Current-induced control of the electron–nuclear spin system in semiconductors on a micrometer scale

Yuansen Chen1, Jungtaek Kim1, Joachim Puls1, Fritz Henneberger1, and Gerd Bacher*1

1 Werkstoffe der Elektrotechnik and CENIDE, University of Duisburg-Essen, Bismarckstr. 81, 47057 Duisburg, Germany
2 Institut für Physik, Humboldt-Universität zu Berlin, Newtonstr. 15, 12489 Berlin, Germany

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Corresponding author: e-mail gerd.bacher@uni-due.de, Phone: +41 203 379 3406, Fax: +41 203 379 3404
Current address: IBM Zurich Research Lab, Rüschlikon, CH-8803, Switzerland.

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The ability of using onchip microcoils to control the electron–nuclear spin system in semiconductors is demonstrated. Electrically generated magnetic fields of several tens of mT can be obtained on a micrometer length scale, which are switchable on a sub-ns time scale due to the low complex coil impedance. This allows one to electrically (i) manipulate the nuclear spins by means of nuclear magnetic resonance in n-GaAs and (ii) control the hyperfine flip-flop rate in CdSe/ZnSe quantum dots.

1 Introduction To obtain local spin control is a fundamental issue toward applications of spin-based semiconductor devices [1]. This can be achieved either by an effective magnetic field [2], such as, e.g., a magnetic exchange field [3, 4], the spin–orbit field [5], and the Overhauser field [6, 7] or by a real external magnetic field, defined, e.g., by a microferromagnet prepared on top of the semiconductor [8, 9] or electrically generated by a microscale current loop [10, 11]. The latter approach permits a transient electrical spin control up to the GHz frequency regime and down to a micrometer length scale [12–14]. Using millimeter-sized external current loops, even the manipulation of single electron spins in individual quantum dots (QD) is achieved [15].

In semiconductors, the electron–nucleus spin coupling is dominated by the Fermi-contact hyperfine interaction [16, 17]. Since the electron wavefunction covers a certain volume in real space, one electron interacts simultaneously with many nuclear spins. In a QD or a shallow donor state, one electron typically experiences over 10^5 nuclear spins in III–V systems, while the nuclear spin number is of the order of 10^2 in the II–VI counterpart [18] due to the smaller lateral extension of the electron wavefunction and the lower abundance of spin-containing nuclei, see, e.g., Table 1.

In the III–V semiconductors GaAs or InAs, the hyperfine interaction constant A_{hf} is of the order of 100 \mu eV. The coupling can generate an Overhauser (B_{Oh}) field up to a magnitude of several Tesla, which has been observed in bulk semiconductors [17], quantum wells [19], and QDs [7, 18, 20]. The pronounced Overhauser field has been typically manifested by an energy shift of the photoluminescence (PL) spectrum [21, 22] and the modification of the electron Larmor precession frequency [6, 23], respectively. On the contrary, in II–VI semiconductors, the Overhauser field is much weaker due to a smaller hyperfine coupling strength and a smaller abundance of lattice atoms carrying nuclear spins. Experimentally, the nuclear field has been determined to be of the order of a few 10 mT only [24, 25].

As a result of the hyperfine coupling with a nuclear spin ensemble, localized electron spins can be affected in different aspects. For example, in the case of a finite nuclear spin polarization, the electrons can experience a nuclear field, namely the Overhauser field [26], and the electron spin relaxation can be strongly influenced by nuclear spin fluctuations, in particular in QDs [27]. Therefore, by either manipulating the Overhauser field or by changing the hyperfine flip-flop rate of the electron–nuclear spin system, the electron spin can be manipulated. These two aspects
are addressed here for two model systems, n-GaAs and CdSe/ZnSe QDs, respectively, using an onchip microcoil technique for electrical manipulation and pulsed laser techniques for optical read-out.

In the first type of experiment, the Overhauser field is dynamically monitored by optically detecting the electron Larmor precession frequency by means of time-resolved Kerr rotation (TRKR) in n-GaAs [28–30]. By electrically generating a resonant RF magnetic field using an onchip microcoil, the nuclear field can be coherently manipulated [31, 32]. As a consequence, the electron spin precession frequency is modified. In the second type of experiments performed on CdSe/ZnSe QDs [33–35], the magnetic field of the microcoil is switched on and off on a sub-ns time scale to decouple/couple the electron and nuclear spin systems, and to achieve an equalization of the electron spin with the average nuclear spin within the time scale of a single spin flip-flop, respectively. The latter offers the unique possibility to read out the nuclear spin optically [13]. In both experiments, we have achieved sub-ns electrical control and fast optical read out of the semiconductor electron–nuclear spin system.

2 Samples and experimental approach

2.1 Layer sequence of the samples

The Si-doped n-GaAs has been grown with molecular beam epitaxy on a GaAs (100) substrate. The active epitaxial GaAs layer has a thickness of 2 μm. An electron density of $5 \times 10^{16}$ cm$^{-3}$ with a mobility of $\mu = 4115$ cm$^2$ V$^{-1}$ s$^{-1}$ is obtained at room temperature. The active region is embedded between an undoped 50 nm thick GaAs layer below and a Si-doped GaAs layer above, which is gradually doped up to an electron density of $5 \times 10^{18}$ cm$^{-3}$ within 15 nm and covered by 15 nm GaAs with a constant doping level of $5 \times 10^{18}$ cm$^{-3}$. The doping sequence avoids band-bending effects and ensures that the nuclear spins probed experience a macroscopically homogeneous electrical environment [32].

The CdSe QDs have been grown by molecular beam epitaxy on a several hundred nm thick, relaxed ZnSe buffer on top of a [001]-oriented GaAs substrate. To achieve a coherent Frank–van-der-Merve growth of up to three monolayers of CdSe, the temperature has to be lowered by around 100 K compared to the ZnSe growth. During a thermal activation, the pseudomorphic film undergoes a striking transformation into a Stranski–Krastanov-like QD morphology [36, 37]. For the present studies, the structures are capped by a 50 nm thick ZnSe layer. The size of the lens-shaped CdSe islands is below 2 nm in height and below 10 nm in lateral extension with a core of practically 100% CdSe [38]. Under appropriate stoichiometry conditions, n-doping of the ZnSe barrier is achieved where each QD captures about one electron on average.

2.2 Microcoil preparation

The technique for generating an onchip magnetic field via microcoils requires a two-step electron beam lithography and lift-off technique. Single-turn gold microcoils with variable aperture sizes ranging from 3 to 20 μm are defined on top of the semiconductors [11, 39]. Because of the small coil length scale, a low complex impedance is achieved. This permits a switchable magnetic field up to tens of mT with a transition time less than 400 ps [12–14].

A schematic sample layout is shown in the upper part of Fig. 1a, whereas the lower part of the figure shows a scanning electron beam micrograph of a typical microcoil used in this work. The structure dimension of a micrometer length scale is chosen as the best compromise for achieving sufficiently large local magnetic fields, allowing optical access using a focused laser beam, and avoiding metal-induced strain in the area under study. For a current flow through the microcoil, the magnetic field is dominant in the out-of-plane direction, while the inplane field is negligible in the microcoil center. From numerical simulations (see Fig. 1b and c), the z-component of the current-induced magnetic field is $\sim 10$ mT for $I = 120$ mA in the coil center.

Table 1 Hyperfine parameters in QDs: $I$ – nuclear spin, $a$ – isotope abundance, $A_{hf}$ – hyperfine coupling constants, $N$ – number of nuclei with magnetic moment seen by the electron.

|            | $I$ | $a$        | $A_{hf}$ (μeV) | $N$     |
|------------|-----|------------|----------------|---------|
| GaAs       | 3/2 | 100% (As)  | 100            | 10$^3$–10$^6$ |
|            | 3/2 | 60% (60Ga) | 100            | 10$^3$–10$^6$ |
|            |     | 40% (60Ga) | 100            | 10$^3$–10$^6$ |
| CdSe       | 1/2 | 25% (Cd)   | $-10$          | 10$^2$  |
|            | 1/2 | 7% (Se)    | $-10$          | 10$^2$  |
| ZnO        | 5/2 | 4% (Zn)    | 6              | some tens |

Figure 1 (a) Schematic description of an Au-microcoil on a semiconductor. In the lower part of the figure, a scanning electron micrograph of a microcoil structure is shown. (b) Numerical simulation of the out-of-plane component $B_{1z}$ of the current-induced magnetic field; (c) Numerical simulation for the inplane component $B_{1x}$. For the simulations, the following numbers have been used: $I = 120$ mA, coil inner diameter: $8.5$ μm, coil width: $5.5$ μm, depth: $-70$ nm below the semiconductor surface.
while the $x$-component of the current-induced magnetic field is almost negligible there.

### 2.3 Optical measurements

A change of the Overhauser field $B_N$ results in a variation of the electron Larmor precession frequency in GaAs, which can be monitored by TRKR. In order to generate a noticeable Overhauser field, circularly polarized light is used to excite the semiconductor in oblique geometry, i.e., the pump beam is tilted from the in-plane magnetic field direction, as indicated in Fig. 2a [17, 26].

The sample is mounted strain-free in a liquid helium flow cryostat and cooled down to a temperature of $5 \text{ K}$. A static external in-plane magnetic field of $B_{\text{ext}} = 374 \text{ mT}$ is applied. Spin-polarized electrons are generated by a circularly polarized pump beam, which deviates by about $\theta_1 \sim 10^\circ$ from normal incidence. The coherent electron spin precession is probed by a linearly polarized beam perpendicular to the sample surface. The Larmor precession frequency $\omega_f$ is given by the total field $B_{\text{tot}} = B_{\text{ext}} + B_N$. The spin dynamics is traced via the polar Kerr rotation of the probe beam detecting the $S_z$ component of the electron spin $S$. The helicity of the pump beam is modulated by an electro-optic modulator with a repetition rate of $200 \text{ kHz}$ and 50% duty cycle and the probe beam is modulated by a mechanical chopper with a frequency of $110 \text{ Hz}$ for double lock-in detection.

As an excitation source a tunable mode-locked Ti: sapphire laser with a pulse width of 2 ps, a repetition rate of $76 \text{ MHz}$, and a photon energy of $1.5166 \text{ eV}$ is used. The power of the pump and the probe beam is 2.9 and 0.55 mW, respectively, focused to less than $10 \text{ m} \mu \text{m}$ inside the microcoil. The microcoil used in the experiments has an inner diameter of $19 \text{ m} \mu \text{m}$ and a strip line width of $6 \text{ m} \mu \text{m}$. The electrical resistance of the microcoil is negligible at low temperatures compared to the matched $50 \Omega$ terminal resistor. As the microcoil is connected between ground and the $50 \Omega$ resistor, the electric potential of the metal stripes is too low to produce a significant electric field, thus avoiding possible nuclear quadrupole (NQ) splitting effects in the studied GaAs active layer by the microcoil. An arbitrary function generator (Tektronix, AFG 3252) is used as a RF source.

The spin dynamics of resident QD electrons in the CdSe/ZnSe system is studied by a special type of time-resolved saturation spectroscopy. Under quasiresonant excitation with circularly polarized light, the hole spin relaxation in the optically excited charged exciton $X^-$ and the subsequent recombination yields an increased depopulation of the initial state with the addressed electron spin. As sketched in the left panel of Fig. 2b, a $\sigma^+$-polarized photon excites the $X^-$ state with a $+3/2$ hole spin state along the right arm and decays after hole spin flip (dotted line) along the left arm pumping the downward oriented electron spin. Depletion of the initial spin-up electron state reduces the optical excitation rate and, by this, the PL signal re-emitted from the QD (right panel).

Excitation is performed with the spectrally tunable output of a dye laser or a selected wavelength of an Ar-ion laser. The time-dependent polarization and intensity is controlled by a combination of a linear polarizer, an acousto-optical modulator, a Pockels cell, both with a rise time of less than $10 \text{ ns}$, and a quarter-wave retardation plate.

The microcoils are driven by electrical pulses with $350 \text{ ps}$ rise/fall time via a coaxial transmission line terminated by a serial connection of a resistor complementing $R_L$ to a total load of $Z = 50 \Omega$. From the reflected electrical signal, a total inductance of about $L_{\text{tot}} \sim 100 \text{ nH}$ can be derived arising mainly from the bonding wires connecting the microcoil to contacts on the sample holder. The chosen sequence of intensity, circular polarization, and current through the microcoils is provided by a combination of pulse generators. The PL of the chosen circular polarization is spectrally dispersed by a monochromator and detected by a combination of fast photomultiplier and different multiscaler units with a time resolution of 1 and $10 \text{ ns}$, respectively. Measurements were performed at $T = 4.3 \text{ K}$ in a He flow cryostat allowing for confocal optical access to the sample with a spatial resolution of $1 \mu \text{m}$. A static magnetic field of up to $5 \text{ T}$ and $100 \text{ mT}$ is alternatively supplied by an OXFORD Micromag system and a Helmholtz coil pair, respectively.

### 3 Results

#### 3.1 Electron–nuclear spin control in n-GaAs

As shown in Fig. 3a, the Larmor frequency $\Omega_L$ strongly changes...
with polarization of the exciting laser and with RF field. For example, an increase from 13.9 to 15.2 GHz is obtained for $\sigma^+$-polarized pumping (upper red curve) as compared to unpolarized excitation (black curve). Obviously, a nuclear field $B_N$ parallel to $B_{ext}$ is built up for $\sigma^+$-polarized excitation thus increasing the Larmor frequency. If $\sigma^-$-polarized excitation is used for the pump beam instead, the situation is reversed: The Overhauser field $B_N$ should now be anti-parallel to $B_{ext}$, and indeed the Larmor frequency is reduced to 12.5 GHz.

Switching on the RF field by applying a RF current through the microcoil, we expect depolarization of the nuclear spins under resonance conditions. Indeed, choosing a frequency of 2.72 MHz, which is expected to be the resonance frequency for the $^{75}$As isotopes at $B_{ext} = 374$ mT, a distinct change of the Larmor frequency can be seen in the TRKR data. For $\sigma^+$-polarized pumping, the Larmor frequency changes to $\omega_l = 14.5$ GHz (lower red curve in Fig. 3a), whereas for $\sigma^-$-polarized excitation, a Larmor frequency of 13.1 GHz is found (upper blue curve in Fig. 3a). This indicates a significant depolarization of the nuclear spins by the RF field. The Larmor frequency does not reach the value measured for unpolarized pumping as only the $^{75}$As nuclear spins are depolarized, whereas the nuclei of the $^{69}$Ga and $^{71}$Ga isotopes still keep their polarization [31]. Note, the depolarization level is almost the same for both polarizations, i.e., $|\Delta \omega_l| = 0.7$ GHz for $\sigma^+$-polarized excitation and $|\Delta \omega_l| = 0.6$ GHz for $\sigma^-$-polarized excitation. This means the nuclear-field amplitude is actually independent of the optical excitation helicity. This gives evidence that the equilibrium electron spin is negligible compared with the optically injected electron spin, and thus the nuclear field for unpolarized excitation can be ignored.

TRKR measurements are performed for different RF frequency values as presented in Fig. 3b. The dynamic nuclear polarization (DNP) formation was ensured to be saturated prior to switching on the RF field. For each RF frequency, the RF excitation is always on during the measurements and the TRKR scanning was performed ten times, which totally takes around 10 min. Since the DNP formation time has been found to be about 10 min for three isotope species [29, 31], it is expected that the DNP formation of each isotope is shorter than the total TRKR scanning time. For the measurements, the current amplitude through the microcoil was set to 30 mA.

The NMR spectra are given in terms of the extracted Larmor frequency values, which are plotted in the right panel of Fig. 3b for $\sigma^+$-polarized and $\sigma^-$-polarized optical excitation, respectively. NMR features of all three species of isotopes are observed and the center frequency of each NMR peak follows well the relation of $f_a = \gamma_a B_{ext}$, where $\gamma_a$ is the nuclear gyromagnetic ratio for the isotope $\alpha$ ($^{69}$Ga, $^{71}$Ga, and $^{75}$As). The fact that the NMR spectral linewidth (between 20 and 40 kHz) is much larger than expected due to the dipole–dipole broadening in bulk GaAs [40] suggests resonance saturation conditions [16], i.e., $\gamma^2 B_{ext}^2 T_1 T_2 \gg 1$.

Here, the longitudinal spin relaxation time $T_1$ is of the order of DNP time in the studied case and the transverse spin relaxation time of the order of $T_2 \sim 100$ $\mu$s in GaAs [23].

Decreasing the current amplitude to 0.5 mA (left panel), the resonance linewidth of the $^{75}$As isotope is reduced down to $\sim 7$ kHz due to the much weaker RF excitation, while the maximum $B_N$ reduction does not vary significantly. This indicates that under resonance conditions even at a small RF-field of about $B_{RF} \sim 0.01$ mT (evaluated from the Rabi oscillation results discussed in Fig. 5), the nuclear spin depolarization is much faster than the DNP formation, which in our sample is found to occur on a scale of $\sim 10$ min.

For the resonances of each isotope $\alpha$, the total Overhauser field $B_N$ is reduced by $B_{Na}$, with the sum of $B_{Na}$ for all three isotope species being comparable to $B_N$. Apparently, complete spin depolarization of each specific isotope is achieved in the case of resonance. Comparing the measured ratio $B_{Na}/B_N$ with the theoretically expected one, a clear deviation is obtained. For example, according to Ref. [41], $B_{71}$-Ga is expected to be less than $B_{69}$ Ga in pure GaAs crystals, in contrast to our data. As outlined in Ref. [31], this is attributed to an isotope specific nuclear spin leakage factor $f_{\text{leak}}$. As the spin leakage factor $f_{\text{leak}}$ depends on
the nuclear spin polarization time $\tau_{\text{hf-pol}}(\alpha)$ and the depolarization time constant $\tau_{\text{hf-dep}}(\alpha)$, isotope specific values of $\tau_{\text{hf-pol}}(\alpha)/\tau_{\text{hf-dep}}(\alpha)$ apparently result in a larger leakage factor for the $^{71}$Ga as compared to the $^{69}$Ga isotope, in agreement with earlier results [42].

By varying the frequency of the RF field over a larger range, nuclear spin depolarization is observed at different resonance frequency positions, as can be seen in Fig. 4a. Several types of NMR resonances can be found [32]: (i) fundamental NMR at frequencies of $f_\alpha = \gamma_\alpha B_{\text{ex}}$; (ii) two-spin NMR involving one isotope species at $2f_\alpha$; (iii) two-spin NMR involving different species of isotopes at $(f_\alpha + f_{\alpha_2})$; (iv) half-harmonic NMR at $1/2f_\alpha$.

A theoretical understanding of the nonfundamental NMR can be achieved by considering the nuclear spin Hamiltonian, which is written as [32]:

$$H = H_Z + H_{\text{hf}} + H_{\text{DD}} + H_{\text{NQ}}.$$  

$H_Z$ refers to the Zeeman energy and $H_{\text{hf}}$ corresponds to the hyperfine interaction. $H_{\text{hf}}$ describes the perturbation due to the RF magnetic field, $H_{\text{DD}}$ the dipole–dipole (DD) interaction, and $H_{\text{NQ}}$ the NQ interaction. Both, DD and NQ interaction cause a mixture of nuclear spin states and thus spin transitions with $\Delta f = 2$ become allowed under resonance conditions [17, 23, 43]. In the case of DD interaction, the spins of neighboring nuclei of either the same or different isotopes can couple leading to a NMR peak at $2f_\alpha$ and at $(f_\alpha + f_{\alpha_2})$, respectively. In contrast, NQ interaction only involves one nucleus and consequently induces the $2f_\alpha$ resonance.

The half-harmonic resonance at $1/2 f_\alpha$ is due to two-quanta RF absorption for spin transitions of $\Delta f = 1$, which stems from an oscillating RF field oblique to the nuclear field [44]. Indeed, from our calculations we found that the transverse RF-field component achieves values of the order of 1 mT, which we therefore attribute to be responsible for the $1/2f_\alpha$ NMR.

In Fig. 4b, the variation of the Larmor frequency $\omega_L$ with lab time is plotted for a few selected resonances. Time zero marks the moment of switching on the RF field. All curves can approximately be fitted by a monoeponential decay. For the fundamental resonances, the characteristic time constant $\tau_{\text{NSR}}(\text{exp})$ is much shorter than 1 min and limited by the time needed for measurement of the TRKR curve. The nuclear spin depolarization at the fundamental resonance is expected to be on the order of 100 $\mu$s, determined by the nuclear spin–spin relaxation time [16]. This is in good agreement with the dephasing time found in our Rabi-type of experiments (see Fig. 5). For both, the $2f_\alpha$ and the $1/2f_\alpha$ resonances, $\tau_{\text{NSR}}(\text{exp})$ is found to be about 1 min, while for the $(f_\alpha + f_{\alpha_2})$ resonances, time constants of a few minutes are obtained. The nuclear spin relaxation rate is here mainly determined by the depolarization rate from resonant RF absorption [32]. Due to its smaller perturbation strength (approximately 0.01 mT or less in GaAs [43]), the DD-induced spin depolarization rate is much smaller than the one caused by the transverse RF field component and the NQ interaction, respectively.

Figure 5 shows Rabi oscillation measurements performed for the $^{75}$As nuclei. To suppress the current-generated lateral RF field inhomogeneity [11], a pinhole was used for probing the spin dynamics within a range of less than 5 $\mu$m in diameter in the center of the microcoil. As the

![Figure 4](http://www.pss-b.com)

**Figure 4** (a) Optically detected NMR under $\sigma^+\chi$-polarized optical excitation (balls) and $\sigma^-\chi$-polarized optical excitation (squares). The horizontal dashed line indicates the Larmor frequency measured under unpolarized excitation. (b) Larmor frequency versus lab time for different resonances (symbols) after switching on the RF field. The solid lines are monoeponential fits. [32]

![Figure 5](http://www.pss-b.com)

**Figure 5** (a) 2D plot of the TRKR data for $^{75}$As Rabi oscillations. The upper panel is obtained with $\sigma^+\chi$-polarized excitation and the lower panel is obtained with $\sigma^-\chi$-polarized excitation. The RF current amplitude is 30 mA at a frequency of 2.72 MHz. (b) Extracted Larmor frequency versus RF pulse width [31]. The extracted $\omega_L$ data are fitted by damped cosine functions (lines). The RF pulse is indicated in the inset.
external magnetic field energetically splits the nuclear spin states, the RF pulse triggers a coherent absorption-emission cycle between the nuclear spin states of the $^{75}\text{As}$ nuclei [16], allowing a coherent control of the nuclear magnetization direction.

In Fig. 5a, TRKR data are depicted in a color-coded plot versus time delay for different RF pulse widths $\tau_{RF}$ changing along the vertical axis. Pronounced phase oscillations of the Kerr rotation signal can be seen by varying $\tau_{RF}$, under either $\sigma^+$-polarized excitation or $\sigma^-$-polarized excitation. Fig. 5b presents the extracted Larmor frequency versus $\tau_{RF}$. By fitting the data with damped cosine functions, it could be determined that the $^{75}\text{As}$ nuclear magnetization coherently nutates with a Rabi frequency of $\gamma_{Rabi}^{75}\text{As} \approx 4\text{kHz}$. On applying a $\pi$ pulse ($\sim 100\mu s$), a sign reversal of the nuclear magnetization is achieved. The effective dephasing time $\tau_{\text{Rabi}}^{75}\text{As}$ is found to be $\sim 200\mu s$, in good agreement with the literature [23].

3.2 Electron–nuclear spin control in CdSe/ZnSe quantum dots Static manipulations of the resident electron spin dynamics in charged CdSe QDs has been previously demonstrated on CdSe/ZnSe QDs [35] by external magnetic fields with field strengths $B$ in the range of a few tens of mT. Such fields can be produced at reasonable current levels by our miniaturized coils introduced above. Figure 6 compares the action of a static external magnetic field $B_z$ and a static current $I_{\text{coil}}$ through the coil with an inner diameter of 5.5 $\mu m$ and an outer diameter of 9.9 $\mu m$. The time-resolved $\sigma^+$-polarized PL is measured for a cycle of alternating $\sigma^+$ and $\sigma^+$-polarized excitation with $\mu m$ spatial resolution at the center of the coil. Equally long chosen pumping intervals $t_{\text{up}}$ and $t_{\text{down}}$ as well as a total cycle duration of a few $\mu s$ ensures that the spin pumping does not result in a formation of a dynamical nuclear polarization [34] here. Without DNP, the PL transient sketched in Fig. 2b reflects directly the depopulation of the respective resident electron spin state and the amplitude $A$ is related to the change of the nonequilibrium spin during the oppositely circular polarized excitation. The observed increase of $A$ with growing $B_z$ is a consequence of the slowing down of the hyperfine flip-flop rate with magnetic field found in CdSe QDs [35] with the relation $\tau_{\text{hf}}(B) = 0.6\mu s [1 + 8B^2/(mT)^2]$. Since the $A(B_z)$ and $A(I_{\text{coil}})$ dependencies are measured on the same QD ensemble in the coil center under identical conditions, the data in Fig. 6 can be used to determine the longitudinal field strength generated by the coil at the given $I_{\text{coil}}$. A value of up to 10 mT is obtained for the highest current applied in very good agreement with the calculated field component in Fig. 1b. The equal increase of $A$ independent of the direction of $I_{\text{coil}}$ is a consequence of the quadratic field dependence of $\tau_{\text{hf}}$.

The main advantage of the microcoil compared to a conventional solenoid is the ability to control the field strength on a sub-ns time scale. This is demonstrated in Fig. 7, where an optically pumped nonequilibrium electron spin state is erased by switching on the coupling to the nuclear spin system via hyperfine interaction. The appropriate sequence of optical and current pulses is schematized in the upper panel. The PL rise at $t=0$ is caused by the switch to co-polarized excitation with a preferential spin-up orientation from the prior $\sigma^+$ pump interval followed by a partial signal decay due to spin reorientation to the down alignment. A significant spin polarization in both orientations is achieved in the presence of a current-induced magnetic field of $B_z = 40$ mT. Please note that the switching from $\sigma^- \rightarrow \sigma^+$ excitation changes the finally reached electron spin from $S_{\text{eq}}$ to $-S_{\text{eq}}$ so that the total absolute change is $2S_{\text{eq}}$. After $I_{\text{coil}}$ is switched off at about $t=1\mu s$, $\tau_{\text{hf}}$ becomes so short that the optically induced electron spin polarization drops completely. This is caused by the drastic increase of the hyperfine flip-flop rate if the magnetic field is switched off.

![Figure 6](image6.png)  
Figure 6 Amplitude $A$ versus field $B_z$ generated by a superconducting magnet (circles) and versus $I_{\text{coil}}$ in both current directions (up and down triangles) [13]. $t_{\text{up}} = t_{\text{down}} = 2.5\mu s$.

![Figure 7](image7.png)  
Figure 7 Optical pumping of the electron spin at $I_{\text{coil}} = 100$ mA and erasure via hyperfine interaction at vanishing $I_{\text{coil}}$ [13]. Upper panel: schematics of the sequence of optical excitation polarization degree $\mu_{\text{exc}}$ and current pulse. A microcoil with 2.7 and 6.3 $\mu m$ inner and outer diameter, respectively, was used, producing a large $B_z$ of up to 40 mT. Lower panel: PL transients observed.
The degree of spin erasure is measured experimentally by $\eta = \Delta A/A$, where $A$ is the amplitude in the presence of $B_z$ here. The experimental value of $\eta = 0.52 \pm 0.08$ is within the accuracy of the setup equal to the expected value for a complete erasure, $\eta = 1/2$. The time constant of electron spin erasure seen at the $\Delta A$ switch is limited by the time-resolution of the photon detection system of about 10 ns proving that the switching time of the magnetic field is shorter. From the total inductance of microcoil and bonding and the load resistance, a reasonable switching time of $\tau_B = L_{\text{coil}}/Z = 2$ ns follows for the experiments here. An improved layout of the wiring will speed up the erasure time to get close to the limit set by the single electron spin flip-flop time $\tau_{sf}/N$.

It should be emphasized that the result in Fig. 7 represents the situation for vanishing DNP, because the hyperfine interaction really equalizes the electron spin $S$ and the average nuclear spin $J$ [13], which in this case is zero due to the equally long pumping intervals $t_{\sigma^-}$ and $t_{\sigma^+}$ and the total cycle duration of only a few $\mu$s. In fact, the combination of optical and electrical pulses can be utilized to read-out the nuclear spin state optically. To achieve a given DNP, an asymmetric optical pumping with a varying ratio between $t_{\sigma^-}$ and $t_{\sigma^+}$, but constant total cycle time, is applied as sketched in the upper part of Fig. 8. While the net nuclear polarization produced in a single cycle is negligibly small, the repetitive application of a sufficiently large number of periods generates an appreciable steady-state value of $J$.

The sign of $J$ itself is determined by the sign of $t_{\sigma^+} - t_{\sigma^-}$. Experimentally, some $10^5$ cycles are awaited before the PL transient is taken. Now, switching off $I_{\text{coil}}$ and by this the magnetic field, the electron spin does not recover to zero but to $S = J$ with vanishing hyperfine flip-flop net rate. The resulting electron spin state, and therefore the uncovered nuclear spin, is manifested by a respective value of $\Delta A$. The normalized change $\Delta A/A$ is larger than 0.5 for an electron spin orientation created by optical pumping opposite to that of the nuclear spin and smaller than 0.5 in the opposite case. As shown in lower part of Fig. 8, the experimentally observed dependence follows this expectation.

3.3 Spin effects in ZnO In the present context, ZnO-related quantum structures have recently come into consideration. Compared to sulfides and selenides, oxides exhibit a drastically reduced spin–orbit (SO) coupling [45, 46] giving rise to the expectation that ZnO represents an interesting material for spintronics or even quantum computing. The prediction of room-temperature ferromagnetism for ZnO doped with magnetic transition metal ions [47] nearly 10 years ago triggered an ongoing research in this field. The advantage of a reduced SO interaction offers the principal possibility of enlarged electron and hole spin coherence time. For example, only the D’yakonov–Perel mechanism is expected to contribute significantly to the electron spin relaxation in the case of n-type ZnO [48, 49], since the Bir–Aronov–Pikus mechanism is absent due to missing holes and the Elliot–Yafet mechanism is strongly weakened by the combination of large bandgap energy and small SO coupling, respectively. Recently, the growth of ZnO/(Zn,Mg)O multiple quantum wells (MQWs) with an inhomogeneous broadening of clearly below 10 meV has become possible, where the charged exciton transition could be unambiguously identified by magneto-optical spectroscopy [50].

The small SO interaction can imply itself an important drawback. Any optical spin pumping or probing via interband transitions are based on the SO coupling fixing the hole spin with the orbital part of the wave function that is responsible for the selection rules of interaction with properly polarized light. Indeed, a marginal inplane strain component found in MQWs grown along the c-axis on sapphire results in a complete reorganization of the spin–orbit coupled states. As a result, the absorption and the PL intensity depend on the linear polarization parallel and perpendicular to the substrate c-axis [51]. Calculations using the Bir–Pikus Hamiltonian have shown that the $X^-$ states with half integer spin remain Kramers-degenerate coupled now to oppositely elliptically polarized light instead of oppositely circularly polarized one. For the free excitons with integer spin, the degeneracy of the optically allowed $\Gamma_5$ states becomes completely lifted and the states do not belong longer to pure spin states [51].

With the use of a c-plane sapphire substrate, such modifications can be avoided. The confirmation of $X^-$ transitions makes the optical spin pumping of a two-
dimensional electron gas in a QW achievable, in a similar way as it should be in principle possible for the donor electrons via the neutral bound exciton D\\(^3\)X. Time-resolved studies on optical orientation on both X\(^-\) in a MQW and D\\(^3\)X in a ZnO epilayer have shown that the longitudinal hole spin relaxation t\(^{\text{h,s}}\) amounts to \(\sim 140\) and \(\sim 100\) ps, respectively, somewhat shorter than the lifetime \(\tau_0\) of the respective exciton complexes. In addition, TRKR experiments exhibit an electron spin coherence time \(T_2^*\) of up to several ns in a MQW, if excited at the neutral exciton resonance. The D\\(^3\)X data are consistent with other studies \[52, 53\] and the observation of quite similar t\(^{\text{h,s}}\) in the bulk and QW situation seems to be characteristic for ZnO structures. In any case, the relation t\(^{\text{h,s}}\) \(\approx \tau_0\) ensures an effective spin pumping of resident electrons according to the scheme in Fig. 2b. Strikingly, we found that an external field strength of only about 2 mT is sufficient to suppress the hyperfine flip-flops. It can be concluded that the Overhauser field in ZnO is more than 1 order of magnitude smaller than in CdSe, in accordance with Table 1, being a consequence of the reduced abundance of magnetic moment carrying nuclei and the reduced volume of exciton complexes. The data prove that studies with lithographically made microcoils on ZnO-related structures may be of high interest, especially with respect to spintronic effects.

4 Conclusions In conclusion, an electrical control of the electron–nuclear spin system has been demonstrated by using lithographically defined microcoils. The onchip microcoil technique allows one electrically to manipulate the nuclear spins by means of NMR in n-GaAs and to control the hyperfine flip-flop in CdSe/ZnSe QDs, so that an electrical control of either the electron spin precession or the electron spin relaxation is achieved. The results underline the potential of such microcoils for ultrafast (sub-ns) electrical spin control and spin manipulation as a general method in semiconductor nano-/microstructures. Further improvements regarding the electrical and geometrical design will make it possible to achieve even shorter switching times and to access single QDs.

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