Split-off states in tunnel-coupled semiconductor heterostructures for ultrafast modulation of spin and optical polarization

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We present a theoretical analysis of the split-off states emerging due to a tunnel coupling between a remote bound state and a semiconductor quantum well (QW). The on-site Coulomb repulsion and the spin splitting of the bound state have been considered. The split-off states emerge in the band gap of the QW and reveal themselves as two solitary peaks in the photoluminescence (PL) from the QW. The peaks have opposite circular polarization and their spectral position strongly depends on the tunnel coupling strength. We suggest a mechanism of ultrafast PL polarization switching by means of electrical modulation of the tunnel coupling by an external gate. The obtained results open a new possibility for the spin and optical polarization control in nanoscale systems.

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I. INTRODUCTION

Application of electron spin in spintronic devices and for data storage requires formation and precise manipulation of spin polarization [1–10]. A common and effective way to probe spin polarization is to detect circular light polarization of luminescence. However, it has been realized rather recently that this spin to light coupling has another advantage useful for practical applications. While the characteristic time of a radiative recombination in a semiconductor usually lies in the range of nanoseconds, the polarization of the light can be modulated much faster and at a low energy cost. The idea underlies the ongrowing interest in semiconductor spin lasers [11,12], which are considered as a promising concept to boost optical data transmission by ultrafast spin and optical polarization dynamics [12,13]. Thus, efficient and fast mechanisms of creation and control of the electron spin polarization in semiconductor heterostructures are of a great importance not only as an essential part of semiconductor spintronics [14,15], but, in particular, for increasing the speed and efficiency of optical information processing.

Significant progress has been made in studies of stationary spin-polarized carriers transport in tunnel junctions in the presence of spin-orbit and exchange interaction [16,17]. Spin-polarized current sources that use nonmagnetic materials are particulary attractive as they avoid the presence of stray magnetic fields that may cause undesirable effects on the spin transport. For example, spin-filter devices capable of generating spin-polarized current without using magnetic properties of materials were proposed in Refs. [18,19].

A notably different mechanism of the spin injection is possible in heterostructures with a QW and remote bound states at magnetic ions separated by a thin spacer [20–23]. In such hybrid heterostructures the magnetic properties and spin polarization of the carriers in the QW can be controlled via the spacer thickness and shape [24–26]. Such systems reveal the mechanism of dynamic spin polarization of electrons in the QW due to the spin-dependent tunneling between the semiconductor QW and the spin-split bound state, which was demonstrated experimentally [27,28] and analyzed theoretically [7,8,29]. A control of the polarization sign by pumping the QW with a laser pulse matching the one of the bound state spin sublevels has been suggested [8].

While the tunnel hybridization of a QW with spin-split bound state gives rise to nonstationary spin and optical polarization, there is another stationary effect related to the spin-dependent hybridization which is very different from the nonstationary effects discussed in Refs. [7,8,29]. Along with the modification of the 2D continuum spectra the hybridization leads to formation of the split-off states in the band gap of the QW. A role of the split-off states in the photoluminescence from the QW and their potential impact on the spin polarization in the hybrid structures has not been considered so far.

In this paper we analyze theoretically the emergence of circular polarized PL in the QW due to the formation of split-off spin polarized localized states in the QW band gap. These states are formed due to the interaction between the 2D electrons in the QW and localized electrons in the spin-split energy levels of the remote bound state. Our results show the possibility to obtain fully circularly polarized PL signal with the split-off states and to control the sign of electrons spin polarization in the QW.

II. THEORETICAL MODEL

Let us consider a model system formed by a QW with one size quantization subband coupled to a spin-split bound state with an energy ε_0 separated from the QW by a tunnel barrier. QW is optically or electrically pumped with unpolarized nonequilibrium 2D electrons with the energies ε_k, where k is the in-plane wave vector. The model system shown in

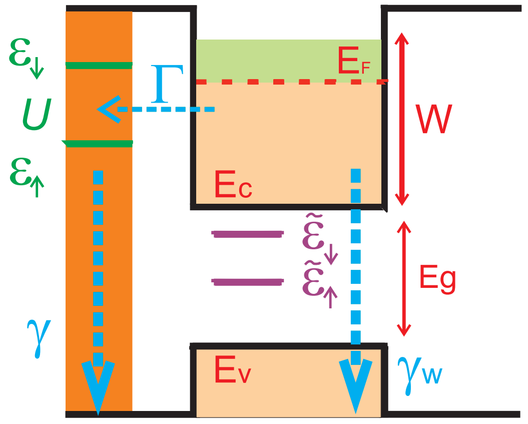

FIG. 1. Scheme of the hybrid system semiconductor QW with remote spin-split bound state. The split-off states are shown in the band gap of the QW. The energy scale is schematic, in semiconductor heterostructures $W \gg E_F$.

Fig. 1 corresponds to the experimentally studied (Ga,Mn)As heterostructures [27], in which the bound state is formed by Mn ions in interstitial positions providing donorlike states. The remote Mn doping layer is located at a distance of 3–10 nm from the QW. The spin splitting of the donorlike bound state occurs due to the effective exchange field in the Mn layer. There is also a fast relaxation from the dopant donor levels due to nonradiative recombination with the holes in the low-temperature grown (Ga,Mn)As layer [6,29].

The Hamiltonian of the system can be written in the following form:

$$\hat{H} = \hat{H}_{\mathrm{QW}} + \hat{H}_{\mathrm{imp}} + \hat{H}_T, \tag{1}$$

where $\hat{H}_{\mathrm{QW}}$ describes nonequilibrium 2D electrons in the QW:

$$\hat{H}_{\mathrm{QW}} = \sum_{k\sigma} \varepsilon_k \hat{c}^+_{k\sigma}\hat{c}_{k\sigma}, \tag{2}$$

$\hat{H}_{\mathrm{imp}}$ corresponds to the bound state with the on-site Coulomb repulsion given by the Hubbard term:

$$\hat{H}_{\mathrm{imp}} = \sum_{\sigma} \varepsilon_0 \hat{n}_\sigma + U\hat{n}_\sigma\hat{n}_{-\sigma}, \tag{3}$$

and $\hat{H}_T$ is the tunneling part describing coupling between the QW and the bound state:

$$\hat{H}_T = \tau \sum_{k\sigma} (\hat{c}^+_{k\sigma}\hat{c}_\sigma + \hat{c}^+_\sigma \hat{c}_{k\sigma}). \tag{4}$$

Here index k labels continuous spectrum states in the QW, and τ is the tunneling transfer amplitude between the QW states and the bound state. We consider τ to be independent on the momentum and spin. The bound state energy level ε_0 can be split due to an exchange interaction or an external magnetic field into two spin sublevels with the energies $\varepsilon_\sigma = \varepsilon_0 \mp \Delta_0/2$, where $\sigma = \uparrow, \downarrow$ denotes the electron spin projection and Δ_0 is the spin splitting energy. Operator $\hat{c}^+_{k\sigma}/\hat{c}_{k\sigma}$ creates/annihilates electrons in the QW, $\hat{n}_\sigma = \hat{c}^+_\sigma\hat{c}_\sigma$ is the occupation number operator for the bound state, operator $\hat{c}_\sigma$ destroys the electron at the bound state with the spin projection σ. U is the on-site Coulomb repulsion energy for the doubly occupied bound state.

Self-consistent description of the tunneling processes in the considered system could be made using the nonequilibrium Keldysh diagram technique. Let us introduce the operators:

$$\hat{G}^R_\sigma(t,t') = -i\Theta(t-t')[\hat{c}_\sigma(t)\hat{c}^+_\sigma(t') + \hat{c}^+_\sigma(t')\hat{c}_\sigma(t)],$$
$$\hat{G}^A_\sigma(t,t') = \hat{G}^{R+}_\sigma(t',t). \tag{5}$$

Retarded $G^R_\sigma(t,t')$ and advanced $G^A_\sigma(t,t')$ Green's functions are obtained from (5) using the standard averaging procedure [30,31]. Performing Fourier transformation over (5) and following the standard nonequilibrium Keldysh Green's functions formalism, one can derive the closed system of Dyson equations, which determines the retarded bound state Green's function $G^R_{dd}(\omega)$:

$$G^R_{kk'} = G^{0R}_{kk'} + G^{0R}_{kk}\tau G^R_{dk'},$$
$$G^R_{dk'} = G^{0R}_{dd}\tau \sum_k G^R_{kk'},$$
$$G^R_{kd} = G^{0R}_{kk}\tau G^R_{dd},$$
$$G^R_{dd} = G^{0R}_{dd} + G^{0R}_{dd}\tau \sum_k G^R_{kd}. \tag{6}$$

Unperturbed bound state Green's function $G^{0R}_{dd}(\omega)$ has the following form:

$$G^{0R}_{dd}(\omega) = (\omega - \varepsilon_\sigma + i\gamma)^{-1} \tag{7}$$

with γ being the nonradiative recombination rate from the bound state. The system of equations (6) can be solved exactly, the solution is the bound state Green's function $G^R_{dd}(\omega)$ accounting for the interaction with 2D electrons in the QW:

$$G^R_{dd}(\omega) = \frac{1}{\omega - \varepsilon_\sigma + i[\gamma + \Gamma(\omega)] - \tau^2 \cdot N(\omega)}, \tag{8}$$

where

$$N(\omega) = Re \sum_k G^{0R}_{kk}(\omega). \tag{9}$$

The tunneling rate $\Gamma(\omega) = \nu_0(\omega)\tau^2$ corresponds to the electron transitions between the QW and the bound state, $\nu_0(\omega)$ is the unperturbed density of states in the QW. As we assume the negligibly weak dependence of τ on k, for the 2D density of states in the QW the tunneling rate is a constant Γ, which we take as a parameter. The Green's functions $G^{0R}_{kk}(\omega)$ for electrons in the QW can be written as:

$$G^{0R}_{kk}(\omega) = (\omega - \varepsilon_k + i\gamma_w)^{-1}, \tag{10}$$

where γ_w describes radiative recombination processes in the QW. The modification of the continuum spectrum due to the coupling with the bound state expressed by Eq. (8) is, of course, the well known Fano-Anderson problem [32]. Along with the modification of the continuum spectrum the coupling results in appearance of split-off states below the edge of the continuum. The energy of these states for both spin directions are given by the equation:

$$\omega - \varepsilon_\sigma - \tau^2 N(\omega) = 0 \tag{11}$$

or

$$\omega - \varepsilon_\sigma - \tau^2 \int d\varepsilon \frac{\nu_0(\varepsilon)}{\omega - \varepsilon} = 0. \quad (12)$$

Equation (12) has a real solution in the bad gap $\omega = \tilde{\varepsilon}_\sigma \in (E_v, E_c)$, where E_c and E_v are the conduction band and the valence band edges in the QW, respectively (see Fig. 1). The energy of the split-off states depends on the tunneling transfer amplitude τ and the unperturbed density of states in the QW. In the 2D case $\nu_0(\omega) = W^{-1}$, where W is the conduction band width. For $\omega < E_c$ one obtains the energy of the split-off states

$$\tilde{\varepsilon}_\sigma = E_c - W \cdot e^{-\frac{\varepsilon_\sigma - E_c}{\Gamma}},$$
$$\tilde{\varepsilon}_{U\sigma} = E_c - W \cdot e^{-\frac{\varepsilon_\sigma + U - E_c}{\Gamma}}. \quad (13)$$

These expressions are valid for $E_c - \varepsilon_\sigma > \Gamma$. In the case $E_c - \varepsilon_\sigma < \Gamma$ one should replace E_c by $\tilde{\varepsilon}_\sigma$ in the exponent and solve the nonlinear equation on $\tilde{\varepsilon}_\sigma$. In both cases the split-off states $\tilde{\varepsilon}_\sigma$ reside in the band gap rather close to the band bottom E_c as $\Gamma \sim W^{-1}$. In the absence of Coulomb interaction in the case of 2D band in the QW for ω close to the split-off state energy $\tilde{\varepsilon}_\sigma$ the Green's functions of the bound state can be written as:

$$G^R_{dd\sigma}(\omega) \simeq \frac{A_\sigma}{\omega - \tilde{\varepsilon}_\sigma + i\gamma}, \quad (14)$$

where the coefficient A_σ in the nominator exponentially depends on the value of the bound state energy ε_σ

$$A_\sigma = W/\Gamma \cdot e^{-\frac{\varepsilon_\sigma - E_c}{\Gamma}}. \quad (15)$$

The density of states in the QW is given by

$$\sum_{kk'} G^R_{kk'} = \frac{B_\sigma}{\omega - \tilde{\varepsilon}_\sigma + i\gamma} \quad (16)$$

with the coefficient B_σ being

$$B_\sigma = \frac{(E_c - \varepsilon_\sigma)^2}{\Gamma^2} \cdot e^{-\frac{\varepsilon_\sigma - E_c}{\Gamma}}. \quad (17)$$

In the presence of Coulomb correlations one can get an expression for the retarded Green's function $G^R_{dd\sigma}(\omega)$:

$$G^R_{dd\sigma}(\omega) = (1 - \tilde{n}_{-\sigma}) \times \left\{ \frac{[1 - A_\sigma(\tilde{\varepsilon}_\sigma)] \cdot \Theta(W - |\omega|)}{\omega - \varepsilon_\sigma + i(\Gamma + \gamma)} + \frac{A_\sigma(\tilde{\varepsilon}_\sigma) \cdot \Theta(|\omega| - W)}{\omega - \tilde{\varepsilon}_\sigma + i\gamma} \right\}$$
$$+ \tilde{n}_{-\sigma} \times \left\{ \frac{[1 - A_{U\sigma}(\tilde{\varepsilon}_\sigma)] \cdot \Theta(W - |\omega|)}{\omega - \varepsilon_\sigma - U + i(\Gamma + \gamma)} + \frac{A_{U\sigma}(\tilde{\varepsilon}_\sigma) \cdot \Theta(|\omega| - W)}{\omega - \tilde{\varepsilon}_{U\sigma} + i\gamma} \right\}, \quad (18)$$

where $\tilde{n}_\sigma$ is the occupation of the bound state renormalized by the presence of the split-off states. One should substitute ε_σ by $\varepsilon_\sigma + U$ in expression (15) to get $A_{U\sigma}$

$$A_{U\sigma} = W/\Gamma \cdot e^{-\frac{\varepsilon_\sigma + U - E_c}{\Gamma}}. \quad (19)$$

System of equations for time evolution of the split-off states occupation numbers has the following form [33]:

$$\frac{\partial n_\sigma}{\partial t} = -2\Gamma \left[n_\sigma - \int \frac{d\omega}{2\pi} \left(G^A_{dd\sigma}(\omega) - G^R_{dd\sigma}(\omega) \right) f_k(\omega) \right] - \gamma n_\sigma. \quad (20)$$

Considering Eqs. (18) and solving the system of equations (20) in the stationary case ($\frac{\partial n_\sigma}{\partial t} = \frac{\partial n_{-\sigma}}{\partial t} = 0$) one can get stationary occupation numbers for the bound state:

$$n^{st}_\sigma = \frac{\Phi(\varepsilon_\sigma) - \Delta\Phi^{-\sigma} \cdot \Phi(\varepsilon_\sigma)}{1 - \Delta\Phi^\sigma \cdot \Delta\Phi^{-\sigma}}, \quad (21)$$

where

$$\Phi(\varepsilon_\sigma) = (1 - A_\sigma)N(\varepsilon_\sigma) + A_\sigma \tilde{N}(\tilde{\varepsilon}_\sigma),$$
$$\Phi_U(\varepsilon_\sigma) = (1 - A_{U\sigma})N(\varepsilon_\sigma + U) + A_{U\sigma}\tilde{N}(\tilde{\varepsilon}_{U\sigma}) \quad (22)$$

and

$$\Delta\Phi^\sigma = \Phi(\varepsilon_\sigma) - \Phi_U(\varepsilon_\sigma). \quad (23)$$

QW occupation functions $N(\varepsilon)$ and $\tilde{N}(\varepsilon)$ have the form:

$$N(\varepsilon) = \int d\varepsilon_k \frac{f_k(\varepsilon_k) \cdot (\Gamma + \gamma)}{(\varepsilon_k - \varepsilon)^2 + (\Gamma + \gamma)^2},$$
$$\tilde{N}(\varepsilon) = \int d\varepsilon_k \frac{f_k(\varepsilon_k) \cdot \gamma}{(\varepsilon_k - \varepsilon)^2 + \gamma^2}. \quad (24)$$

III. RESULTS AND DISCUSSION

Let us first clarify the effect of the tunnel coupling between the spin degenerate bound state and the QW on the occupation number of the bound state. Figure 2 shows the occupation

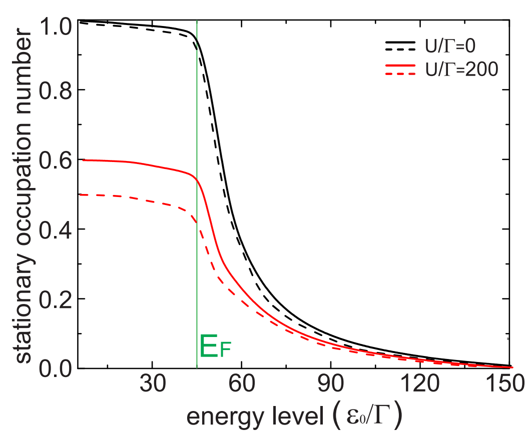

FIG. 2. Stationary occupation of the spin degenerate bound state energy level as a function of energy level position ε relative to the Fermi level $E_F/\Gamma = 45$. Dashed lines correspond to the absence of the tunnel coupling; solid lines correspond to the presence of the tunnel coupling and formation of the split-off state. Black curves correspond to the absence of Coulomb correlations $U/\Gamma = 0$; red curves are for the case when Coulomb correlations are present at the bound state $U/\Gamma = 200$. Parameters $W/\Gamma = 1000$, $\gamma = 0.5$, and $\Gamma = 1$ are the same for all the figures. The green vertical line is the Fermi level position.

number of the bound state calculated using Eq. (21) as a function of the bound state energy. Black curves correspond to the case when on-site Coulomb interaction at the bound state is neglected. In the absence of the tunnel coupling (dashed black curve in Fig. 2) the bound state is not hybridized with the QW continuum. It still has a finite width being broadened by the nonradiative relaxation rate γ as clearly seen from Eq. (7). Therefore, its occupation is simply the Fermi distribution function integrated with the Lorentz-like density of states of the broadened bound state level. Assuming zero temperature the occupation number is equal to one when the bound state energy lies far below the Fermi level E_F and decreases when the bound state energy approaches E_F due to the broadening of the state, then it rapidly goes to zero when the bound state is well above the Fermi level E_F. When the tunnel coupling is enabled, the broadening of the bound state increases by Γ and also the split-off state emerges below the continuum. The effect of the latter appears to be significant as it slightly increases the bound state occupation despite the enhanced broadening as can be seen in Fig. 2, solid black curve. Its role becomes even more pronounced in the case of the on-site Coulomb interactions as discussed below. The results of the calculations with the on-site Coulomb interaction taken $U = 200\Gamma$ are shown by the red curves in Fig. 2. The presence of Coulomb correlations lead to the strong decrease of the energy level occupation when it is localized below the Fermi level. It occurs due to the Coulomb blockade effect as tunneling of the electron from the QW requires an additional energy cost if the bound state is occupied. The presence of the split-off state leads to the additional contribution to the bound state occupation and results in a substantial increase of the bound state occupation when it is localized below the Fermi level. Note that the effect is stronger than in the $U = 0$ case. That is because the double-occupied energy level located well above E_F also produces a split-off replica below the band bottom, hence, occupied.

Let us now discuss the case when the bound state is split in the electron spin projection. Consequently, the tunneling from the QW is possible into two different bound state energy levels corresponding to the opposite spin projections. Naturally, for the most efficient spin dependent tunneling the Fermi level should fall between the spin-split bound state levels. It could be done by doping of the QW or external gating. Panels (a) and (b) in Fig. 3 demonstrate stationary occupation of the bound state energy levels with given spin projection in the presence of Coulomb correlations as a function of a difference between the bound state spin sublevels $\Delta_0 = (\varepsilon_\downarrow - \varepsilon_\uparrow)$. The sign of Δ_0 could be controlled by the external magnetic field applied to the system. We do not account for the spin splitting of the electron states in the QW. This is the typical case as usually there is a strong difference of g factors for magnetic ions and electrons in the QW in heterostructures of the considered type [34].

For the calculation results presented in Fig. 3 the spin-up energy level position $\varepsilon_\uparrow$ is fixed either slightly below the Fermi level E_F (a) or well below E_F (b). All the curves in Fig. 3 are with account for the on-site Coulomb correlations; the color is related to the sign of the spin projection. Dashed curves correspond to the absence of the tunnel coupling; solid curves account for the tunnel coupling and the split-off states.

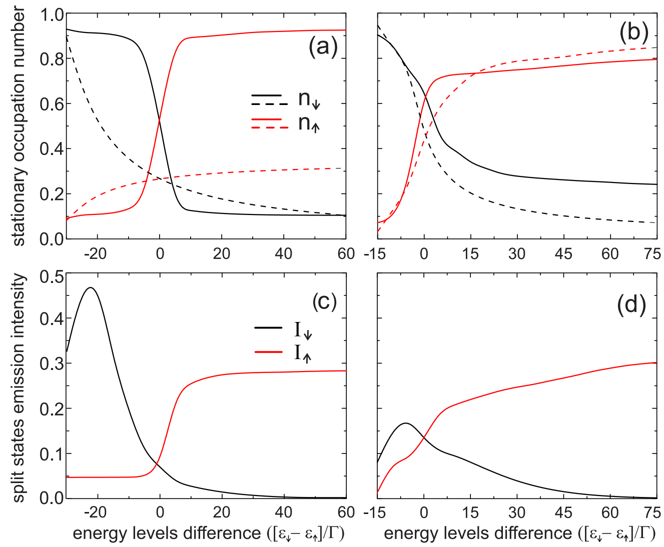

FIG. 3. Stationary occupation of the spin-split bound state energy levels $\varepsilon_\uparrow$ and $\varepsilon_\downarrow$ (a), (b) and the split-off states PL intensity (c), (d). In panels (a), (b) solid curves correspond to the situation when the split-off states are present in the system and dashed curves demonstrate results in the case when the split-off states are neglected. In panels (c), (d) black curves correspond to the $\tilde{\varepsilon}_\downarrow/\Gamma$ state and red curves describe $\tilde{\varepsilon}_\uparrow/\Gamma$ state. For panels (a), (c) $\varepsilon_\uparrow/\Gamma = 40$, for panels (b), (d) $\varepsilon_\uparrow/\Gamma = 15$. Parameters $W/\Gamma = 1000$, $U/\Gamma = 200$, $\varepsilon_F/\Gamma = 45$, $\gamma = 0.5$, and $\Gamma = 1$ are the same for all the figures.

For both situations the behavior of stationary occupation is quite similar: The lower energy level is more occupied than the higher one. The most interesting fact, however, is that the presence of the split-off states substantially modifies the occupation of the bound state.

The split-off states are located below the bottom of the band (E_c). Being occupied, they partially contribute to the bound state occupation a_σ and the QW occupation b_σ. As follows from Eqs. (14)–(17), these partial contributions of the split-off states obey

$$\frac{a_\sigma}{b_\sigma} = \frac{A_\sigma}{B_\sigma} = \frac{W\Gamma}{(E_c - \varepsilon_\sigma)^2}. \tag{25}$$

Let us now, for simplicity, assume single occupation of the split-off state $a_\sigma + b_\sigma = 1$, we then get

$$a_\sigma = \frac{W\Gamma}{(E_c - \varepsilon_\sigma)^2 + W\Gamma}. \tag{26}$$

If $n = 1$ is the occupation number of the bound state then the average energy of an electron located at the bound state reads

$$\varepsilon = [(1 - a_\sigma)\varepsilon_\sigma + a_\sigma \tilde{\varepsilon}_\sigma], \tag{27}$$

where the energy of the split-off state $\tilde{\varepsilon}_\sigma$ is close to the band bottom and can be estimated as $\tilde{\varepsilon}_\sigma \approx E_c$. Substituting (26) into (27) we get:

$$\varepsilon - \varepsilon_\sigma = -(\varepsilon_\sigma - E_c)\left[1 + \frac{(\varepsilon_\sigma - E_c)^2}{W\Gamma}\right]^{-1}. \tag{28}$$

As (28) suggests, the strong enough tunnel coupling renormalizes the energy of the electron localized at the bound state, the effect is more pronounced when ε_σ is further away from the band bottom E_c provided that $\varepsilon_\sigma - E_c < \sqrt{W\Gamma}$. That qualitatively explains the behavior observed in Figs. 3(a) and 3(b). When the spin up energy level $\varepsilon_\uparrow$ is localized slightly below the Fermi level [see panel (a) in Fig. 3], the split-off state significantly increases the bound state occupation for the spin down energy level as the renormalization of the bound state level is substantial. Simultaneously, occupation of the spin-up bound state energy level becomes smaller. When the sign of the detuning Δ_0 is changed, the situation is reversed, the split-off states result in the increase of the spin up energy level occupation and decrease of the spin down level occupation.

For the energy levels ε_σ located well below the Fermi level, thus closer to the band bottom, the effect of the split-off states becomes weaker as the energy renormalization is smaller. Indeed, if a bound state coincides with the band bottom, there is no renormalization at all. This is the case for the most left point in Fig. 3(b); when the level $\varepsilon_\downarrow$ reaches E_c the tunnel coupling does not affect the occupation at all except for the small difference due to enhanced broadening of the bound state level due to Γ. Also as the Coulomb interaction is accounted for, the occupation number would never exceed unity in Fig. 3.

The effect of the tunnel coupling on the system spectrum and occupation is also reflected in the intensity of the PL from the QW, which is presented in Fig. 3 panels (c) and (d). The emergence of the split-off states below the edge of the conductance band leads to the appearance of additional peaks in the frequency range corresponding to the QW band gap. The calculated structure and occupation of these states allowed us to obtain the PL intensity, which is proportional to A_σ [see Eq. (15)]. An important feature of the band gap peaks is that the PL corresponding to each of the two split-off states is fully circularly polarized. Depending on the tunable parameters of the system both peaks can have similar amplitudes when the split-off states are close to each other and there also exists a possibility for the amplitude of one peak to exceed the amplitude of the other one. Parameters for the panels (c) and (d) are the same as for the panels (a) and (b), respectively.

In the considered model the PL spectra from the QW is constant right above the band gap frequency. While in this work we focus on the split-off states, it is worth noting that the continuous part of the spectra has features at the frequencies corresponding to the bound state levels ε_σ which would be also polarized as the resonance position is different for the spin-split sublevels. For the detailed study see Refs. [35,36].

The circular polarization of the band gap PL peaks corresponds to the electron spin projection at the bound state. In the QW we assume a radiative recombination of the electrons with 2D heavy holes with the projections of total angular momentum $j = \pm 3/2$ recombining with electrons $j = \mp 1/2$ and emitting right- (σ^+) and left- (σ^-) circularly polarized light, respectively. The position of the peaks relative to the continuous PL spectra is given by $\tilde{\varepsilon}_\sigma$ [see Eq. (13)]. An example of the calculated PL peaks indicating their polarization is presented in Fig. 4.

To estimate the peaks amplitudes and positions we consider parameters relevant for GaAs heterostructures doped with Mn [29]. We take $\Gamma = 13$ meV, $\epsilon_\sigma - E_c \approx 50$ meV, $\Delta_0 =$ 2.5 meV, $W = 2$ eV. The amplitude of both peaks reach ≈30 percent of the continuous QW PL spectra intensity and their position appear to be 39 meV and 47 meV below E_c for σ^+ and σ^- polarizations, respectively.

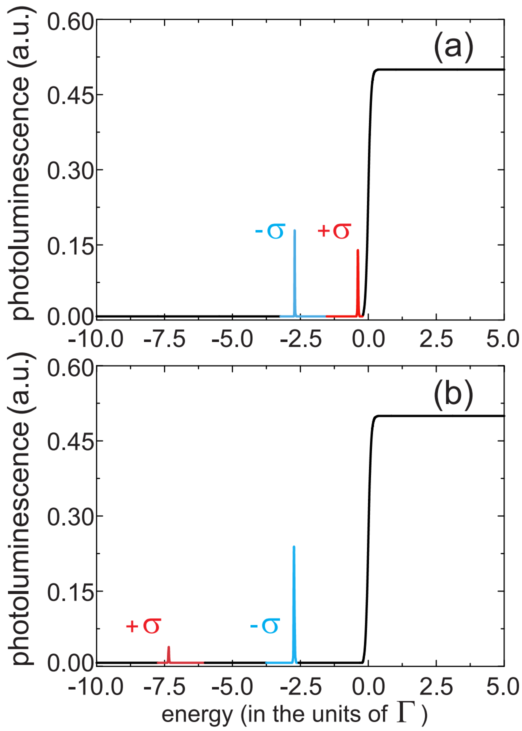

FIG. 4. Photoluminescence as a function of frequency demonstrating two narrow well resolved peaks with right and left circular polarization. (a) $\varepsilon_\downarrow/\Gamma = 30$ and $\varepsilon_\uparrow/\Gamma = 20$; (b) $\varepsilon_\downarrow/\Gamma = 10$ and $\varepsilon_\uparrow/\Gamma = 20$. Parameters $E_c = 0$, $W/\Gamma = 1000$, $U/\Gamma = 200$, $\gamma = 0.5$ and $\Gamma = 1$ are the same for all the figures.

We would like to emphasize two important features of the solitary PL peaks corresponding to the split-off states. They have opposite circular polarization and their spectral position strongly depends on the tunnel coupling parameter τ. These properties enable us to come up with an idea of an ultrafast modulation of the PL polarization by electrical means. For the taken parameters a change of the tunneling rate Γ by a factor of 1.5 results in the shift of the peak positions by ≈30 meV, quite enough for the experimental realization. Such modulation of the tunnel barrier transparency can be achieved by application of an electric field by means of an external gate [37]. With fixing the output frequency such modulation would allow for switching of the PL polarization as the spin-up and spin-down split-off states position is alternated with the change of the tunnel coupling. As the characteristic tunneling time is of a picosecond order, this mechanism allows for ultrafast modulation of the PL polarization and could be promising for applications in spintronics, in particular, for ultrafast polarization modulation in spin lasers.

IV. SUMMARY

We present a theoretical analysis of the split-off states and their influence on the spin polarization in the semiconductor QW coupled to a remote bound state with account for on-site

Coulomb correlations and spin splitting. The split-off states are formed in the band gap of the QW due to the tunnel coupling between the QW and the bound state. It has been shown that the split-off states strongly affect the occupation of the bound state and modify the PL signal from the QW. We have shown that the split-off states reveal themselves as solitary peaks, which are fully circularly polarized with their spectral position being very sensitive to the tunnel coupling. The obtained results open a new possibility to control spin polarization in nanoscale systems. In particular, a gate assisted modulation of the tunnel barrier transparency allows for ultrafast switching of the PL circular polarization from the QW. Therefore, the suggested mechanism is promising for polarization modulation devices such as spin lasers.

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