding to the A sublattices, and $(v_o - \delta)$ is the potential depth corresponding to the B sublattices. In principle, the energy detuning, δ , between the A and B sublattices can be induced by locally adjusting the thickness of the cavity or by nonresonantly pumping one of the sublattices, where the excitonic reservoir will induce a blue shift locally [29]. Alternatively, micropillars corresponding to the A and B sublattices could be engineered with different diameters [30]. In order to determine the characteristics of the system, we consider the linear band structure by setting $F_{\pm} = 0$ and $\Gamma = 0$ in Eq. (1) and, after applying the Bloch theorem, the wave function can be written as $\psi_{\sigma_{\pm}}(x, y) = u_{\sigma_{\pm}}(x, y) \exp(ikx)$. Here, $u_{\sigma_{\pm}}(x, y)$ are the Bloch wave functions, which are periodic along the x direction, and k is the wave vector in the x direction. Upon substitution, Eq. (1) leads to an eigenvalue problem, which can be solved to obtain the band structure. In Figs. 2(a) and 2(b), and Figs. 2(c) and 2(d), we plot the dispersion without and with the on-site term, respectively. For $\delta = 0$, we obtain the band structure corresponding to the polariton topological insulator (TI) [see Fig. 2(a)], where the time-reversal symmetry is broken by the Zeeman splitting. Here, the bulk modes are gapped and in the band gap, robust edge states appear. In this case, the edge states come in counterpropagating pairs, meaning that within the band gap of each propagating edge state there is another one propagating in the opposite direction. The inclusion of δ

FIG. 2. The band structure of the system under consideration of on-site energy $\delta = 0$ in (a),(b) and $\delta \neq 0$ in (c),(d). In (a) and (c), the states residing in bulk are shown in blue, whereas red and green correspond to the states located near the two edges. In (b) and (d), the same dispersions with spin contribution are plotted, where red corresponds to the states having spin σ_+ and green corresponds to the states having spin σ_- . The dashed lines in (c) and (d) represent the energy window where the edge state is surrounded by the bulk states with opposite spins and also there is no counterpropagating edge state. (e) A comparison of the localization of the edge states between the polariton SF and polariton TI.

breaks the inversion symmetry and the parameters can be chosen such that the band gap at one of the Dirac points gets closed, resulting in a topologically trivial system. The dispersion of the bulk modes is similar to those in bilayer graphene, with a perpendicular bias [31], although in this case the edge states are completely separated in energy [see Fig. 2(b)]. Note that there is no counterpropagating edge

mode present in the energy window of interest. We also plot the localization of the edge states in Fig. 2(e), which shows that the edge states in the spin-filter (SF) regime are more localized than the standard polariton TI.

To emphasize the interest in the obtained states, we plot the band structures and color code them depending upon the contribution of the spins. In the case of a polariton TI [see Fig. 2(b)], we find that within the band gap, the edge states have mixed spin. On the other hand, in the SF regime [Fig. 2(d)] the edge states have almost pure spin and, more importantly, the topmost edge state is surrounded by only the bulk modes with opposite spin polarization. We can then expect σ_+ spin-polarized polaritons to propagate through the edge mode without back scattering, manifesting a robust SF.

Considering typical parameters, we have set the following: $a = 3 \mu\text{m}$; $m = 2.5 \times 10^{-5} m_e$, where m_e is the free electron mass; $v_o = 16.5 \text{ meV}$; $\delta = 0.54 \text{ meV}$; $\sigma = 0.45 \mu\text{m}$; $\Delta_T = 0.15 \text{ meV } \mu\text{m}^2$; and $\Delta_Z = 0.18 \text{ meV}$. This gives an energy interval of the SF edge state of around 0.12 meV , which is well above the line width of modern samples. Generally, the SF regime can be found for $\varepsilon_g \simeq 2\Delta_Z \leq \Delta_T$, where ε_g is the trivial band gap induced by the broken inversion symmetry [30] and Δ_T determines the group velocity of the edge state. It should be stressed that instead of the magnetic field (Zeeman splitting), the non-resonant excitation can itself break the symmetry between the two spins [32,33]. For the rest of the paper, we focus on the topmost edge state and although here only σ_+ spin-polarized polaritons are filtered, the same can be done for the σ_- spin-polarized polaritons [30].

III. SPIN FILTERING

To see the effect of spin filtering on a wave packet, we excite the topmost edge state resonantly with a continuous linearly polarized Gaussian-shaped pump by setting $F_{\sigma_+} = F_{\sigma_-}$. The spin-polarization degree is defined as

$$S(x,y) = \frac{|\psi_{\sigma_+}(x,y)|^2 - |\psi_{\sigma_-}(x,y)|^2}{|\psi_{\sigma_+}(x,y)|^2 + |\psi_{\sigma_-}(x,y)|^2}. \quad (3)$$

The spatial distribution of $S(x,y)$ is plotted in Fig. 3(a) in the SF regime, which shows that although we have excited both spin states, only σ_+ spin-polarized polaritons populate the top edge and propagate over $100 \mu\text{m}$ from the spot of excitation. Due to the finite lifetime, the intensity decreases with an increase of the propagation distance. σ_- spin-polarized polaritons are instead scattered randomly into the bulk. To demonstrate that there is no state of σ_+ spin-polarized polaritons to propagate in the reverse direction, we consider injecting both spins from the opposite corner of the top edge [see Fig. 3(c)]. We observe that the σ_+ spin-polarized polaritons cannot propagate from the left to the right of the system; instead, they are converted to

FIG. 3. The spatial distribution of the spin degree defined in Eq. (3) for the polariton SF is shown in (a), where polaritons having both spins are pumped but only σ_+ spin-polarized polaritons flow one way. The same for polariton TI is shown in (b), where polaritons with mixed spins flow counterclockwise through the edges of the sample. (c) The feedback-suppression effect, where the σ_+ spin-polarized polaritons cannot flow from left to right; instead, they are converted to σ_- spin-polarized polaritons and scatter randomly into the bulk. The red arrows indicate the position of the continuous pump. Parameters: pump width = $9 \mu\text{m}$; $\Gamma = 0.01 \text{ meV}$; $k_0 = 2.5/a$ and $\omega_p = -1.65 \text{ meV}/\hbar$ are used in (a) and (c); and $k_0 = 2.35/a$ and $\omega_p = -1.98 \text{ meV}/\hbar$ are used in (b).

σ_- polaritons and scatter into the bulk. This confirms that the system can act as a feedback-free unidirectional spin filter. On the other hand, in a polariton TI, the polaritons with mixed spins rotate through the edges of the sample in the anticlockwise direction [see Fig. 3(b)]. Although the exact contribution of the spins in $S(x, y)$ can be changed slightly depending upon the state that is being excited, there is always a significant contribution from both the spins.

In all our calculations, we take a lifetime around 60 ps, close to that reported in Refs. [7,29]. The SF would also work with a shorter lifetime, although the propagation distance would decrease [30]. Higher lifetimes such as those reported in Ref. [34] are thus the ideal.

FIG. 4. (a)–(d) The propagation of a σ_+ coherent pulse in a honeycomb lattice strip with one of the sites removed. (e) The propagation of the same pulse in a honeycomb lattice strip with spatial disorder. The color scale represents the σ_+ polarized intensity, integrated over the strip width (the y coordinate). The green curve indicates the average position of the wave packet and the dashed lines the variance. These are defined as $x_m = \int x |\psi_{\sigma_+}|^2 dx dy / I_0$ and $x_v^2 = \int x^2 |\psi_{\sigma_+}|^2 dx dy / I_0 - x_m^2$, respectively, where $I_0 = \int |\psi_{\sigma_+}|^2 dx dy$. The parameters are the same as in Fig. 3(a).

IV. PULSE PROPAGATION AROUND DEFECT

Figures 4(a)–4(d) consider the propagation of a σ_+ polarized pulse (obtained by multiplying the coherent excitation with a Gaussian in time) in the case where a micropillar is removed from the top edge. We observe that near the defect, the pulse spreads a little but does not back scatter and after some time the pulse passes the defect. This further proves the robustness of the edge state.

V. PROPAGATION IN THE PRESENCE OF DISORDER

In real systems, disorder is present throughout. Figure 4(e) shows the propagation of a σ_+ polarized pulse in the presence of a continuous disorder potential, which is characterized by a root-mean-square amplitude of $30 \mu\text{eV}$ and a correlation length of $3 \mu\text{m}$ (corresponding to typical experimental values [35]). Again, no back scattering of the σ_+ polarized wave packet is observed.

VI. STRIP WIDTH

In addition to supporting one-way propagation, a channel for connecting information-processing elements should not be too wide, so as not to limit the footprint area of future systems. In Figs. 5(a) and 5(b), the band structure for a narrow strip having three unit cells along the y direction

FIG. 5. The band structure corresponding to a very short strip consisting of three unit cells along the y direction for TI shown in (a) and the same for SF shown in (b). The edge states in the topological gap corresponding to the TI disappear. (c) The spatial distribution of the spin degree for SF. The red arrow indicates the position of the continuous pump. The parameters are kept the same as those in Fig. 3(a).

is plotted for a polariton TI and SF, respectively. Due to the lack of bulk states, the well-spread edge states in the polariton TI [see Fig. 2(e)] vanish, while comparatively more localized edge states in the polariton SF remain unaffected, resulting in σ_+ spin propagation as shown in Fig. 5(c).

VII. DISCUSSION

Modern electronic devices and integrated circuits are limited by heating and time delay of communication between components. Optical devices seem to be promising candidates for fast interconnections and higher fidelity may reduce error correction overhead, resulting in less heating [36]. However, the optics needs to operate at a lower threshold to be relevant, for which nonlinear optical systems are essential. Along with the nonlinearity, the spin degree of freedom of exciton-polaritons makes them extremely promising for nonlinear spintronics devices. The coupling between the exciton-polariton condensates in lattices has been demonstrated but this coupling is always bidirectional and, consequently, there will be unwanted feedback. Although the proposed spin filter is not itself a device, it could become an essential component of complete future circuits or networks.

VIII. CONCLUSION

We propose theoretically a scheme to obtain an optical spin filter based on exciton-polaritons. The naturally present spin-orbit interaction along with Zeeman splitting and an on-site detuning between sublattices in a polariton honeycomb lattice gives rise to a band structure in which all the edge states of the system split in energy and there exists an energy window in which one of the edge states

with pure spin is surrounded by bulk states with opposite spin. Since there is no state with the same spin to which to back scatter, polaritons having one of the spins in the system can propagate robustly through the edge state and thus spin filtering is achieved. The absence of the pairing backward-propagating edge state reduces feedback in the system. Furthermore, the edge state in the polariton SF is more localized than those in a polariton TI, which allows us to operate with very narrow channels. In contrast, regular polariton TI systems would require much wider strips and/or lattices to support robust propagation and would cause mixing between spin-polarized states.

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