Optical initialization and dynamics of spin in a remotely doped quantum well

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The excitation of electron spin polarization and coherence by picosecond light pulses and their dynamics in a wide remotely doped quantum well are studied theoretically and experimentally. Assuming that all electrons in the quantum well are localized, the theory considers the resonant interaction of light pulses with the four-level system formed by the electron spins of the ground state and the hole spins of the trion excited state. The theory describes the effects of spontaneous emission, a transverse magnetic field and hole spin relaxation on the dynamics detected by the Kerr rotation of a probe pulse. Time resolved Kerr rotation experiments were carried out on a remotely doped 14 nm GaAs quantum well in the frequency range of optical transitions to the heavy hole (HH) trion and to the light-hole (LH) trion degenerate with the HH exciton. The experiments on the resonant excitation of the HH trion show a very slow heavy hole spin relaxation and, consequently, a weak electron spin polarization after the trion relaxation. In contrast, the resonant excitation of the LH trion/HH exciton results in a fast hole spin relaxation that increases electron spin polarization.

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I. INTRODUCTION

Currently there is great interest in developing electron spin in semiconductors as a useful property for quantum information technologies and other applications. Some of this effort is focused on exploiting the intimate connection between electron spin and light polarization to optically initiate and control spin states. Previous studies have revealed that optically generated spins in semiconductor quantum wells are influenced by recombination, electron-hole exchange, hyperfine interactions, and localization.[1–5] Most of the published work addressed excitons in undoped quantum wells. These systems exhibit a relatively simple energy-level structure and a good understanding of the physics has been achieved.

Doped semiconductors provide the additional advantage that the electronic spin can exist after recombination of an associated optically excited state.[6] In our previous work on donors in GaAs layers, the low concentration of donors provides an analog of remotely doped quantum dots.[7] However, studies of these doped layers with time-resolved Kerr rotation (TRKR) gave results that were complicated with phase changes and multicomponent decays caused by the mixing of light and heavy holes. Quantum wells offer a simpler case because the light and heavy holes are split in energy. Thus we have made this study of electrons and their associated optical transitions (trions) in remotely doped wide quantum wells. Little previous work has been done in this area.[8,9]

In pump-probe experiments like TRKR, the first (pump) pulse changes the ensemble of quantum systems coherently acting on each single system. While the second (probe) pulse propagates through the ensemble of coherently prepared systems, its polarization is affected by the scattering of each system. The result can be detected by various techniques. Among them are the time resolved Kerr rotation and Faraday rotation of light polarization.[3,5,10,11] Both pump and probe pulses are shorter than the relaxation times and the relaxation processes do not affect the interaction with the pump and probe pulses. The period of time between the pulses can be long enough to observe the effects of the relaxation.

In this paper we present theory and experiments for resonant pump-probe (TRKR) studies of electrons and trions in a remotely doped wide quantum well. The results concentrate on the electron—heavy-hole (HH) singlet trion system. This trion is the lowest-energy optical excitation in the system. This case is especially interesting because it can be easily modeled and the relatively small inhomogeneous broadening maintains a long lifetime for the precessing spin, which allows one to get a significant insight into the physical processes of spin excitation and dynamics. We find that the excitation of the spin of the resident electrons depends on the recombination time, the hole relaxation time, the magnetic field and the energy chosen for the light pulses. The results provide a physical picture of how light pulses polarize the spin of resident electrons in weakly doped quantum wells.

II. MICROSCOPIC PHYSICS FOR THE HEAVY-HOLE TRION

We consider an optical interaction with localized electrons in a remotely doped quantum well. The pump and probe light is tuned to the resonance with the trion state formed by two electrons in a singlet state and a heavy hole. The $\pm 1/2$-spin electron states and the $\pm 3/2$-spin trion states comprise the four-level system depicted in Fig. 1 where the labels show the spin projections in the direction of light propagation. According to the selection rules, σ^+ polarized light excites transitions only between the $+1/2$ electron and $+3/2$ trion states while σ^- polarized light interacts only with the $-1/2$ and $-3/2$ pair of states. Emitting a spontaneous photon, a trion state decays on average within the emission time τ_r. Obeying the selection rules, the $+3/2$ and $-3/2$ trion states decay to the $+1/2$ and $-1/2$ electron states, respectively.

The probe pulse is tuned to the same resonance as the pump pulse between the $\pm 1/2$ and $\pm 3/2$ states. The linearly

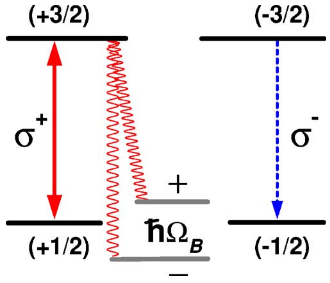

FIG. 1. (Color online) Energy level diagram for the electron and heavy-hole singlet trion. The labels show the electron and trion spin projections in the direction of light propagation. The $\pm 3/2$ and $\pm 1/2$ states are coupled by plus/minus circularly polarized light propagating in the growth direction. A magnetic field applied perpendicular to the growth axis produces superpositions (labeled + and −) of the $\pm 1/2$ states split in energy by ħ times the Larmor frequency Ω_B.

polarized probe is a coherent superposition of two circularly polarized components σ^+ and σ^-, which are scattered by the $(+1/2,+3/2)$ and $(-1/2,-3/2)$ pairs of electron-trion states, respectively. The ensemble of four-level systems, coherently prepared by the pump pulse and driven by the probe pulse, produces a wave front moving in the direction of the probe light. After propagation through the ensemble, the σ^+ and σ^- components of the linearly polarized probe light acquire different phase shifts: φ_+ and φ_- induced by the pump pulse. The probe light initially polarized in the x direction will have both components: $E_x=E_0 \cos\theta_K$ and $E_y=E_0 \sin\theta_K$, where $\theta_K=(\varphi_+-\varphi_-)/2$ is the polarization angle. The phase shift is proportional to the difference in the average population of the states that scatter the probe light: $\varphi_\pm \propto f_{\pm 3/2}-f_{\pm 1/2}$. The populations evolve with time changing the phase shifts and the polarization angle: $\theta_K(t)=(\varphi_+-\varphi_-)/2 \sim f_{-1/2}-f_{+1/2}+f_{+3/2}-f_{-3/2}$.

A transverse magnetic field mixes the $+1/2$ and $-1/2$ states. In general, an arbitrary spin state can be written in the form: $a_{+1/2}|+1/2\rangle + a_{-1/2}|-1/2\rangle$, where coefficients $a_{+1/2}$ and $a_{-1/2}$ determine the spin direction. Half of the population difference between the $\pm 1/2$ states, $S_L=(|a_{+1/2}|^2-|a_{-1/2}|^2)/2$, is the spin polarization in the direction of light propagation. The perpendicular spin polarization, in the plane of the quantum well, is the product of the coefficients: $S_B=a_{+1/2}\cdot a^*_{-1/2}$. In a transverse magnetic field, the coefficients $a^+_{+1/2}=a^+_{-1/2}=1/\sqrt{2}$ and $a^-_{+1/2}=-a^-_{-1/2}=-1/\sqrt{2}$ correspond to the $|+\rangle$ and $|-\rangle$ states, split by the Zeeman energy $\hbar\Omega_B$(see Fig. 1), with the spin parallel and anti-parallel to the magnetic field: $S^\pm_B=a^\pm_{+1/2}\cdot a^\pm_{-1/2}=\pm 1/2$. For both states, the spin in the direction of light propagation is zero, $S^\pm_L=0$.

At a temperature higher than the Zeeman splitting, an ensemble of N spins is unpolarized and individual spins are equally distributed between the $|+\rangle$ and $|-\rangle$ states. The $+1/2$ and $-1/2$ states are also equally populated with the average populations: $f_{\pm 1/2}=\langle |a_{\pm 1/2}|^2\rangle = 1/2$. In all directions, the average spin polarization of the ensemble is zero, $\langle S_L\rangle = \langle S_B\rangle = 0$.

The initially unpolarized ensemble of spins becomes polarized by the σ^+ polarized light. (The σ^+ is chosen for definiteness; one could also start with σ^-.) Acting on a single spin initially in the state $a_{+1/2}|+1/2\rangle + a_{-1/2}\rangle|-1/2\rangle$, the electric field E of the σ^+ polarized light, interacting with the transition dipole d, drives coherently the $+1/2$ component to the $+3/2$ state. After the pulse, the $+1/2$ and $-1/2$ states form a new superposition with the $+3/2$ state: $\bar{a}_{+1/2}|+1/2\rangle + \bar{a}_{-1/2}|-1/2\rangle + \bar{a}_{+3/2}+3/2\rangle$. For a pulse length Δt_p much shorter than the period of spin precession in the magnetic field, $\Delta t_p \ll \Omega_B^{-1}$, the spin rotation during the pulse can be neglected. The population of the $+1/2$ state is reduced by the amount moved to the $+3/2$ state $|\bar{a}_{+1/2}|^2=|a_{+1/2}|^2-|\bar{a}_{+3/2}|^2$. The $-1/2$ state does not interact with the σ^+ light and thus the $\bar{a}_{-1/2}$ coefficient remains equal to $a_{-1/2}$. In particular, if the spin is initially in the $|+\rangle$ or $|-\rangle$ state, the population of the $+3/2$ state, after the pulse, becomes $\bar{a}^2_{+3/2}=(1-\cos\Omega_R\Delta t_p)/4$, where $\Omega_R=2Ed/\hbar$ is the Rabi frequency of the pumping light proportional to the square root of its intensity, $\sqrt{I}$. The excess of the $-1/2$ state population over the $+1/2$ state $|a^\pm_{-1/2}|^2-|\bar{a}_{+1/2}|^2=|\bar{a}_{+3/2}|^2$ can be viewed as the superposition of the $|-\rangle$ and $|+\rangle$ states which are mixed by the pump pulse. The pump pulse changes the original spin polarization of the $|+\rangle$ and $|-\rangle$ states from $S_L=0$ and $S_B=\pm 1/2$ to $S_L=(\cos\Omega_R\Delta t_p-1)/8$ and $S_B=\pm(1/2)\cos(\Omega_R\Delta t_p/2)$, respectively. At low pumping intensity $\Omega_R\Delta t_p \ll 1$, the small deviation of spin from the initial direction is equal to $(\Omega_R\Delta t_p)^2/8 \propto I$. At high intensity, the spin polarization changes nonlinearly. For a π-pulse, $\Omega_R\Delta t_p=\pi$, the S_L spin polarization saturates at $-1/4$ and the S_B polarization turns to zero. A single spin, initially collinear with the magnetic field, becomes completely polarized in the direction perpendicular to the magnetic field and begins precession around it.[12]

The ensemble is changed uniformly by a pump pulse shorter than the radiative decay time, $\Delta t_p \ll \tau_r$. After summing the contribution from each single spin, the average populations and the spin polarizations are

$$
\begin{aligned}
f_{+3/2} &= \langle |\bar{a}_{+3/2}|^2\rangle = \tfrac{1}{4}(1-\cos\Omega_R\Delta t_p),\\
f_{+1/2} &= \langle |a_{+1/2}|^2\rangle - \langle |\bar{a}_{+3/2}|^2\rangle = \tfrac{1}{2} - f_{+3/2},\\
f_{-1/2} &= \langle |a_{-1/2}|^2\rangle = \tfrac{1}{2},\\
\langle S_L\rangle &= \tfrac{1}{8}(\cos\Omega_R\Delta t_p - 1),\\
\langle S_B\rangle &= 0.
\end{aligned}
\tag{1}
$$

Let us consider now the decay of the coherent superposition created by the short optical pulse. First, we assume that the electron and hole spin relaxation is much slower than the spontaneous emission. Interacting with spontaneous photons, the $+3/2$ trion component cannot remain in the coherent superposition with the ground state, $\bar{a}_{+1/2}|+1/2\rangle + a_{-1/2}|-1/2\rangle + \bar{a}_{+3/2}|+3/2\rangle$ created by the pump pulse. Instead, the system emits a spontaneous photon with the total probability $W_B=|\bar{a}_{+3/2}|^2=|a_{+1/2}|^2-|\bar{a}_{+1/2}|^2$ or it stays in the ground state with the probability $W_D=|a_{-1/2}|^2+|\bar{a}_{+1/2}|^2$.

At a low magnetic field, $\Omega_B\tau_r \ll 1$, the precession does not change the ground state wave function on the scale of radiative relaxation time. The ground state wave function can be

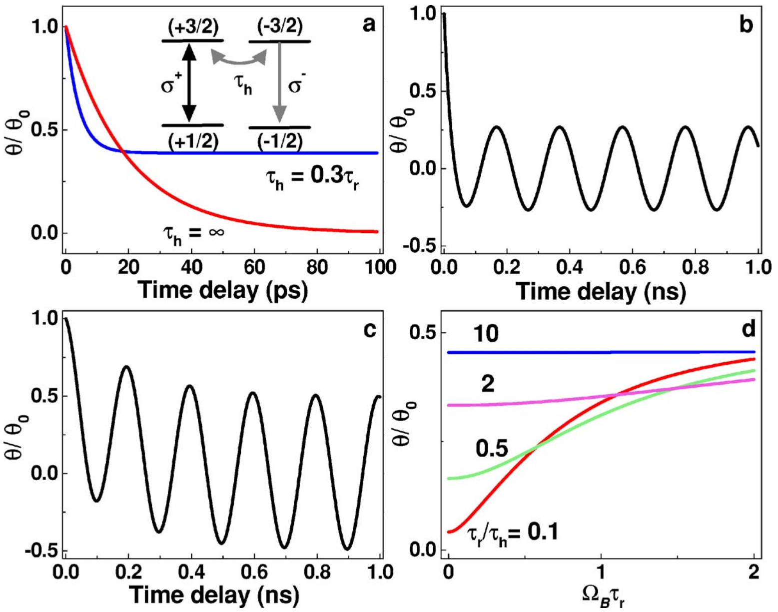

FIG. 2. (Color online) Theoretical dependence of the Kerr rotation angle θ_K on time and magnetic field. (a) The time dependence of the rotation angle at zero magnetic field without hole spin relaxation ($\tau_h=\infty$) and with hole spin relaxation where $\tau_h=0.3\tau_r$. The inset shows the role of hole relaxation on the populations of the states. Parts (b) and (c) show the time dependence of the rotation angle without hole relaxation. In (b) the Larmor precession is slow with respect to the recombination time ($\Omega_B\tau_r=0.3$); in (c) the Larmor precession is fast with respect to the recombination time ($\Omega_B\tau_r=3.0$). (d) The magnetic field dependence of rotation angle after trion relaxation for various hole spin relaxation times.

written as $b_{+1/2}|+1/2\rangle+b_{-1/2}|-1/2\rangle$ where $b_{\pm1/2}=W_D^{-1/2}\bar{a}_{\pm1/2}$. The average population of the $-1/2$ component of the ground state remains unchanged by the pump pulse: $f_{-1/2}=\langle W_D\cdot|b_{-1/2}|^2\rangle=\langle|a_{-1/2}|^2\rangle=0.5$. The pulse reduces the population of the $+1/2$ component by W_B. For the slowly precessing ground state, the radiative decay of the trion populates the $+1/2$ state only.[13] After radiative relaxation, the population of the $+1/2$ state is split between two alternatives with the probabilities W_B and $W_D|b_{+1/2}|^2=|\bar{a}_{+1/2}|^2$, which add up on average to the total $f_{+1/2}=\langle|a_{+1/2}|^2\rangle=0.5$, the same as the population before the pump pulse. The trion relaxation compensates the lack of the $+1/2$ component in the ground state produced by the pump pulse. The average spin polarization returns to zero: $\langle S_L\rangle=\langle S_B\rangle=0$. Figure 2(a) shows how θ_K decreases to zero with delay time from the initial angle of rotation, θ_0. The short-range behavior of the rotation angle is determined by the relatively fast relaxation of the $+3/2$ state: τ_r.

At high magnetic field $\Omega_B\tau_r\gg1$, the period of the ground state precession is much shorter than the radiative decay time of the excited state, which can emit a spontaneous photon with the same probability W_B as at a low magnetic field. The $+3/2$ state, however, cannot coherently follow the fast precession of the $+1/2$ state and the excited state decays into one of the stationary states: $|+\rangle$ or $|-\rangle$, populating the $+1/2$ and $-1/2$ states with equal probability $0.5W_B$. With the probability W_D, the system does not emit a spontaneous photon and remains in the coherently rotating ground state whose wave function is $b_{+1/2}(t)|+1/2\rangle+b_{-1/2}(t)|-1/2\rangle$, where $b_{\pm1/2}(t)=W_D^{-1/2}[\bar{a}_{\pm1/2}\cos(\Omega_Bt/2)+i\bar{a}_{\mp1/2}\sin(\Omega_Bt/2)]$. After radiative decay, the average population of the $+3/2$ state is zero and the populations of the $+1/2$ and $-1/2$ states oscillate as $f_{\pm1/2}=\langle0.5W_B+W_D|b_{\pm1/2}|^2\rangle=0.5(1\pm\delta f\cos\Omega_Bt)$, where $\delta f=\langle|\bar{a}_{+1/2}|^2-|a_{-1/2}|^2\rangle=\langle|\bar{a}_{+1/2}|^2\rangle-0.5$ is the population difference induced by the pump pulse. The average spin polarization is perpendicular to the magnetic field and also depends on time: $\langle S_L\rangle=0.25\cdot\delta f\cdot\cos\Omega_Bt$ and $\langle S_B\rangle=0$. At high magnetic field $\Omega_B\tau_r\gg1$, the rotating spin polarization is not compensated by the radiative decay.[12] Figures 2(b) and 2(c) show the time evolution of the polarization angle for slow and fast precession. In addition to the short-range decay, the rotation angle shows oscillations, which decay on a much longer scale of time. The amplitude of the oscillation grows with magnetic field, which indicates the increase in the spin polarization coherently precessing around the magnetic field.

So far we have assumed that the hole relaxation is slow. Let us now consider the effect of hole relaxation on the population and electron spin polarization. The pump process is not affected by the relaxation because the hole relaxation time, τ_h, is longer than the pulse length. Before radiative decay, however, the hole relaxation has a chance to flip the $+3/2$ trion state to the $-3/2$ state with the probability $P=0.5(1+\tau_h/\tau_r)^{-1}$. After the trion radiative relaxation, the average population of the $+1/2$ state is not restored to its original value of one-half even at a low magnetic field. Instead of one-half, the average populations are $f_{\pm1/2}=0.5\mp\langle|\bar{a}_{+3/2}|^2\rangle P$. The average uncompensated spin is perpendicular to the magnetic field and equal to $\langle S_L\rangle=-\langle|\bar{a}_{+3/2}|^2\rangle P$. Fast hole relaxation leads to strong electron spin polarization. At $\tau_h\ll\tau_r$, the excited state returns to the $+1/2$ and $-1/2$ spin states with equal probabilities: $0.5\langle|\bar{a}_{+3/2}|^2\rangle$.The average spin, $\langle S_L\rangle=-0.5\langle|\bar{a}_{+3/2}|^2\rangle$, is one-half of the excited state population and it does not vary with magnetic field. Without magnetic field, the time evolution of the probe polarization angle is shown in Fig. 2(a). After the hole and trion relaxation, the spin polarization approaches one half of its initial value. The rotation angle θ_K also decreases only by one-half instead of falling to zero as it does when the hole relaxation is slow. As the hole relaxation changes from slow to fast the amplitude of Kerr rotation becomes less dependent on magnetic field [see Fig. 2(d)].

III. EXPERIMENTS AND DISCUSSION

The sample used consists of five quantum wells of different thicknesses grown with interrupts to produce large mono-

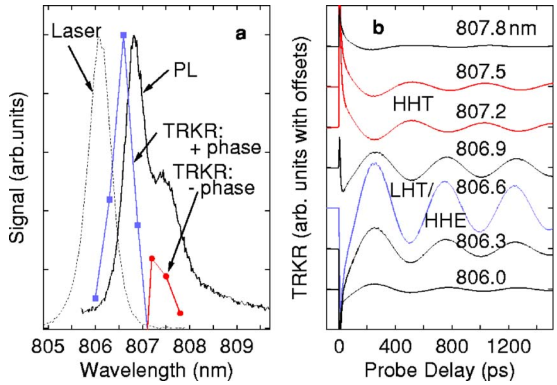

FIG. 3. (Color online) Photoluminescence and spectral dependence of the Kerr rotation. Figure 3(a) shows two partially resolved lines in photoluminescence attributed to the HH trion (807.5 nm) and the LH trion/HH exciton (806.8 nm), the spectral dependence of the amplitude of the Kerr rotation for the HH trion (circles) and the LH trion/HH exciton (squares), and the spectrum of the 1.5 ps light pulse (dotted curve). Figure 3(b) shows the Kerr rotation for different wavelengths with the magnetic field at 500 mT and the temperature at 6 K.

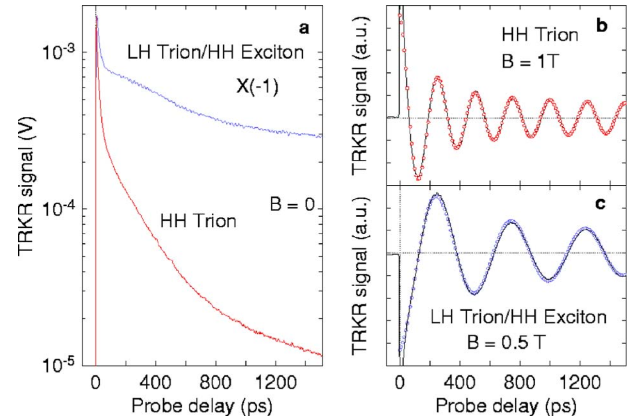

FIG. 4. (Color online) Form of the decays for $B=0$ and small magnetic fields with light resonant with the HH trion and with the LH trion/HH exciton. In Fig. 4(a), the data for $B=0$ show three decays: one fast for $t<100$ ps, one intermediate of about 200 ps, and one in the ns range. In Figs. 4(b) and 4(c), the intermediate and long decays at the fields shown have been fit with an exponential times cosine. The results of the fits are given in Table I.

layer islands. It was Si-modulation doped to produce an electron density in the wells of 3×10^{10} cm^{-2}. For this study we have focused on the 14 nm well. More details on the structure and properties are given in Ref. 14. The time-resolved Kerr rotation (TRKR) experiments were performed in the Voigt geometry with 1.5 ps pulses produced by a mode-locked Ti:sapphire laser with an 82 MHz repetition rate. The pump and probe pulses were degenerate in energy. The pump pulse was modulated from σ^+ to σ^- at 20 kHz to reduce the amount of laser noise from the pump beam and the probe pulse was linearly polarized. The Kerr rotation was measured by splitting the reflected probe into orthogonal polarizations and detecting them with balanced photodiodes. The difference signal from the photodiodes was then lock-in detected in phase with the pump modulation. The sample was placed in the center of a superconducting magnet and cryostat and most of the experiments were done at 6 K. The TRKR setup is similar to that of Crooker *et al.*[15]

The photoluminescence for this quantum well is shown in Fig. 3(a). There are two peaks and excitation-power studies confirm that the shorter wavelength peak comes from the heavy-hole neutral exciton and the longer from the heavy-hole trion. For a well of this thickness, the neutral exciton is expected to be nearly degenerate with the light-hole trion.[4] The dotted curve shows the spectrum of the laser, which is equal in width to the emission lines.

The spectral dependence of the TRKR at 500 mT is shown in Fig. 3(b). As expected, enhancements in the signal occur that can be correlated with the peaks in the photoluminescence spectrum. We identify the strong signal at 806.6 nm that starts off like a minus-cosine function with the LH trion/HH exciton and the secondary signal around 807.35 nm that starts off like a cosine function with the HH trion. The amplitudes of these oscillatory signals are plotted in Fig. 3(a) and show a slight displacement to shorter wavelength (higher energy) from the photoluminescence features. If the spectral dependence of the Kerr effect for these transitions is symmetric,[15] this difference between absorption and emission can be taken as evidence of localization for the exciton and trions.[2]

The Kerr-rotation results exhibit multiple decays. More specifically, there are three. Data taken at zero magnetic field for both resonant wavelengths [Fig. 4(a)] show a fast initial decay of about 30 ps, a middle decay of around 200 ps, and a slow decay of a few nanoseconds. This pattern persists when the magnetic field is applied. Here we expect the decay for each component i to be of the form

$$\theta_{Ki}=A_i\exp(-t/\tau_i)\cos\Omega_i t, \tag{2}$$

where θ_{Ki} is the Kerr rotation, τ_i is the lifetime, $\Omega_i=\mu_B gB$ the Larmor frequency, μ_B is the Bohr magneton, g_i is the effective g-factor, and B is the in-plane magnetic field. Focusing first on the HH trion, Fig. 4(b) shows that the multiple decays are present for $B=1$ T as well, with the slower two now showing oscillation. The oscillatory parts have been fit with the first having $g_1=0.268$ and τ_1 of 170 ps and the second having $g_2=0.286$ and $\tau_2=2.5$ ns. The form of the decay is similar, but not identical, to the form predicted by the model presented in Sec. II. At the position of the LH trion and HH exciton, the signal starts off as a negative cosine. The data for $B=0$ [Fig. 4(a)] and $B=0.5$ T [Fig. 4(c)] show three distinct decays similar to those observed for the HH trion. The initial decay is again less than 50 ps. The longer decays were fit to give $g_1=0.269$ and τ_1 of 250 ps and $g_2=0.289$ and $\tau_2=2.5$ ns. These values are very similar to those obtained for the HH trion (see Table I).

The theory presented for the HH trion can account for most of the results obtained with light resonant with the HH trion in the following way. From the model for σ^+ excitation, the initial signal is comprised of $+3/2$ hole orientation and $-1/2$ electron orientation ($B=0$) or coherent superposition ($B>0$) resulting from depletion of the $+1/2$ state with respect to the $-1/2$ state. Direct recombination removes the

TABLE I. Results of the fitting of the intermediate and long decays for cases with light resonant with the HH trion and LH trion/HH exciton.

	HH exciton		Ground-state electron	
	g^*	t (ps)	g^*	t (ns)
HH trion position (807.35 nm)	0.268	170	0.286	2.5
LH trion/HH exciton position (806.6 nm)	0.269	250	0.289	2.5

oriented hole and cancels the ground-state difference in population. Without any other effect, the signals would be gone quickly following recombination. This is in fact the general form of the signal shown by the fast and slow parts of the experimental data both at $B=0$ and 1 T [see Figs. 4(a) and 4(b)]. From the literature, the radiative lifetime for the trion is 60 ps.[16] This value corresponds roughly to the fast decay seen in the experiments of around 30 ps. We associate the fast decay seen in Kerr rotation with the loss of ground-state difference in population caused by the trion relaxation. The long (2.5 ns) component of the data is readily attributed to the ground-state spin. Both the long lifetime and the g-factor[3] confirm this association. The existence of some ground state spin at long times even in zero magnetic field indicates that there is some hole relaxation competing with the recombination. From the small ratio of final to initial signal seen for $B=0$, we can safely assume that $\tau_h > 10\tau_r$. Thus τ_h is at least 300 ps giving further evidence of localization of the trion under these conditions.[2,17]

The theory for the trion does not account for the middle decay with characteristic time of 170 ps. Since the spectral separation of the HH trion and LH trion/HH exciton in the photoluminescence and Kerr rotation is not complete, we attribute the middle decay to a contribution from the HH exciton—more specifically from the excitonic electron. The decay is longer than that the purely radiative decay of the trion since the exciton has a dark (spin) state.[1]

Power studies performed at 0.5 T confirm that the response is third order in the electric fields (χ^3). The Kerr rotation is linear in pump power and independent of probe power.[18] The magnetic field dependence of the Kerr rotation, however, differs from the usual behavior. When the laser is resonant with the LH trion and HH exciton, the case is normal with the strength of the signal independent of magnetic field [Fig. 5(a)].[19] The decay for $B=0$ T forms an envelope that the decays for finite magnetic fields follow closely. The signal strength depends only on the pump power. In contrast, the amplitude of the signal varies with magnetic field when the light is resonant with the HH trion [Fig. 5(b)]. Beyond 0.2 T, the slowly decaying portion of the signal increases [inset to Fig. 5(b)]. The magnetic field dependence for the HH trion provides a nice confirmation of the model for the spin excitation. Cancellation of the ground-state difference in population requires that the phase of the ground state remain fixed since recombination of the σ^+ trion returns an electron

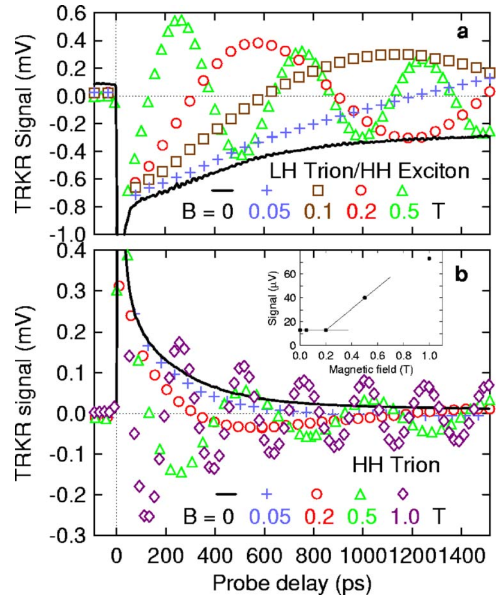

FIG. 5. (Color online) Effect of the magnetic field on the amplitudes of the Kerr rotation signals. In Fig. 5(a), with light resonant with the LH trion and HH exciton, the decay for $B=0$ forms an envelope for the decays at finite magnetic fields. In Fig. 5(b), with light resonant with HH trion, the amplitude of the slowly decaying signal increases with increasing magnetic field. The inset in Fig. 5(b) shows the magnetic field dependence of the amplitude of the slowly decaying signal.

spin at the original phase. When Larmor precession of the ground-state electrons changes the phase significantly (changing the relative populations of 1/2 and −1/2), the cancellation no longer occurs. This leads to a sharp increase in the signal as the Larmor period becomes a significant fraction of the recombination time. If we take $\tau_r=60$ ps, the Larmor period is equal to it when $B=4.2$ T and 0.5 T is a significant fraction. The experimental result corresponds closely to the predicted field-dependence when the hole relaxation time is long as shown in Fig. 2(d).

With an understanding of the processes for light resonant with the HH trion, we proceed to discuss briefly the results for the higher-energy light resonant with the LH trion and HH exciton. Both the LH trion and HH exciton are excited by the pump pulse. For the (singlet) LH trion, the σ^+ polarized excitation excites the transition from the −1/2 electron-spin ground state to the +1/2 hole-spin excited state. This leaves an excess electron spin population in the +1/2 ground state and will give the opposite phase for the Kerr rotation. While the phase can change just from the spectral dependence of the Kerr effect,[15] it seems that the selection rules give the simplest explanation here. The two-dimensional density of states for the LH will overlap that of the HH and rapid relaxation to the lower-energy state will occur. This process will produce extremely rapid hole-spin relaxation that takes less than a picosecond. Thus the ground-state coherence will be preserved and a strong long-lived component is observed. For the HH excitons, if free, they could also find electrons and form trions. However, the behavior observed correlates better with localized excitons.[4] At this energy, the localized HH excitons produce the oscillating decay with a lifetime of 250 ps from their precessing electrons.

IV. CONCLUSION

The theory and experiments presented here show the mechanism by which circularly polarized light creates electron spin coherence of resident electrons in a wide GaAs quantum well. We have shown that the amount of coherent population created by resonant excitation of the singlet HH trion depends critically on the recombination time, the hole spin relaxation time and the magnetic field. Greater coherence occurs at resonant excitation of the singlet LH trion due to the rapid relaxation of the light hole spin. This work conducted on weakly localized electrons has laid the groundwork for the more precise control of the electron spin polarization proposed in Ref. 12 and should lead to its practical realization in quantum dots.

ACKNOWLEDGMENTS

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