Time-resolved and continuous-wave optical spin pumping of semiconductor quantum wells

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Received 7 July 2008
Published 29 October 2008
Online at stacks.iop.org/SST/23/114001

Abstract

Experimental and theoretical studies of all-optical spin pump and probe of resident electrons in CdTe/(Cd,Mg)Te semiconductor quantum wells are reported. A two-color Hanle-MOKE technique (based on continuous-wave excitation) and time-resolved Kerr rotation in the regime of resonant spin amplification (based on pulsed excitation) provide a complementary measure of electron spin relaxation time. The influence of electron localization on long-lived spin coherence is examined by means of spectral and temperature dependences. Various scenarios of spin polarization generation (via the trion and exciton states) are analyzed and difference between continuous-wave and pulsed excitations is considered. Effects related to inhomogeneous distribution of g-factor and anisotropic spin relaxation time on measured quantities are discussed.

(Some figures in this article are in colour only in the electronic version)

1. Introduction

Optical spin pumping of resident electrons in bulk semiconductors has been established since the 1970s [1]. Circularly polarized light induces interband transitions between the hole states in the valence band (VB) and the electron states (e) in the conduction band (CB). Owing to the optical selection rules the allowed photo-generated electron–hole pairs in the case of \( \sigma^+ \) excitation are \((hh, \pm 3/2); (e, \pm 1/2)\) and \((lh, \pm 1/2); (e, \pm 1/2)\). Here, \((hh, \pm 3/2)\) denotes the heavy-hole state with spin projection \(j_z = \pm 3/2\) on the incident direction of the light, \((lh, \pm 1/2)\) denotes the light-hole state with \(j_z = \pm 1/2\) and \((e, \pm 1/2)\) the electron state with \(s_z = \pm 1/2\). Since the matrix elements of the heavy-hole and light-hole transitions relate as \(\sqrt{3}\) to 1, \(\sigma^+\) excitation creates predominantly spin-down electrons, i.e., with \(s_z = -1/2\) (figure 1). If hole spin flip occurs, the hole can recombine with a spin-up electron emitting \(\sigma^-\)-polarized photon. As a result, the total electron spin \(S_e = (n_e - n_-)/2\) after recombination is not zero any more (\(n_{\pm}\) is concentration of \(s_z = \pm 1/2\) electrons). The net spin polarization is equal to

\[
\rho_e = 2S_e/n_e, \quad (1)
\]

where \(n_e\) is the total electron concentration. Typically, a hole loses its spin quickly and the electron spin relaxation time \(\tau_s\) is rather long. As a consequence, the optically created electron spins are accumulated in \(s_z = -1/2\) state under continuous wave (cw) excitation with \(\sigma^+\) polarization. This implies optical spin pumping. The theoretical limit of the steady-state net spin polarization is \(\rho_e = 0.5\) [1]. This is valid for bulk semiconductors with degenerate light-hole and heavy-hole states, like GaAs and CdTe.

In bulk \(n\)-type semiconductors the circular polarization degree of the edge photoluminescence (PL) \(P = (I^+ - I^-)/(I^+ + I^-)\) is proportional to the net spin polarization \(P = \rho_e/2\). Here, \(I^\pm\) is the PL intensity of...
The corresponding optical transition under $\sigma^\pm$ excitation detected in $\sigma^-$polarization component. When an in-plane magnetic field $B$ is applied electron spins start to precess around the magnetic field direction with the Larmor frequency $\omega_L = g_e \mu_B B / \hbar$, where $g_e$ is the electron $g$-factor and $\mu_B$ is the Bohr magneton. As a result, $S_z$ decreases following

$$S_z (B) = \frac{S_z (0)}{1 + (B / B_{L/2})^2}.$$ (2)

Such behavior detected by means of the PL circular polarization $P \propto S_z$ is known as the Hanle effect [1]. The average electron spin drops by a factor of 2 when the electron spin lifetime $\tau_s$ and the Larmor frequency $\omega_L$ satisfy the condition $\omega_L \tau_s = 1$. This condition defines characteristic magnetic field $B_{L/2} = \hbar / (\mu_B g_e \tau_s)$ allowing us to obtain $\tau_s$ if the $g$-factor is known. These experiments are performed under cw optical excitation.

In $n$-type semiconductors an additional spin decay channel is the recombinations of resident electrons with photogenerated holes. As a consequence the spin lifetime depends on pump density $W$ as

$$\tau_s^{-1} = \tau_s^{-1} (1 + W / W_0).$$ (3)

Using equation (2) in the limit of low excitation density (i.e., $W$ much below the characteristic pump density $W_0$) one can find the spin relaxation time of the ‘unperturbed’ electron system $\tau_s$, as $\tau_s \rightarrow \tau_s$.

Time-resolved spectroscopy based on pulsed photoexcitation offers straightforward access to the carrier spin dynamics (for overview see [2]). For example, polarized photoluminescence measured by streak camera allows us to follow spin relaxation during the exciton lifetime. Faraday and Kerr rotation techniques can monitor magnetization induced by spin polarized carriers at time delays not limited by exciton recombination. In external magnetic fields not only spin dynamics but also Zeeman splitting for carriers and excitons can be measured from the period of signal oscillations. Respectively, the direct evaluation of carrier and exciton $g$-factors became possible. Partial limitations of time-resolved techniques are related to relatively high peak density of photoexcitation as compared to cw excitation. This may cause perturbation of the electronic system (carrier heating, nonequilibrium distribution) which may affect the spin dynamics.

In low-dimensional quantum well (QW) heterostructures the optical spin pumping can be modified for following reasons. (i) Quantum confinement and strain lift up the degeneracy of light-hole and heavy-hole states. Hence, their selective excitation becomes possible and the theoretical limit for the net spin polarization in this case is increased to 100% ($\rho_z = 1$). (ii) In $n$-type modulation doped QWs a negatively charged trion ($T$) became stable [3]. It is composed of a hole and two electrons, the latter form the spin singlet. The singlet state itself does not contribute to the spin polarization, but resonant excitation of trions and/or trion formation from excitons provide an efficient mechanism for polarization of resident electrons in QWs [4].

The goal of this paper is to perform a comparative study of electron spin dynamics under cw and pulsed photoexcitation and to clarify whether the same values of spin relaxation times can be achieved by both techniques. Experiments were done for the very same samples, which are CdTe/CdMgTe quantum well structures with diluted concentration of resident electrons.

The paper is organized as follows. We start with model considerations of the generation of spin polarization and spin coherence in quantum wells with diluted two-dimensional electron gas. Then we present experimental results for cw and pulsed excitation, analyze them and conclude.

2. Theory

In the present section different scenarios of spin polarization generation and spin dynamics in transverse magnetic field are discussed. We consider both continuous wave (cw) and pulsed excitation regimes and we discuss one after another the mechanisms of spin pumping under excitation at the trion and exciton resonances.

2.1. Spin coherence excitation

Here we analyze the processes responsible for the generation of the electron spin coherence in QWs with low density two-dimensional electron gas (2DEG). It is assumed that the temperature measured in energy units, $k_B T$, is smaller than both exciton and trion binding energies and that the resident electron concentration is so small that trion and exciton states are stable (i.e. these states are not screened by electron–electron interactions). More detailed analysis is published in [4].

2.1.1. Trion resonant pumping

Continuous wave excitation. First, we consider the case where the trion is resonantly excited. The trion ground state forms the singlet spin structure, i.e., the spins of the two electrons involved in the trion are oriented antiparallel [5, 6]. A $\sigma^+$ polarized photon creates an electron–hole pair ($hh, +3/2; e, -1/2$) and simultaneously an electron with the spin projection $s_z = +1/2$ is picked out from the resident ensemble in order to form a trion (figure 2). Thus, the resident ensemble becomes depleted of electrons with the spin projection opposite to that of photocreates ones, i.e. in our
case the number of spin \( s_z = 1/2 \) electrons decreases. This means that a spin polarization of resident carriers \( S_z \neq 0 \) appears. The electrons returning to the ensemble after trion recombination can compensate this polarization. However, if the hole spin relaxation time in the trion is much shorter than the trion radiative lifetime the returned electrons are spin unpolarized and the spin polarization generated at the moment of absorption is conserved.

From the considerations above it is clear that at the initial moment each absorbed \( \sigma^+ \) photon increases the spin of resident electrons by \( \mp 1/2 \). In order to study the pump power dependence of the spin polarization generation we use the system of coupled kinetic equations describing the populations of electrons with different spins \( n_{\sigma} \), and populations of singlet trions \( T_{\sigma} \) (with the heavy-hole spin \( \pm 3/2 \))

\[
\begin{align*}
\frac{d}{dt} T_+ &= \frac{T_+}{\tau_0} - \frac{T_+}{2 \tau_s^T} + n_+ G = 0, \\
\frac{d}{dt} T_- &= \frac{T_-}{\tau_0} - \frac{T_-}{2 \tau_s^T} + n_- G = 0, \\
\frac{d}{dt} n_+ &= \frac{n_+}{\tau_s} + \frac{n_+}{2 \tau_s} - \frac{T_+}{\tau_0} - n_+ G = 0, \\
\frac{d}{dt} n_- &= \frac{n_-}{\tau_s} + \frac{n_-}{2 \tau_s} - \frac{T_-}{\tau_0} - n_- G = 0.
\end{align*}
\]

Here it is assumed that the pumping is \( \sigma^+ \) polarized, \( G \) is the generation rate (being proportional to the pump density \( W \)), \( \tau_0^T \) is the trion radiative lifetime, \( \tau_s^T \) is the spin relaxation time of a hole in a trion (trion spin relaxation time) and \( \tau_s \) is the spin relaxation time of resident electrons. System (4) should be complemented with the condition

\[
T_+ + T_- + n_+ + n_- = n_e,
\]

which describes the conservation of resident electrons \( n_e \).

The total steady-state spin of resident electrons can be written as

\[
S_z = \frac{n_+ - n_-}{2} = -\frac{n_e}{2} \frac{G \tau_s}{\tau_s + \tau_s^T + \tau_T + \tau_0^T}.
\]

It is worth noting that the absolute value of the electron spin \( |S_z| \) decreases with an increase of trion spin relaxation time \( \tau_s^T \). It is in agreement with the qualitative picture outlined above: in order to induce the spin coherence of resident electrons the spin flip of a hole in a trion is required. That means that \( \tau_s^T \) should be comparable to or shorter than the trion lifetime \( \tau_0^T = 50 \) ps. With further increase of the pump density the total spin of resident electrons grows sublinearly (figure 3). This is because the larger \( |S_z| \) is, the fewer resident electrons with proper spin required to form the trion (i.e., \( s_z = 1/2 \)) are left.

At high pump densities \( G \tau_s \gg \tau_s^T / \tau_T \), i.e., where the spin polarization is small and saturation effects are observed the electron spin grows linearly with the pumping density as

\[
S_z = -\frac{n_e}{4} \frac{G \tau_s \tau_T}{\tau_s^T / \tau_T - \tau_0^T}.
\]

This means that almost all resident electrons become spin polarized: \( p_e \approx 1 \). The negative sign results from the \( \sigma^+ \) pumping when electrons with spin \( s_z = -1/2 \) are created. The minority of electrons \( s_z = +1/2 \) are bound in trions whose density is \( n_e (\tau_0^T / \tau_s^T) \). This implies that efficient spin pumping of resident electrons may be achieved at relatively low excitation densities. Such unexpected behavior is a direct consequence of the long spin relaxation of resident carriers. With increasing pump density the spin-up electrons

\[\text{Figure 2. Schematic illustration of the optical spin pumping of resident electrons in QWs, resonant excitation to the trion (T).}\]

\[\text{Figure 3. (a) Spin polarization of resident electrons } p_e \text{ as a function of the generation rate in units of } G \tau_s \text{ under } \sigma^+ \text{ excitation. Dashed curve corresponds to the trion resonant pumping. Solid curves correspond to the exciton resonant pumping. They are labeled by the respective electron-in-exciton spin relaxation time, } \tau_s^T. \text{ All curves are obtained for the electron trion relaxation time } \tau_e = 2 \text{ ns, the exciton and trion radiative lifetimes } \tau_0^T = \tau_0^X = 50 \) \text{ ps, the spin relaxation time of a hole in a trion } \tau_s^T = 10 \) \text{ ps. The exciton to trion conversion coefficient is } \gamma n_e = 1 \) \text{ ps. (b) The spin polarization of resident electrons (under the exciton resonant pumping) shown for a wide range of } G \tau_s \text{ (for } \tau_s^T = 10 \) \text{ ps).}\]
(sz = +1/2) are captured to the trions and return back unpolarized. The stronger the pumping, the fewer electrons have appropriate spin orientation to form a trion and the number of trions decreases.

**Pulsed excitation.** Under pulsed excitation the physical processes governing electron polarization generation are essentially the same as for cw excitation [4]. After the trion formation the electron gas becomes depleted of electrons with the spin projection opposite to that of photogenerated ones. The spin flip of a hole in a trion results in the imbalance of the spins of resident electrons and those returning after trion radiative recombination.

It is worth noting that the initial number of photogenerated trions under pulsed resonant excitation (nT0) cannot exceed ne/2, where ne is the density of resident electrons. Thus, nT0 increases linearly with the pump intensity for small excitation densities and then saturates at the value ne/2. The total generated electron spin at τT0 ≫ τs can be evaluated as

\[ S_z = -\frac{n_e}{4} G \tau_0 T / (1 + G \tau_0 T), \]  

where G is the generation rate, proportional to the pump power.

It is instructive to compare this result with the result of equation (8) obtained for the steady-state pumping. In the case of pulsed excitation the maximum total spin of resident electrons is -n_e/4, i.e., twice as small than the steady-state value.

2.1.2. **Exciton resonant pumping.** **Continuous wave excitation.** Next, we discuss the case of exciton resonant pumping. Photocreated excitons can capture resident electrons and form trions. This process is schematically shown in figure 4. Under the assumption that the hole in the exciton loses its spin rapidly one may label concentrations of X, σ+, and σ- excitons in accordance with the electron spin projection in the exciton. The trions formed from these excitons are unpolarized because the spin of the paired electrons in the trion is zero and the hole spin is lost prior to trion formation. The kinetic equations describing such a situation are written

\[ \frac{X_+}{\tau_0} - \frac{X_+}{2\tau_s} + \frac{X_+}{2\tau_X} = \gamma n_e X_+ + G_X = 0, \]

\[ \frac{X_-}{\tau_0} - \frac{X_-}{2\tau_s} + \frac{X_-}{2\tau_X} = \gamma n_e X_- + G_X = 0, \]

\[ \frac{T}{\tau_0} + \gamma n_e X_+ + \gamma n_e X_- = 0, \]

\[ \frac{T}{\tau_0} + \gamma n_e X_+ + \gamma n_e X_- = 0. \]

\[ \frac{n_+}{2\tau_s} + \frac{n_-}{2\tau_s} + \frac{T}{2\tau_0} - \gamma n_e X_+ = 0, \]

\[ \frac{n_+}{2\tau_s} + \frac{n_-}{2\tau_s} + \frac{T}{2\tau_0} - \gamma n_e X_- = 0. \]  

Here τX is the exciton radiative lifetime, T is the total concentration of excitons, and G_X is the exciton generation rate (σ+ polarized pumping is assumed).

Solid lines in figure 3 show the spin of resident electrons as a function of the pump density calculated for various values of G_X. Contrary to the case of trion resonant excitation, an initial increase of the net spin polarization is followed by its decrease to zero for very high pump densities (shown in panel b). Qualitatively, such behavior can be explained as follows. In the absence of electron spin relaxation, the formation of trions is prevented if the exciton concentration exceeds half of the electron concentration. However, the spin flip on an electron-in-exciton can provide the formation of trion even for higher exciton concentrations (schematically shown in figure 5). This is accompanied by a decrease of the net spin polarization of excitons [4]. At very high pumping densities, G_X τ_e/ n_e ≫ 1, all resident electrons become bound to unpolarized trions and the system will consist of n_e unpolarized trions and X excitons (no unbound excitons) and X = G_X τ_e/ n_e excitons with the total electron-in-exciton spin (X_+ - X_-)/2 = G_X τ_e/ n_e τ_s/ (τ_0 + τ_s^X).

At low pump densities most excitons form trions (the trion formation process can be considered as the fastest one when γ n_e/2 ≫ τ_0^X, τ_s^X), which results in the spin polarization of resident excitons. At G_X τ_e/ n_e ≪ 1

\[ S_z = -\frac{1}{2} G_X \tau_s, \]  

and each absorbed photon participates in the exciton formation either directly or via the exciton state. Therefore, the efficiency of the electronic spin polarization is expected to be nearly independent of the pump energy. Comparing equation (11) and equation (7) one can build a relation between the trion generation rate G entering equation (4) and exciton generation rate in equation (10) as G_X = n_e G/2. The spin polarization degree at high pump powers is predicted to be strongly dependent on the excitation energy and demonstrate qualitatively different behavior.

**Pulsed excitation.** In time-resolved experiments at low pump densities shortly after the pulsed optical excitation all excitons are bound to trions and the QW contains n_T0 trions (where n_T0 =
\[ G \times \tau_{\text{pulse}} \text{ with } \tau_{\text{pulse}} \text{ being the pulse duration, } \tau_{\text{pulse}} \ll \tau_X^0, \text{ is the number of photogenerated excitons and } n_e = n_0^X \text{ resident electrons with a total spin} \]

\[ |S_z| = n_0^X / 2, \quad n_0^X \leq n_e / 2. \]  

(12)

As a result, \( n_0^X \) spins of resident electrons contribute to the Kerr rotation.

At higher excitation intensity \( \left(n_0^X \geq n_e / 2 \right) \) all \( n_e / 2 \) resident electrons with \( s_e = +1/2 \) are bound to trions. Therefore, in the absence of electron-in-exciton spin relaxation processes the trion density cannot exceed \( n_e / 2 \), thus the total spin density of the electron gas is limited by \( |S_z| < n_e / 4 \). The electron-in-exciton spin relaxation allows us to convert the remaining \( n_0^X = (n_e / 2) \) excitons into trions. Obviously, the maximum number of formed trions cannot exceed the concentration of background electrons, \( n_e \). The total spin of resident electrons after the excitons and trions have recombined can be estimated as (provided that the holes are unpolarized)

\[ |S_z| \approx \frac{1}{4} \left( n_e - \frac{2n_0^X - n_e}{1 + (2\tau_e^X / \tau_0^X)} , \quad n_e > \frac{2n_0^X - n_e}{1 + (2\tau_e^X / \tau_0^X)} \right) \quad \text{otherwise}. \]  

(13)

This equation is valid both for \( B = 0 \) and \( B \neq 0 \) when \( n_0^X \geq n_e / 2 \), otherwise equation (12) holds. At \( n_0^X = n_e / 2 \), the values of \( |S_z| \) given by equations (12) and (13) coincide and are equal to \( n_e / 4 \). An initial linear increase of \( |S_z| \) followed by a decrease is seen from equation (13). The decrease of initial electron spin as a function of pump intensity is steeper for smaller values of \( \tau_e^X / \tau_0^X \), i.e. for shorter hole spin relaxation times. It is worth stressing that in this regime the electron spin polarization vanishes at very high pumping whereas, under resonant trion excitation, \( |S_z| \) monotonically increases with increasing pump power and saturates at \( n_e / 4 \).

### 2.2. Spin dynamics in transverse magnetic fields

Application of a transverse magnetic field (i.e., perpendicular to the initial spin direction) induces electron spin precession about the magnetic field, which is often called spin beats. In the general case time evolution of \( S_z \) is described by

\[ S_z(t) \sim e^{-i/T^z} \cos(\tilde{\Omega}t + \psi), \]  

(14)

where \( \psi \) is an initial phase, \( T^z \) is the dephasing time of electron spin ensemble. The dephasing time is contributed by the coherence time of individual spins \( T^z_2 \) and by inhomogeneous spin relaxation time \( T^{z\text{inh}}_2 \) caused by, e.g., variation in electron g-factors: \( 1/T^z_2 = 1/T^z_2 + 1/T^{z\text{inh}}_2 \) (see chapter 6 in [2]). In the simplest case of isotropic and homogeneous systems the spin relaxation of resident electrons is characterized by a single time constant \( \tau_s = T^z_2 \). In this case \( T^z_2 = \tau_s \), and the frequency of spin beats is equal to the Larmor frequency \( \tilde{\Omega} = \omega_L \). In the case of QWs experimental behavior appears to be somewhat more complicated for few reasons: (i) spin relaxation can be anisotropic [7], (ii) the main mechanism of spin dephasing can be caused by g-factor inhomogeneous distribution, (iii) the initial phase can depend on the magnetic field.

#### 2.2.1. Continuous wave excitation

In many cases the regime of cw spin pumping in transverse magnetic fields (and the resultant Hanle effect) may be described in terms of the spin relaxation time \( \tau_s \). The corrections due to anisotropy and inhomogeneity of structures are discussed in section 2.2.2.

**Trion resonant pumping.** The system of kinetic equations describing spin dynamics of resident electrons and trions at low pump densities under trion resonant excitation has the following form [4]:

\[ \frac{dS_x}{dt} = \omega_L S_x - \frac{S_x}{\tau_s} + \frac{n_e}{4} G(t), \]

\[ \frac{dS_y}{dt} = -\omega_L S_y - \frac{S_y}{\tau_s}, \]

\[ \frac{dS_z}{dt} = \frac{S_y}{\tau_f} + \frac{n_e}{4} G(t), \]  

(15)

where \( \tau_f = (T_s - \tau) / 2 \) is the effective spin density of the trions, \( S_z \) and \( S_y \) are the electron spin density components. Function \( G(t) \) in equation (15) describes the trion generation rate. It is assumed that the \( z \)-axis coincides with the growth direction of the QW structure and the magnetic field is applied along the \( x \)-axis. The solution for \( S_z \), \( S_y \) and \( S_T \) in the steady-state \( (G(t) \text{ is a constant}) \) is readily written

\[ S_z = -\frac{n_e G \tau_T^z}{4 \tau_f^z + 1 + (\omega_L \tau_f)^2}, \]

\[ S_y = -S_z \omega_L \tau_s, \]

\[ S_T = \frac{n_e}{4} G \tau_L^z. \]

(16)

Field dependence of \( S_z \) is described by the Lorentz curve, in agreement with equation (2). The first line of equation (16) at \( \omega_L = 0 \) agrees with equation (7). In the absence of trion spin relaxation \( (\tau_f^z \rightarrow \infty) \) the application of an in-plane magnetic field does not result in steady-state spin polarization. This is in contrast with time-resolved experiments, where spin beats could be observed even for \( \tau_f^z \rightarrow \infty \) (see figure 6 in section 2.2.2 and [4]).

Typically, the trion spin relaxation time is much shorter than the electron spin relaxation time, \( \tau_f^z \ll \tau_s \). Therefore, in moderate magnetic fields \( \omega_L \tau_s \sim 1 \) the steady-state spin of electrons significantly exceeds the trion spin, \( |S_z| \gg |S_T| \).

**Exciton resonant pumping.** In this case the total spin of both resident electrons \( S_z \) and electrons in excitons \( S_X = (X_e - X_-)/2 \) should be considered. Using equations (10) in the rotating (with the frequency \( \omega_L \)) frame of reference, one can derive the following equations describing the steady-state spin of electrons and excitons in the weak pumping regime \( (G_X \tau_s \ll 1) \):

\[ \omega_L \times S_X = \frac{S_X}{T^X_2} = \frac{1}{2} G_X, \]

\[ \omega_L \times S = \frac{S_X}{\tau_s} + \gamma n_e = 0, \]

(17)

where \( S_X \) and \( S \) are the vectors of spin of electrons-in-excitons and of resident electrons, respectively, \( G_X \) is a vector directed along the \( z \)-axis with an absolute value \( G_X \) and \( T_{\gamma}^{-1} = \left( T_0^X \right)^{-1} + \left( \frac{T_0^X}{2} \right)^{-1} + \gamma n_e / 2 \) is the lifetime of electron-in-exciton spin.
The reason is schematically illustrated in figure 6. The initial phase of spin beats in the case of trion resonant excitation strongly depends on magnetic fields. Initial phase.

After pulsed excitation some of the resident electrons with a certain spin projection are picked out to form trions. In the absence of hole spin relaxation the returning electrons after trion radiative recombination exactly compensate the generated spin of resident carriers. The transverse magnetic field breaks such a balance: the spins of electrons in trions do not precess while the spins of resident electrons rotate around the magnetic field, therefore the compensation is not complete. Detailed analysis shows that the electron spin as a function of time can be recast as [4]

\[ S_z(t) = \frac{n^T_0}{2} [1 - \eta \sin(\omega_0 t - \varphi) e^{-t/\tau_s} - \eta' e^{-t/\tau_T}] \]

where \( \eta' \) and \( \eta'' \) are the real and imaginary parts of \( \eta = (\tau_0^T)^{-1}/[(\tau_T^T)^{-1} - \tau_s^{-1} - i\omega_0] \). Hence, the initial phase is a non-monotonic function of the magnetic field.

**Inhomogeneous distribution of electron g-factor.** In order to take into account inhomogeneous broadening we assume that electron states are localized and characterized by a distribution of g-factors, \( f(g) \). Hopping between different states is neglected, therefore it suffices to average the single electron spin precession \( \exp(-t/\tau_s) \cos(\omega_0 t + \varphi) \) over this distribution.

In the case of cw pumping the Hanle curve, (16), should be averaged over \( f(g) \) as well.

It is possible to obtain the analytical results when \( f(g) \) is described by the Lorentzian function

\[ f(g) = \frac{\alpha}{\pi[(g - g_0)^2 + \sigma^2]} \]  

where \( g_0 \) is the average value of g-factor and \( \sigma \) is the distribution width. Instead of (16) we have for \( \langle S_z \rangle \)

\[ \langle S_z \rangle \sim \frac{1 + \sigma \mu_B B \tau_s/\hbar}{(1 + \sigma \mu_B B \tau_s/\hbar)^2 + (g_0 \mu_B B \tau_s/\hbar)^2} \]

We note that equation (22) deviates from the standard Hanle curve in the magnetic field \( B \) where \( \sigma \mu_B B \tau_s/\hbar \) becomes of the order of unity. Typically, in QWs the g-factor spread is small, so that \( \sigma \ll g_0 \). As a result, inhomogeneous broadening does not strongly modify the Hanle curves in the region of \( g_e \mu_B B \tau_s/\hbar \sim 1 \).

The influence of the inhomogeneous distribution of g-factors on spin beats is different. They can still be observed in high magnetic fields when \( g_e \mu_B B \tau_s/\hbar \gg 1 \). We obtain averaging with equation (21)

\[ \langle S_z(t) \rangle \sim \exp(-\sigma \mu_B B t/\hbar - t/\tau_s) \cos(\omega_0 t + \varphi) \]  

Comparison with equation (14) suggests that inhomogeneous broadening of the electron g-factor results in an additional spin decay channel characterized by a time constant \( T_2^{inh}(B) = (\sigma \mu_B B / \hbar)^{-1} \). In strong enough magnetic fields (provided \( T_2^{inh}(B) \ll \tau_s \)) this spin decay channel dominates the spin beat dephasing process \( T_2^* \approx T_2^{inh} \).

**Anisotropy effects.** In semiconductor quantum wells spin relaxation is known to be anisotropic, see [7] and references therein. In this case the inverse spin relaxation times are described by a tensor. In the general case of an asymmetric quantum well there are three linearly-independent components \( \tau_{xx}, \tau_{yy} \) and \( \tau_{zz} \), where \( x||[110], y||[110] \) and \( z||[001] \).
spin relaxation time $T_2$ and the precession frequency $\hat{\Omega}$ entering equation (14) are then written as follows [8, 9]:

$$\frac{1}{T_2} = \frac{1}{2} \left( \frac{1}{\tau_{zz}} + \frac{1}{\tau_{yy}} \right), \quad \hat{\Omega} = \sqrt{\omega_L^2 - \frac{1}{4} \left( \frac{1}{\tau_{zz}} - \frac{1}{\tau_{yy}} \right)^2}.$$  

Note that the magnetic field is directed along the $x$-axis. The Hanle effect is described by a Lorentzian $\sim [1 + (\omega/T_H)^2]^{-1}$, where $T_H = \sqrt{\tau_{zz}/\tau_{yy}}$ [7]. Thus, characteristic times $\tau_{yy}$ and $\tau_{zz}$ can be found:

$$\tau_{zz} = \frac{T_2}{2} - \sqrt{\frac{2}{T_2} \frac{T_4}{T_2^2} - 2}, \quad \tau_{yy} = \frac{T_2}{2} + \sqrt{\frac{2}{T_2} \frac{T_4}{T_2^2} - 2}.$$  

(25)

Here it is assumed that $\tau_{zz} < \tau_{yy}$, otherwise signs before square roots should be reversed. The rotating magnetic field in the QW plane allows us to restore all components of the spin relaxation tensor.

### 2.2.3. Resonant spin amplification

In time-resolved experiments periodic trains of laser pulses with repetition period of about $T_{\text{rep}} = 12$ ns are commonly used. In the case where the spin dephasing time is longer than $T_{\text{rep}}$ the signal does not fully decay from the previous pulse and overlaps with the signal generated by the following pulse. Depending on the magnetic field strength it may cause constructive or destructive interference of these contributions, which complicates the evaluation of the dephasing times. To overcome this complication the technique of resonant spin amplification has been suggested [10].

In resonant spin amplification (RSA) experiments the magnetic field dependence of the electron spin is analyzed at a fixed pump–probe delay. If the pulse repetition period is commensurable with the spin precession in the magnetic field the spin polarization is enhanced because the spin is injected in the system at the moments when the precessing spin of resident carriers is parallel to the spin of photocreated electrons.

The electron spin polarization can be written as

$$S_e(\Delta t) = \frac{S_0}{2} e^{-(t_{\text{rep}} + \Delta t)/T_2} e^{T_{\text{rep}}/T_2} \frac{\cos(\Omega T_{\text{rep}})}{\cos(\Omega T_{\text{rep}})} C(2\Omega(\Delta t)),$$

where $\Delta t \in [-T_{\text{rep}}, 0]$ is the delay between the probe pulse and the following pump pulse,

$$C(x) = \cos x - \frac{1}{2\Omega} \left( \frac{1}{\tau_{zz}} - \frac{1}{\tau_{yy}} \right) \sin x.$$  

(26)

According to equation (26) and figure 7 the electron spin polarization represents a series of sharp peaks as a function of the magnetic field. These peaks correspond to the commensurability of the spin precession period and the pulse repetition period. The analysis shows that the ratio of the zero-field maximum to the next maxima (with not too large numbers $N \ll T_{\text{rep}}/|\Delta t|$) is $\eta = (e^{T_{\text{rep}}/T_2} - 1)(e^{T_{\text{rep}}/T_2} - 1)$ [9]. It is equal to 1 in the case of isotropic spin relaxation, $\tau_{zz} = \tau_{yy}$, and smaller than 1 if $\tau_{zz} < \tau_{yy}$.

The inhomogeneous distribution of $g$-factor values can also affect the resonant spin amplification. As can be shown [9], the shape of the zero-field peak is not affected by the $g$-factor distribution. However, peaks with larger numbers become wider and smaller as the effective spin decay constant $T_2^* \approx T_2^{\text{inh}}(B)$ decreases with $B$ (see equation (23)).

### 3. Experiment

#### 3.1. Samples

We present characteristic experimental results based on two representative samples of CdTe/Cd0.7Mg0.3Te QWs grown by molecular beam epitaxy on (0 0 1)-oriented GaAs substrates [11].

Sample 1 is a single 80 Å CdTe/Cd0.7Mg0.3Te QW (the PL spectrum is shown in figure 8(a)). Free electrons in the QW...
are provided due to modulation doping by iodine donors in the Cd$_{0.7}$Mg$_{0.3}$Te barrier at a distance of 100 Å from the CdTe QW. An electron density in the QW of $n_e = 8 \times 10^{10}$ cm$^{-2}$ was evaluated by optical means [12].

Sample 2 is a multiple quantum well (MQW) structure (the PL spectrum is shown in figure 8(b)). It contains five 200 Å wide CdTe QWs separated by 1100 Å thick Cd$_{0.78}$Mg$_{0.22}$Te barriers. The sample is nominally undoped. The low concentrations of electrons $n_e = 1.3 \times 10^{10}$ cm$^{-2}$ in the CdTe QWs are due to residual $n$-type doping of barriers.

### 3.2. Experimental techniques

#### 3.2.1. Continuous wave excitation

For cw pumping, a tunable dye laser is used. Its radiation wavelength can be set in resonance with either the trion ($T$) or exciton ($X$) states (or one of some other high energy states) of a QW sample. The excitation is modulated between ($\sigma^+$ or one of some other high energy states) of a QW sample.

**Hanle effect.** The degree of circular polarization of the photoluminescence (PL) $P = (I_{\sigma^+} - I_{\sigma^-})/(I_{\sigma^+} + I_{\sigma^-})$ is detected by a Si-based avalanche photodiode and a two-channel photon counter. Here, $I_{\sigma^\pm}$ is the PL intensity under $\sigma^\pm$ excitation detected in $\sigma^\pm$ circular polarization component. PL spectra are dispersed by a 1 m monochromator. The degree of circular polarization can be selectively detected on either the trion or exciton emission lines, denoted by $P_T$ and $P_X$, respectively. When the external magnetic field is applied in the QW plane $P$ is suppressed to zero (Hanle effect). We note, the circular polarization in emission is sensitive to the polarizations of resident electrons ($\rho_e$) and excitons ($\rho_T$).

**Two-color Hanle-MOKE.** Alternatively, the net spin polarization of resident electrons can be probed by the magneto-optical Kerr effect (MOKE) [13–15]. In this experiment, one tunable dye laser is used for spin pumping and another tunable Ti:sapphire laser is used to probe spins. The latter is linearly polarized, and the photoinduced Kerr rotation $\theta$ is measured by a balanced diode detector and demodulated by a lock-in amplifier. Pump and probe energies ($E_{\text{pump}}$ and $E_{\text{probe}}$) can be tuned independently, referred to as the two-color mode [15]. By analogy with the Hanle effect, when external magnetic fields are applied in the QW plane $\theta$ is suppressed to zero. Also note, in the cw pumping regime concentrations of excitons and trions are much less than the concentration of resident electrons. Therefore, the photoinduced Kerr rotation is proportional to the total spin of resident electrons only, $\theta \propto S_z$ (i.e., contributions from electrons in excitons and trions can be neglected).

#### 3.2.2. Pulsed excitation

Complementary, coherent spin dynamics can be studied by means of a time-resolved pump–probe Kerr rotation technique [2, 16]. It allows direct monitoring of the evolution of the spin coherence of carriers generated by the pump pulses.

**Time-resolved Kerr rotation (TRKR).** A Ti:sapphire laser generates 1.5 ps pulses at a repetition frequency of 75.6 MHz. The spectral width of pulses is 1.5 meV allowing us to selectively excite different resonances in QWs. The laser beam is split into pump and probe beams (one color mode) and the time delay $\Delta t$ between them is varied by a mechanical delay line. The pump beam is circularly polarized by means of a photoelastic modulator operated at 50 kHz. The probe beam was linearly polarized, and the Kerr rotation $\theta$ is measured by a balanced photodetector. From an analysis of the decay of the Kerr amplitude $\theta(t)$ the spin dephasing time of electrons $T_2^*$ can be extracted [17, 18]. And the precession frequency of $\theta(t)$ allows us to determine the electron $g$-factor.

The TRKR technique can be further improved by using two synchronized Ti:sapphire lasers [4]. This allows us to tune independently the pump and probe energies (two-color mode).

Note, for short delay times (in the case of pulsed excitation) the concentration of photocreated excitons can be comparable with the concentration of resident electrons. Therefore, the Kerr signal is proportional to the total spin of electrons-in-excitons ($S_X$) and of resident electrons ($S_e$). This is in contrast to the cw Kerr signal, where contribution from excitons can be neglected.

**Resonant spin amplification (RSA).** It allows us to measure $T_2^*$ times in the zero magnetic field limit [10, 17]. In this method an external magnetic field is scanned from small negative ($\sim 20$ mT) to small positive fields ($+20$ mT) and a small negative time delay of the probe pulse is chosen ($\Delta t = -80$ ps). The RSA is very suitable for the study of long lived spin beats, when the Kerr signal has considerable amplitude at negative delays and interferes with the signal at positive delays.

### 3.3. Experimental results: continuous wave excitation

#### 3.3.1. Hanle effect

The typical photoluminescence (PL) spectrum of sample 1 is presented in figure 8(a). It consists of two lines separated by 3.8 meV. The high-energy line is attributed to the neutral exciton ($X$). The low-energy line originates from the negatively charged trion ($T$) [5], consisting of a hole and two electrons in the singlet state [19].

A Hanle curve detected on the exciton line (pump energy $E_{\text{pump}} = 1.648$ eV, $W = 2$ W cm$^{-2}$) is shown in figure 9(a). It has two regions: (i) in the high-field region the polarization decreases slowly from 5.5% down to 3.5%. (ii) In the low-field region $|B| < 3$ mT a sharp peak is observed, where the polarization drops from ca 6% down to 5.5%. The origin of the complex line shape of the Hanle curves in QWs has been qualitatively understood [20]. The point is that the trion is formed from the exciton and a resident electron $X + e \rightarrow T$, and this formation appears to be spin dependent. Hence, the degree of circular polarization detected on the exciton ($P_X$) and trion ($P_T$) emission lines can be written as [20]

$$P_X = \rho_X - \gamma \rho_e, \quad P_T = \rho_X + \rho_e, \quad (27)$$
where $\rho_X \propto S_X + S_X^{\perp}$ and $\rho_S \propto S_S^z$ are spin polarizations of excitons and resident electrons, respectively, and $\tilde{g}$ is a constant describing the efficiency of trion formation from an exciton. The total resident electron and electron-in-exciton spins ($S_S$ and $S_X^{\perp}$) are introduced in section 2.2.1. The total hole-in-exciton spin denoted by $S_X^{\perp}$ results in incomplete depolarization of the exciton, as it is not affected by in-plane magnetic fields (in-plane hole $g$-factor is close to zero). The magnetic field dependences of $S_X$ and $S_S$ are described by equation (18). In the limit $T_X / \tau_S \ll 1$, depolarization of each of them follows the standard Hanle curve of equation (2) but with different characteristic magnetic field $B_1/2$.

The best fit to equations (27) and (18) is shown by the solid line in figure 9(a). From this fit we find for region (i) $B_1/2 \approx 280 \text{ mT}$. It is ascribed to the depolarization of an electron within the exciton. Due to the electron–hole exchange interaction (enhanced within the exciton) the electron spin relaxation time $t_{\tau}^s$ shortens following the hole spin flips characterized by a time constant $\tau_{\tau}^h$. This provides in the case of the strong electron–hole exchange interaction the width of the exciton Hanle curve $B_1/2 = \sqrt{\tau_{\tau}^h / \tau_0^s} \Delta_0 / |g_e| \mu_B$, where $\Delta_0$ is the electron–hole exchange splitting [21].

The regime (ii) in figure 9(a) is ascribed to the depolarization of resident electrons [20, 22]. It is measured with higher resolution for excitation density $W = 18 \text{ W cm}^{-2}$ as presented in figure 9(b). Fitting to equation (2) (solid line) gives a characteristic magnetic field $B_1/2 \approx 0.75 \text{ mT}$. We find that $B_1/2$ depends on the exciton density in accord with equation (3) [22]. Extrapolation of $B_1/2$ to zero excitation density ($W \to 0$) allows us to find the electron spin relaxation time $\tau_{\tau} = 19 \text{ ns}$ using the electron $g$-factor $|g_e| = 1.35$ (deduced from the spin beats as discussed is section 3.4.1).

3.3.2. Two-color MOKE. We note that the contribution of the net spin polarization of resident electrons in the exciton or trion emission may be rather weak and difficult to detect. As an alternative, the two-color MOKE offers a highly sensitive technique. We found that the Kerr angle obtained in the cw regime is maximum when the pump energy coincides or is slightly above and the probe energy is slightly below the $X$ or $T$ resonance [23]. Hence, independent tuning of pump and probe energies is an essential requirement for the two-color MOKE method. A typical dependence of the Kerr angle $\theta$ on the pump density $W$, obtained by this method in sample 2, is shown in figure 10 [23]. The pump ($E_{\text{pump}} = 1.598 \text{ eV}$) and the probe ($E_{\text{probe}} = 1.597 \text{ eV}$) energies relate to the trion resonance (the PL spectrum of this sample is presented in figure 8(b)). The concentration of pumped electron spins $S_z \propto \theta^2$ increases monotonically with $W$ with a tendency to saturation for the pump densities exceeding $1 \text{ W cm}^{-2}$. The theoretical calculation of equation (5) for the trion resonant pumping (presented in figure 3(a)) is shown by the dotted curve in figure 10. It follows the experimental data quite well.

Figure 11(b) demonstrates efficiency of spin pumping in sample 2 ($E_{\text{pump}} = 1.597 \text{ eV}$) as a function of pump energy. Reflectivity (absorption) spectrum is presented in figure 11(a) for comparison. A pair of resonances associated with the neutral (heavy-hole) exciton $X$ and negatively charged trion $T$ is well resolved [5, 19]. Also several additional high-energy resonances are clearly seen in the reflectivity spectrum. Clear correlations in optical spectra of figure 11 are seen. At the low excitation density ($W = 0.25 \text{ W cm}^{-2}$) used the efficiency of spin pumping is nearly independent of whether the pumping is resonant with the heavy-hole exciton ($E_{\text{pump}} = 1.600 \text{ eV}$) or the trion ($E_{\text{pump}} = 1.598 \text{ eV}$), in accord with theoretical consideration in section 2.1 (equations (7) and (11)). The net spin polarization of resident electrons changes its sign when pumping is resonant with the light-hole exciton ($E_{\text{pump}} = 1.614 \text{ eV}$) as expected from the optical selection rules [1]. Remarkably, the spin pumping is still efficient without the tendency to decrease for pump energies up to $50 \text{ meV}$ above the bottom of the conduction band, when electrons with large kinetic energy are created [24].

Figure 12 shows the Hanle-MOKE curves for different probe energies while the pump energy is fixed at the exciton resonance, $E_{\text{pump}} = 1.600 \text{ eV}$. These data are recorded at low excitation (pump density $W = 0.25 \text{ W cm}^{-2}$). When the
We fit these data (solid line) by a sum of two Lorentz curves of equation (2) with different $B_{1/2}$ (dashed lines). Details of fitting procedure are discussed elsewhere [23].

We find spin relaxation time of $\tau_s = 30$ ns in addition to $\tau_s' = 5$ ns. As the spin pump conditions are identical in both cases (i.e., figures 12(a) and (b)), the difference can be ascribed only to the selective sensitivity depending on the probe energy. A possible explanation originates from a spatially inhomogeneous distribution of resident electrons in the QW plane [26]. The trions being charged are more sensitive to the localization (for instance, in the electrostatic potential of ionized barrier donors) as compared to the neutral excitons. As a result, when the detection energy is set in resonance with the trion state (solid symbols in figure 13). With increasing temperature localized electrons become delocalized, and as a consequence their spin relaxation rate $1/\tau_s$ rises. In section 3.4.3 we compare these results with data obtained in TRKR experiments and discuss mechanisms of spin relaxation.

### 3.4. Experimental results: pulsed excitation

We turn now to time-resolved experiments performed on sample 2 by means of pump–probe Kerr rotation.

#### 3.4.1. Time-resolved Kerr rotation (TRKR): one-color mode

Spin beats for different magnetic fields (pump and probe are set in resonance with the trion) in sample 2 are shown in figure 14. From fits to equation (14) one can find the spin dephasing time $T^*_\varphi$ the precession frequency $\Omega$ and the initial phase $\varphi$. In the case when anisotropy effects of equation (24) can be neglected (high magnetic fields) $\Omega$ is equal to the Larmor frequency $\omega_L = g_\nu \mu_B B / h$ allowing us to determine precisely the electron $g$-factor, $|g_\nu| = 1.64$. The spin dephasing time $T^*_\varphi$ decreases with magnetic field as $1/B$ (see inset) in accord with equation (23). Hence, in order to avoid the contribution of inhomogeneous distribution of $g$-factor values to the spin
The initial phase $\phi$ depends on the external magnetic field. The point is that the spins of the electrons forming a trion are in a singlet state and therefore do not precess around the magnetic field as the resident electrons do. Recombination of trions returns partially $z$-polarized electrons to the 2DEG. Their spin orientation differs from that of the precessing electrons, which results in a shift of $\phi$ [4]. We have analyzed the initial phase by extrapolating the spin beats measured at longer delays back toward zero delay. The results are given in figure 15. A pronounced minimum is seen at $B \approx 0.7$ T in qualitative agreement with the model calculations shown by the solid line (equation (20)), performed with a trion lifetime $\tau_T^0 = 30$ ps (obtained from the PL decay time [4]) and a spin relaxation time of the hole in a trion $\tau_T^s = 60$ ps (used as the only fitting parameter). We also assumed an electron spin dephasing time $T^*_2 = 2$ ns (which corresponds to the experimental value at $B = 3$ T). The model demonstrates qualitative agreement with the experimental data.

The time-resolved Kerr rotation signal measured for resonant pumping of either trion or exciton resonances is shown in figure 16, probe energy coincides with the pump one. Only long-lived spin beats of resident electrons are observed when the laser is resonant with the trion and two component decay is seen for the laser resonant with the exciton. The experimental conditions (namely, low pump density $1.5$ W cm$^{-2}$ and weak magnetic field $B = 0.25$ T) are chosen here to achieve longer electron spin coherence times [4, 17]. The delay time range between pump and probe in figure 16 covers 6 ns. Under these conditions the spin dephasing time of the resident electrons reaches 13.7 ns for pumping in the trion resonance and 4.2 ns for pumping in the exciton resonance. Moreover the spin coherence does not fully decay during the time interval of 13.2 ns between the pump pulses as is clearly seen by the beats at negative delays.

One can also see that the signal measured on the trion resonance contains only the long-lived spin beats of resident electrons. However the exciton signal clearly shows two component decays. The fast decay with a time of 50 ps is due to electron precession in photogenerated excitons. It disappears with the radiative decay of excitons. The long-lived component with a time of 4.2 ns is due to the resident electrons, in which spin coherence is generated by excitons captured to trions shortly after photogeneration.

Theoretical analysis predicts qualitatively different dependences of the Kerr rotation amplitude on excitation density for the excitons and trions (see equations (9), (13)). It is confirmed by experimental data from figure 17. Here the values of the Kerr signals normalized to their maximum values...
according to equation (13). In the latter case, the fitting parameter is the electron spin density equal to \( n_e/4 \) which, according to our theoretical predictions, correspond to the electron spin density equal to \( n_e/4 \) are shown. This allows us to compare the efficiency of spin coherence generation. One can see that in the low pumping regime the spin coherence generation efficiency per absorbed photon is practically the same for the laser tuned either to the exciton or the trion resonance. This is in good agreement with cw experiments of figure 10 (see also discussion in section 3.3.2) and the theory (sections 2.1.1 and 2.1.2): each absorbed photon creates a trion either directly or via an intermediate excitonic state, thus an electron with a given spin orientation is removed from the 2DEG. The strong pumping regime is different for exciton and trion excitation. In the case of the laser tuned to the trion resonance the spin of the 2DEG saturates (the small decrease of the Kerr rotation amplitude can be attributed to heating of the 2DEG) while for the laser tuned to the exciton resonance a strong decrease of the spin coherence generation efficiency is seen.

The curves in figure 17 are the results of theoretical calculations based on the models outlined in sections 2.1.1 and 2.1.2. The dashed curve corresponds to trion resonant excitation, while the solid line is for exciton resonant excitation. For the dashed curve the only fitting parameter was the saturation level. For the solid line the only fitting parameter was the ratio between the electron-in-exciton spin relaxation time and the exciton radiative lifetime, which is \( T_\chi^X/T_0^X = 10 \). The spin relaxation of an electron in the exciton \( T_\chi^X \sim 0.5 \text{ ns} \) appears to be shorter than that for the resident electrons. It has been shown that the electron–hole exchange interaction within the exciton provides an efficient spin decay channel in the external magnetic field [21].

### 3.4.2. Time-resolved Kerr rotation (TRKR): two-color mode

In QWs with low electron concentration (such as sample 2) the Fermi energy smaller than the typical localization potential is caused by well width fluctuations. As a result, at low temperatures a major fraction of electrons is localized. In order to have a deeper insight into the effects of electron localization on the spin coherence two-color pump–probe measurements have been performed in the regime where the longest dephasing times have been achieved, i.e. using a low pump density in a weak magnetic field of 0.25 T [4]. Figure 18 shows the results of such an experiment, where the pump beam is resonant with the exciton transition and the Kerr rotation signal is detected at either the trion or the exciton energy. As the excitation conditions are identical for both signals in figure 18, one would expect to detect the same dephasing times for the 2DEG, irrespective of the probe energy. Therefore, it is rather surprising to observe that the dephasing times of the electron coherence differ by a factor of 2: \( T_\chi^X = 10.8 \text{ ns} \) and \( 5.3 \text{ ns} \) for probing at the trion and the exciton resonance, respectively. To explain this difference we suggest that different fractions of the resident electrons contribute to the Kerr rotation signal measured at the trion or the exciton energies. This can be understood if we take into account that different mechanisms lead to the Kerr rotation signal at the exciton and trion energy. For detection at the trion resonance the effect is contributed by the variation of the trion oscillator strength, which is directly proportional to the concentration of electrons with a specific spin orientation [4, 12]. The trion stability (i.e., its oscillator strength) increases when resident electrons are localized in the QW plane [26], and these electrons possess a longer spin dephasing time. The Kerr rotation signal detected at the exciton energy monitors 2DEG spin beats mostly due to the spin-dependent exciton–electron scattering [27]. The possibility for this scattering to occur implies the existence of free electrons. Therefore, at the exciton energy we address free or quasi-free resident electrons which have shorter spin dephasing times. This result is in good agreement with experimental data and conclusions drawn from the cw experiments of figure 12.
3.4.3. Resonant spin amplification. In certain magnetic fields the spins excited by the pulse train come in phase resulting in their amplification [10]. Such resonant spin amplification (RSA) in sample 2, presented in figure 19, was obtained for a small negative time delay of −80 ps between the probe and pump pulses. The spin dephasing time $T_2^*$ time can be directly evaluated from the fit of RSA peak to equation (26). The spin dephasing time of $T_2^* = 30$ ns is evaluated.

**Figure 19.** (a) Resonant spin amplification measured at a negative delay of −80 ps in sample 2. Pump density is 0.05 W cm$^{-2}$. (b) The fit (dashed line) of a RSA peak to equation (26). Spin dephasing time of $T_2^* = 30$ ns is evaluated. $T = 2$ K.

Spin beats and spin decay of resident electrons in external magnetic fields are shown to be a subject of spin relaxation [28, 29]. The linear dependence of spin relaxation rate 1 $\propto \beta \tau_p$, where the signal is still detectable, is the momentum relaxation time and $\beta$ is a coefficient characterizing dependence of the electron spin splitting on its wave vector. The effects of structure anisotropy on spin relaxation are neglected. The linear dependence measured experimentally allows us to suggest that in the studied sample $\tau_p$ is independent of temperature for $T < 100$ K. This is in agreement with the fact that the electron mobility in CdTe-based QWs is rather low (usually does not exceed a few tens of thousands V cm$^{-2}$ s$^{-1}$) and $\tau_p$ is rather short and falls in the picosecond range.

4. Conclusions

With the use of various techniques based on continuous wave and pulsed excitation we perform detailed experimental studies on optical spin pumping, spin relaxation and spin coherence of resident electrons in CdTe/Cd, Mg/Te quantum wells. The experimental results are substantiated by a theoretical model based on the classical approach to spins. Polarization of the resident electrons and generation of their spin coherence is provided by the capture of the resident electrons into trions. The trions in turn can be either photogenerated by resonant excitation or formed from the excitons. Variation of pump energy (tuning to resonance with either the trion or the exciton) highlights details of the polarization process. Pump density dependences for the two excitation conditions are nearly the same for low densities (as each absorbed photon participates in spin pumping) and substantially different for high densities (for the trion pump the spin polarization tends to 100% while for the exciton pump the spin polarization decreases to zero).

Spin beats and spin decay of resident electrons in external magnetic fields are shown to be a subject of spin relaxation anisotropy and $g$-factor inhomogeneity. Independent variation of pump and probe energies proves that electron localization provides an increase of the spin relaxation time, as also confirmed by temperature-dependent experiments. The spin relaxation time of localized electrons can achieve 30 ns, as independently found in the continuous wave experiment in the limit of low pump density and in pulsed experiments by means of resonant spin amplification. Note that this is the longest spin relaxation time in QW heterostructures reported to date.

Acknowledgments

We acknowledge fruitful discussions with E L Ivchenko. Samples for these study have been grown in the Institute of Physics, Warsaw by G Karczewski, T Wojtowicz and J Kossut. This work was supported by the Deutsche Forschungsgemeinschaft (SPP1285 Spintronik via grants nos 958/0-1) as well as by the Russian Foundation for Basic Research. MMG was partially supported by the ‘Dynasty’ foundation—ICFPM.

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