Dynamic spin injection into a quantum well coupled to a spin-split bound state

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We present a theoretical analysis of dynamic spin injection due to spin-dependent tunneling between a quantum well (QW) and a bound state split in spin projection due to an exchange interaction or external magnetic field. We focus on the impact of Coulomb correlations at the bound state on spin polarization and sheet density kinetics of the charge carriers in the QW. The theoretical approach is based on kinetic equations for the electron occupation numbers taking into account high order correlation functions for the bound state electrons. It is shown that the on-site Coulomb repulsion leads to an enhanced dynamic spin polarization of the electrons in the QW and a delay in the carriers tunneling into the bound state. The interplay of these two effects leads to non-trivial dependence of the spin polarization degree, which can be probed experimentally using time-resolved photoluminescence experiments. It is demonstrated that the influence of the Coulomb interactions can be controlled by adjusting the relaxation rates. These findings open a new way of studying the Hubbard-like electron interactions experimentally.

I. INTRODUCTION

The research field of spintronics continues to grow covering various spin phenomena in solid-state physics. The first generation or "metallic" spintronics is associated with magnetism and exchange interaction. It has succeeded in suggesting practical applications, the vivid example is the giant magnetoresistance effect widely used in hard drives. Semiconductor spintronics of the second generation is focused on the effects based on spin-orbit interaction, which locks a particle motion with its spin. This locking is key to many attracting practical applications. In the traditional semiconductor materials the effective spin-orbit interaction is relatively weak, so the latest research in the field of spintronics develops in two overlapping directions: new semiconductors with stronger spin-orbit interaction, including topological insulators, and new spin phenomena based on the exchange interaction. Our work contributes to the second direction, in the present paper we focus on a dynamical spin injection into a quantum well (QW) through a tunnel barrier with account for the exchange interaction in the leads.

The spin injection into a semiconductor remains the cornerstone of modern spintronics. The conductivity mismatch prevents an efficient spin injection from a ferromagnetic metal into a semiconductor. The widely discussed solutions of this problem (apart from those based on spin-orbit interaction) include using dilute magnetic semiconductor as a spin injector, spin injection from a ferromagnet through a tunnel barrier, superdiffusive spin transport. In our paper we consider a spin injection into a semiconductor due to spin-dependent relaxation of initially unpolarized ensemble of charge carriers. In this sense it is similar to spin-dependent recombination phenomena.

In our model a non-equilibrium distribution of the non-polarized carriers is assumed created instantaneously in the QW. We analyse theoretically the subsequent kinetics of the spin polarization. The study of this physical model is motivated by the experimental studied reported in Refs. In these experiments an InGaAs based heterostructures with a QW and remote Mn doping layer was optically pumped with non-polarized carriers; the experimentally measured time-resolved photoluminescence indicated a development of non-stationary spin polarization of the 2D carriers in the QW on the time scale smaller than the radiative recombination time. The origin of this phenomena was explained theoretically as electron tunneling from the QW into a Mn dopant layer. The spin splitting of the impurity bound state in the dopant layer due to exchange interaction results in the spin-dependent tunneling rate and thus leads to the spin polarization of the carriers remaining in the QW.

In this paper we generalize the theory by developing non-stationary formalism to describe the charge and spin kinetics in the QW coupled to a bound state. An important feature of our study is that in addition to the conventional Heisenberg exchange interaction we account for the Coulomb correlations at the bound state which can contribute to the spin splitting as it is well known from Anderson and Stoner model. We show how the kinetics of the spin polarization in the QW depends on the relaxation rates and the strength of the Coulomb on-site correlation at the bound state.

II. THEORETICAL MODEL

Although motivated by the experiments, in this paper we do not restrict ourselves to a particular semiconductor heterostructure design. Let us consider two
The model systems shown in Fig. 1 correspond to the design of experimentally studied (Ga,In,Mn)As heterostructures with bound states formed by paramagnetic impurities, (a) is alternative design with bound states formed by a quantum dot (QD).

A somewhat different model system is shown in Fig. 1b. Here the bound state is formed by Mn ions in interstitial position providing donor-like states. The delta-doping Mn layer is located at a distance of 3-10 nm from the QW. The spin splitting is due to the effective exchange field in the doping Mn layer and the relaxation from the impurity donor levels is due to a very fast non-radiative recombination with the holes in the low-temperature grown (Ga,Mn)As layer.

Fig. 1. Two possible realizations of the considered system: (a) corresponds to the design of experimentally studied (Ga,In,Mn)As heterostructures with bound states formed by paramagnetic impurities, (b) is an alternative design with bound states formed by a quantum dot (QD).

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

$$\hat{H} = \hat{H}_{QW} + \hat{H}_1 + \hat{H}_{int}, \quad (1)$$

where $\hat{H}_{QW}$ describes the QW:

$$\hat{H}_{QW} = \sum_{\sigma,k} \varepsilon_k c_{1\sigma}^\dagger c_{1\sigma}, \quad (2)$$

$\hat{H}_1$ describes the bound state with the Hubbard term for on-site Coulomb repulsion:

$$\hat{H}_1 = \sum_{\sigma} \varepsilon_1 n_{1\sigma}^\sigma + U n_{1\sigma}^\sigma n_{1\sigma}^{-\sigma}, \quad (3)$$

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

$$\hat{H}_{T} = \sum_{k\sigma} t_k (c_{1\sigma}^\dagger c_{k\sigma} + c_{k\sigma}^\dagger c_{1\sigma}).$$

Here index $k$ labels continuous spectrum states in the QW, $t_k$ is the tunneling transfer amplitude between QW states and the bound state. The bound state is characterized by the energy level $\varepsilon_1$, which can be split due to an exchange interaction or an external magnetic field into two spin sublevels with the energies $\varepsilon_\sigma = \varepsilon_1 + \sigma \Delta$, where $\sigma = \pm 1/2$ is the electron spin projection and $\Delta$ is the energy gap. Operators $c_{1\sigma}^\dagger, c_{1\sigma}$ are the creation and annihilation operators for the electrons in the QW. $n_{1\sigma} = c_{1\sigma}^\dagger c_{1\sigma}$ is the occupation number operator for the bound state, operator $c_{1\sigma}$ destroys electron in the bound state with the spin projection $\sigma$. $U$ is the on-site Coulomb repulsion energy for the doubly occupied bound state. Further analysis deals with the low temperature regime when the Fermi level is well defined and the temperature is much lower than all other energy scales in the system.

III. NON-STATIONARY ELECTRONIC TRANSPORT FORMALISM

Let us further consider $\hbar = 1$ and $e = 1$ elsewhere, so the equations of motion for the electron operators products $\hat{n}_{1\sigma}^\sigma = c_{1\sigma}^\dagger c_{1\sigma}, \hat{n}_{1\sigma} = c_{1\sigma}^\dagger c_{1\sigma}$ and $\hat{n}_{k\sigma}^\sigma = c_{k\sigma}^\dagger c_{k\sigma}$ can be written as:

$$i \frac{\partial \hat{n}_{1\sigma}^\sigma}{\partial t} = - \sum_{k,\sigma} t_k \cdot (\hat{n}_{k\sigma} - \hat{n}_{1\sigma}^\sigma), \quad (4)$$

$$i \frac{\partial \hat{n}_{1\sigma}^\sigma}{\partial t} = - (\varepsilon_{1\sigma}^\sigma - \varepsilon_k^\sigma) \cdot \hat{n}_{1\sigma}^\sigma - U \cdot \hat{n}_{1\sigma}^{-\sigma} \hat{n}_{1\sigma}^\sigma + t_k \cdot \hat{n}_{1\sigma}^\sigma - \sum_{k' \neq k} t_{k'} \cdot \hat{n}_{k'\sigma}^\sigma, \quad (5)$$

$$i \frac{\partial \hat{n}_{k\sigma}^\sigma}{\partial t} = - (\varepsilon_{k\sigma}^\sigma - \varepsilon_{k'\sigma}^\sigma) \cdot \hat{n}_{k\sigma}^\sigma - t_{k'} \cdot \hat{n}_{k'\sigma}^\sigma + t_{k}\cdot\hat{n}_{k'\sigma}^\sigma. \quad (6)$$
Following the logic of Ref. 19 we substitute the solution of Eq. (6) into Eq. (5) to obtain:

$$i \frac{\partial n_{1k}^\sigma}{\partial t} = -(\epsilon_1^\sigma - \epsilon_k + i\Gamma_k) n_{1k}^\sigma - U n_{11}^{-\sigma} n_{1k}^\sigma + t_k(n_{1k}^\sigma - \bar{n}_k^\sigma) + i \sum_{k' \neq k} t_{k'k} \int_0^t e^{i(\epsilon_k' - \epsilon_k)(t-t')} n_{k'1}^\sigma \, dt', \quad \text{Eq. (7)}$$

where $\Gamma_k = \pi \nu_0 (\epsilon_k) \Gamma_k^2$ and $\nu_0 (\epsilon_k)$ is the unperturbed density of states in the QW. Further we assume that the tunneling parameter $t_k$ has a negligibly weak dependence on $k$, so for 2D density of states in the QW the tunneling relaxation rate is a constant $\Gamma_k \equiv \Gamma$, which we take as a parameter. If condition $\frac{\epsilon_k - \epsilon_1^-}{\Gamma} \gg 1$ is fulfilled, $\bar{n}_k^\sigma$ is a slowly varying quantity in comparison with $\hat{n}_{1k}^\sigma$. Consequently, it is reasonable to consider that:

$$\frac{\partial}{\partial t} \hat{n}_{1+}^{\sigma} \hat{n}_{1k}^- \sim \hat{n}_{1-}^{\sigma} \frac{\partial}{\partial t} \hat{n}_{1k}^+,$$

(8)

Taking into account that $(\hat{n}_{1k}^+)^2 = \hat{n}_{1k}^+ \hat{n}_{1k}^- \cdot (1 - \hat{n}_{1k}^-) \cdot \hat{n}_{1k}^+ = 0$, one can find expressions for $(1 - \hat{n}_{1}^{-\sigma}) \hat{n}_{1k}^{\sigma}$, $\hat{n}_{1}^{-\sigma} \hat{n}_{1k}^{\sigma}$ and obtain the equations for the time evolution of the particle number operators $\hat{n}_{1}^{\sigma}$, $\hat{n}_{1k}^{\sigma}$ for the bound state and QW, respectively. The suggested theoretical approach allows one to analyze dynamic spin injection due to the spin-dependent tunneling with account for high order correlation functions for the bound state electrons. Therefore, it gives possibility to analyze the effects of the on-site Coulomb repulsion.

We also add explicitly the spin-independent relaxation terms describing recombination in the QW with the rate $\gamma_k$ and relaxation at the bound state due to non-radiative recombination (Fig. 1a) or tunnel leakage into the lead (Fig. 1b) with the rate $\gamma_1$. Thus, we account for the effect of Coulomb correlations on the tunneling between QW and the bound state and also for the additional bound state broadening due to the relaxation into a lead or reservoir. This approach neglects the influence of the QW and bound state relaxation channels on each other, but allows for decoupling of the QW and bound state equations of motion. Therefore, we obtain:

$$\frac{\partial \hat{n}_{1k}^\sigma}{\partial t} = -2 \Gamma \cdot (\hat{n}_{1k}^\sigma - (1 - \hat{n}_{1}^{-\sigma}) \cdot \hat{\Phi}(\epsilon_{\sigma}) - \hat{n}_{1}^{-\sigma} \cdot \hat{\Phi}(\epsilon_{\sigma} + U) - \gamma_1 \cdot \hat{n}_{1k}^\sigma),$$

$$\frac{\partial \hat{n}_{1k}^\sigma}{\partial t} = \frac{2 \Gamma}{\nu_0 \pi} \cdot [(1 - \hat{n}_{1}^{-\sigma})(\hat{n}_{1k}^\sigma - \hat{n}_{1k}^\sigma)] \frac{\tilde{\Gamma}}{(\epsilon_{\sigma} - \epsilon_k)^2 + \Gamma^2} + \frac{\hat{n}_{1}^{-\sigma} \cdot (\hat{n}_{1k}^\sigma - \hat{n}_{1k}^\sigma) \tilde{\Gamma}}{(\epsilon_{\sigma} + U - \epsilon_k)^2 + \Gamma^2} - \gamma_k \cdot \hat{n}_{1k}^\sigma.$$

(9)

Here we introduced the QW occupation operators $\hat{\Phi}(\epsilon_{\sigma})$ and $\hat{\Phi}(\epsilon_{\sigma} + U)$ as:

$$\hat{\Phi}(\epsilon_{\sigma}) = \int d\epsilon_k \cdot \hat{f}_{k}^\sigma(\epsilon_k) \cdot \frac{1}{\pi} \frac{\tilde{\Gamma}}{(\epsilon_{\sigma} - \epsilon_k)^2 + \Gamma^2},$$

$$\hat{\Phi}(\epsilon_{\sigma} + U) = \int d\epsilon_k \cdot \hat{f}_{k}^\sigma(\epsilon_k + U) \cdot \frac{1}{\pi} \frac{\tilde{\Gamma}}{(\epsilon_{\sigma} + U - \epsilon_k)^2 + \Gamma^2}.$$

(10)

where $\tilde{\Gamma} = \Gamma + \gamma_1$. Note, that we introduced $\tilde{\Gamma}$ in order to properly account for the structure of the bound state, which is affected both by the hybridization with the QW and the separate relaxation channel with the rate $\gamma_1$.

One can obtain equations for the bound state occupation numbers $n_{1k}^\sigma$ by averaging equations for the operators and by decoupling electron occupation numbers for the QW states from the bound state occupation numbers. Such decoupling procedure is reasonable if one considers that Kondo correlations can be neglected. After decoupling the QW occupation numbers operators $\hat{n}_{1k}^\sigma$ in Eqs. (9)-(10) have to be replaced by the distribution functions $f_{k}^\sigma$. In order to take into account spin-independent relaxation processes from the QW and the bound state we add the corresponding rates $\gamma_k$ and $\gamma_1$ into kinetic equations for the bound state occupation numbers and the QW electron distribution function. Assuming that equilibrium state corresponds to the empty bound state and empty QW we obtain the following equations:

$$\frac{\partial n_{1k}^\sigma}{\partial t} = -2 \cdot \Gamma \cdot I_{k}^\sigma - \gamma_1 \cdot n_{1k}^\sigma,$$

$$\frac{\partial f_{k}^\sigma}{\partial t} = 2 \cdot \Gamma \cdot J_{k}^\sigma - \gamma_k \cdot f_{k}^\sigma,$$

(11)

where

$$I_{k}^\sigma = n_{1k}^\sigma - (1 - n_{1}^{-\sigma}) \cdot \hat{\Phi}(\epsilon_{\sigma}) - n_{1}^{-\sigma} \cdot \hat{\Phi}(\epsilon_{\sigma} + U)$$

$$J_{k}^\sigma = \frac{1}{\nu_0 \pi} \frac{[(1 - n_{1}^{-\sigma})(n_{1k}^\sigma - f_{k}^\sigma)]}{(\epsilon_{\sigma} - \epsilon_k)^2 + \Gamma^2} \frac{\tilde{\Gamma}}{(\epsilon_{\sigma} + U - \epsilon_k)^2 + \Gamma^2},$$

(12)

and QW occupation functions $\hat{\Phi}(\epsilon_{\sigma})$ and $\hat{\Phi}(\epsilon_{\sigma} + U)$ read

$$\hat{\Phi}(\epsilon_{\sigma}) = \int d\epsilon_k \cdot f_{k}^\sigma(\epsilon_k) \cdot \frac{1}{\pi} \frac{\tilde{\Gamma}}{(\epsilon_{\sigma} - \epsilon_k)^2 + \Gamma^2},$$

$$\hat{\Phi}(\epsilon_{\sigma} + U) = \int d\epsilon_k \cdot f_{k}^\sigma(\epsilon_k + U) \cdot \frac{1}{\pi} \frac{\tilde{\Gamma}}{(\epsilon_{\sigma} + U - \epsilon_k)^2 + \Gamma^2}.$$

(13)

We further solve Eqs. (11)-(12) implying the following initial conditions at $t = 0$: the bound state is empty, therefore $n_{1k}^\sigma = \hat{n}_{1k}^{-\sigma} = 0$; the QW is filled by the photoexcited carriers with a non-equilibrium energy distribution function characterized by chemical potential $\mu^*$ and electron temperature $T$: $f_{k}(0) = \frac{1}{e^{(\epsilon_k - \mu^*)/\hbar \epsilon_k} + 1}$. 

The spin polarization of the electrons remaining in the QW is given by $N^\uparrow - N^\downarrow$, where $N_\sigma = \int f_\sigma (\varepsilon_k) d\varepsilon_k$ it manifests itself in the circular polarization of the photoluminescence (PL) from the QW which can be measured. The spin polarization degree which would be measured by optical means is defined as:

$$P = \frac{N^\uparrow - N^\downarrow}{N^\uparrow + N^\downarrow}.$$  

The polarization degree $P$ is negative when electrons with spin projection $\sigma = -\frac{1}{2}$ prevail. The considered theoretical approach is more general than in Ref.\cite{10} as Eqs. (11) cover both cases of resonant and non-resonant tunneling between the QW and the bound state.

IV. RESULTS AND DISCUSSION

The spin kinetics calculated according to the theory described above is shown in Figs. 2-4. In all the calculations we assume the same transparency of the tunnel barrier.

The spin polarization and QW sheet density for different QW relaxation rates $\gamma_k$. For solid lines $U/\Gamma = 35$, for dashed lines $U/\Gamma = 0$. Parameters $\varepsilon^\uparrow/\Gamma = 2$, $\varepsilon^\downarrow/\Gamma = -2$, $\mu^*/\Gamma = 0$, $\gamma_k/\Gamma = 0.15$ and $\Gamma = 1$ are the same for all the figures.

FIG. 2. (Color online) Time evolution of spin polarization and QW sheet density for different QW relaxation rates $\gamma_k$. For solid lines $U/\Gamma = 35$, for dashed lines $U/\Gamma = 0$. Parameters $\varepsilon^\uparrow/\Gamma = 2$, $\varepsilon^\downarrow/\Gamma = -2$, $\mu^*/\Gamma = 0$, $\gamma_k/\Gamma = 0.15$ and $\Gamma = 1$ are the same for all the figures.

FIG. 3. (Color online) Time evolution of spin polarization and QW sheet density for different bound state relaxation rates $\gamma_1$. Insert in the panel (a) demonstrates that for large values of $\gamma_1/\Gamma$ spin polarization exhibits the same behavior with (black curve) and without (red dashed curve) Coulomb interaction. Parameters $\varepsilon^\uparrow/\Gamma = 2$, $\varepsilon^\downarrow/\Gamma = -2$, $\mu^*/\Gamma = 0$, $\gamma_k/\Gamma = 1.5$ and $\Gamma = 1$ are the same for all the figures.

FIG. 4. (Color online) Time evolution of spin polarization degree. Parameters $\varepsilon^\uparrow/\Gamma = 2$, $\varepsilon^\downarrow/\Gamma = -2$, $\mu^*/\Gamma = 0$, $\gamma_k/\Gamma = 1.5$ and $\Gamma = 1$ are the same for all the figures. Colors of the curves and values of parameters $U/\Gamma$ and $\gamma_1/\Gamma$ in the main panel correspond to the colors and values shown in Fig. 3. Insert shows the behavior of black and blue curves at the beginning of the time evolution.
characterized by the tunneling rate, which is taken $\Gamma = 1$. Fig. [2] shows the time evolution of the spin polarization (a) and total number of electrons (b) in the QW with account for the on-site Coulomb repulsion $U$ for the electrons in the bound state. The calculation results are presented for various relaxation rates $\gamma_k$, which describe the radiative recombination in the QW. As shown in Fig. [2] a, the total number of electrons in the QW is decreased as the initial non-equilibrium concentration of electrons relaxes due to the tunneling and radiative recombination. In the absence of Coulomb correlations ($U = 0$), the total relaxation rate for the electrons in the QW is simply the sum of the two $\gamma_{QW} = \gamma_k + \Gamma$. This conclusion holds also for the case $U \neq 0$ if $\gamma_k > \Gamma$. The electrons mostly relax through the recombination channel in the QW and do not have enough time to be affected by the correlations at the bound state. This situation is illustrated by blue and red curves in Fig. [2]. However, if $\Gamma > \gamma_k$ and $U \neq 0$, the decrease of the carriers sheet density in the QW due to the tunneling is delayed as now the tunneling of an electron requires an additional energy cost if the final state is occupied. This case corresponds to the solid and dashed black lines in Fig. [2]

Since the bound state is split in spin projection, the relaxation rate through the tunneling channel is spin-dependent. Therefore, the spin polarization of the electrons in the QW shown in Fig. [2] a increases with time. The increase is linear at $t < (\gamma_k + \Gamma)$ in agreement with Ref.[12], later on the spin polarization decays as QW becomes empty. As can be clearly seen in Fig. [2] a, when the Coulomb correlations at the bound state become important, that is $U \neq 0$, $\Gamma > \gamma_k$, the maximum of the spin polarization is substantially increased. For the parameters used for Fig. [2] the enhancement is more than two times. The position of the maximum on time scale is also substantially shifted to larger times. Thus, the strong Coulomb correlations lead to a stronger dynamic spin injection into the QW and the delayed kinetics, consequently, the spin polarization in the QW is preserved for a longer time.

The effect of the Coulomb correlations on the spin polarization in the QW also depends on the spin-independent relaxation rate at the bound state $\gamma_1$ (assumed to be the same for both spin sublevels). This influence is shown in Fig. [3]. Obviously, if the bound state sublevels are emptied faster than the rate of the incoming tunneling electrons from the QW, the Coulomb correlations shouldn’t play a role as the bound state would never get doubly occupied. Indeed, for $\gamma_1/\Gamma > 1$ the evolution of the total sheet density and the spin polarization in the QW is the same for $U/\Gamma = 30$ (red curve) and for $U = 0$. In the latter case the magnitude of $\gamma_1$ does not matter as occupation of the bound state is not accompanied with an additional energy cost. The effect of the Coulomb correlations becomes important as $\gamma_1$ is enhanced so that the electrons are less effectively removed from the bound state. The spin injection in this case in enhanced as can be clearly seen in Fig. [3] a, blue line. The QW total occupation dynamics is also affected by the Coulomb interactions, which lead to a decrease in the decay rate analogously to what was seen in Fig. [3] b. However, one can note that the discrepancy between different lines in Fig. [3] a develops at times $t > \Gamma$. That is, when the bound state becomes significantly populated with the tunneling electrons so that the correlations become important.

Finally, Fig. [4] shows the spin polarization degree $P$ introduced in Eq. [14]. It is this quantity that can be measured experimentally as a degree of circular polarization of the photoluminescence from the QW. Its time evolution in the presence of the Coulomb correlations is somewhat non-trivial. The linear growth of the spin polarization degree at $0 < t < \Gamma$ is common for the cases with and without on-site Coulomb correlations. Starting from $t = \Gamma$ the increase of $P$ is suppressed by the Coulomb correlations (blue and green lines in Fig. [4]). This is a net effect of the two: the spin polarization, which is the numerator in [14] is enhanced due to an effectively large spin splitting of the bound state but the total occupation of QW, which is the denominator [14] remains larger as the electrons are stuck in the QW. The Coulomb correlations do not manifest themselves if $\gamma_1 \gg \Gamma$ as was discussed above (red line in Fig. [4]). As the total number of non-equilibrium carriers in the system decreases the Coulomb correlations effect on polarization degree vanishes and all the curves converge at larger times in Fig. [4]. The characteristic time evolution of the polarization degree demonstrated in Fig. [4] has been never reported before. In our opinion, it gives a good opportunity to verify the role of the Coulomb correlations in the systems of the considered type experimentally. It is also clear, that the influence of the correlations can be well controlled by adjusting the system parameters. In particular, for the system design shown in Fig. [1] b the bound state relaxation rate $\gamma_1$ is directly related to the transparency of the barrier on the left, which can be tuned by changing the barrier height or its thickness.

V. SUMMARY

We have studied dynamic spin injection by the mechanism of spin-dependent relaxation in a quantum well coupled to the spin-split bound state. In this work for the first time the impact of the Coulomb correlations at the bound state on the spin and sheet density kinetics in the QW were analyzed. As supported by our analysis, the effect of the Coulomb correlations is twofold. Firstly, the on-site Coulomb repulsion leads to an effectively larger spin splitting and, consequently, an enhanced spin polarization of the electrons remaining in the QW. Secondly, it increases the characteristic time of the carriers relaxation in the QW since it reduces the electron tunneling into the bound state. We predict that the interplay of these two effects would lead to the non-trivial dependence of a circular polarization degree of photoluminescence from the QW. This characteristic dependence will
allow probing the strength of the on-site Coulomb correlations experimentally. As shown by our analysis, the effect of the Coulomb correlations can be controlled by affecting the relaxation times. For example, the bound state relaxation time can be tuned by a tunnel barrier separating it from the lead. This opens a way of studying the Hubbard-like electron-electron interactions experimentally.

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1 A. Hoffmann and S. D. Bader, Phys. Rev. Applied 4, 047001 (2015).
2 I. Zutić, J. Fabian, and S. Das Sarma, Rev. Mod. Phys. 76, 323 (2004).
3 J. Sinova and I. Zutić, Nature Materials 11, 368 (2012).
4 C. H. Li, O. M. J. van t Erve, J. T. Robinson, Y. Liu, and L. Li, Nature Nanotechnology 9, 218 (2014).
5 D. Pesin and A. H. MacDonald, Nature Materials 11, 409 (2012).
6 S. Bader and S. Parkin, Annual Review of Condensed Matter Physics 1, 71 (2010).
7 A. Fert and H. Jaffrès, Phys. Rev. B 64, 184420 (2001).
8 M. Oltscher, M. Ciorga, M. Utz, D. Schuh, D. Bougeard, and D. Weiss, Phys. Rev. Lett. 113, 236602 (2014).
9 T. Dietl and H. Ohno, Rev. Mod. Phys. 86, 187 (2014).
10 E. I. Rashba, Phys. Rev. B 62, R16267 (2000).
11 M. Battiato and K. Held, Phys. Rev. Lett. 116, 196601 (2016).
12 E. L. Ivchenko, L. A. Bakaleinikov, and V. K. Kalevich, Phys. Rev. B 91, 205202 (2015).
13 E. L. Ivchenko, L. A. Bakaleinikov, M. M. Afanasiev, and V. K. Kalevich, Physics of the Solid State 58, 1539 (2016).
14 V. Korenev, I. Akimov, S. Zaitsev, V. Sapega, L. Langer, D. Yakovlev, Y. A. Danilov, and M. Bayer, Nat. Commun. 3, 959 (2012).
15 I. Akimov, V. L. Korenev, V. F. Sapega, L. Langer, S. V. Zaitsev, Y. A. Danilov, D. R. Yakovlev, and M. Bayer, physica status solidi (b) 251, 1663 (2014).
16 I. V. Rozhansky, K. S. Denisov, N. S. Averkiev, I. A. Akimov, and E. Lähderanta, Phys. Rev. B 92, 125428 (2015).
17 K. S. Denisov, I. V. Rozhansky, N. S. Averkiev, and E. Lähderanta, Semiconductors 51, 43 (2017).
18 A. Hewson, The Kondo Problem to Heavy Fermions (Cambridge University Press, 1993).
19 N. Maslova, P. Arseyev, and V. Mantsevich, Solid State Communications 248, 21 (2016).
20 J. Q. You and H.-Z. Zheng, Phys. Rev. B 60, 8727 (1999).
21 J. Q. You and H. Z. Zheng, Phys. Rev. B 60, 13314 (1999).