Spin transport and spin manipulation in GaAs (110) and (111) quantum wells

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Received 15 September 2013, revised 17 December 2013, accepted 9 January 2014
Published online 14 February 2014

Keywords acoustic transport, GaAs quantum wells, spin dynamics, spin–orbit-interaction

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1 Introduction

The manipulation of spins in semiconductor materials has become an active area of investigation, in particular after different proposals for spin-based electronic information processing have been put forward [1–3]. Device functionalities based on electron spins require processes for the generation, storage, detection, and transport of spins as well as interaction mechanisms to efficiently manipulate the spin vector. Advanced applications further presuppose the interaction between two spins in order to realize spin–spin gates. In this respect, two main challenges are (i) the enhancement of the electron spin coherence times in intrinsic III–V bulk semiconductors, which is typically of only one ns, thus restricting the number of spin manipulation steps that can be realized before decoherence effects set in, and (ii) the development of spin manipulation techniques that do not compromise the spin lifetime.

While isolated electron spins can be efficiently manipulated using a magnetic field, electron spins in a crystal interact with other spins from electrons (the exchange interaction) or nuclei (the hyperfine interaction) as well as with the lattice potential. These interactions can be used for the realization of spin control gates, but at the same time they may also lead to spin relaxation if not appropriately controlled. The electron–hole exchange interaction leads to the so-called Bir–Aronov–Pikus (BAP) spin dephasing mechanism [4], which is particularly important for excitons [5] as well as in highly p-type doped III–V semiconductors at low temperatures [6]. Spin scattering via the hyperfine-interaction is normally negligible for free carriers, but becomes important for localized electrons in quantum dots.

In this contribution, we investigate spin transport and manipulation in III–V semiconductor quantum well (QW) structures grown along non-conventional crystallographic directions, such as the [110] and [111] directions. The motivation for these studies arises from the special characteristics of the interaction of spins with the lattice potential in these structures, the well-known spin–orbit (SO) interaction. The SO-interaction arises from the fact that an electron experiences a varying electric field while moving through a noncentrosymmetric crystal (such as the zinc-blende...
semiconductors). The latter translates into a momentum ($\hbar \mathbf{k}$) dependent effective magnetic field $\mathbf{B}_{\text{SO}}$ in the electron reference frame, which acts on its spin. The SO-coupling can be controlled by external electric [1, 7] and strain fields [8]. As a result, it provides an interesting mechanism for spin generation and detection (i.e., optical orientation [9]) as well as for spin manipulation without the application of magnetic fields. Examples are the generation of non-equilibrium spin populations using electric currents [10–14], as well as the spin-galvanic [15, 16] and spin Hall effects [13, 17, 18].

The SO-interaction can also lead to spin dephasing in an electron ensemble, since electrons with different $\mathbf{k}$s will experience SO-fields of different strengths and directions. Their spins will then precess at different rates, leading to a reduction of the resulting spin ensemble, an effect known as the Dyakonov–Perel (DP) spin dephasing mechanism [19–21]. The SO-interaction can also couple different spin states during electron scattering processes. One example is the Elliott-Yafet spin dephasing mechanism [22], which describes spin-flip transitions induced during electron scattering by impurities or phonons and is expected to be influential in highly doped, low-bandgap materials. Finally, the SO-coupling is also behind the intersubband spin relaxation mechanism [23] in QW structures, where electron scattering between two subbands is accompanied by a spin flip.

Two main approaches have been proposed to control the relaxation and to manipulate electron spins in semiconductors. The first, which will be discussed in detail in Section 2, exploits the dependence of the SO-interaction on the symmetry of QW structures. The second relies on the use of confinement potentials of quantum wires and dots (QDs). Confinement isolates spins from the semiconductor matrix, thus ensuring the long spin coherence times required for the coherent manipulation by external fields [24, 25]. Confinement can also be efficiently employed to control SO-effects via motional narrowing. The latter relies on the fact that the DP spin scattering rate is inversely proportional to the momentum scattering rate [19]: frequent momentum scattering at the borders of the confinement potential randomizes the effective SO-field $\mathbf{B}_{\text{SO}}$ experienced by the spins, thus leading to longer lifetimes [26, 27]. The required confinement dimensions are dictated by the SO length $\ell_{\text{SO}}$, which is defined as the typical distance required for a one rad spin precession under $\mathbf{B}_{\text{SO}}$. $\ell_{\text{SO}}$ can reach a few $\mu$m in wide GaAs QWs, thus making it possible to employ mesoscopic confinement potentials for spin control. Finally, no DP relaxation is expected in one-dimensional systems (1D) since the axis of the fictive SO-field $\mathbf{B}_{\text{SO}}$ is fixed and its direction reversible with $\mathbf{k}$ [28, 29].

The highest degree of spin control and manipulation has so far been achieved in QDs defined either by Stranski–Krastanov growth or by metal gates [25]. In these systems, single electron spins in QDs have been manipulated using a variety of mechanisms including the electrical control of the exchange interaction [30] as well as electron spin resonance (ESR) induced by a varying magnetic field induced by a strip-line [31] or by a periodically varying $\mathbf{B}_{\text{SO}}$ [32]. In the last case, the effective magnetic field associated with the SO-interaction can be varied through an oscillatory spin motion induced by an electric field [33, 34]. Proposals for coherent spin control using ESR generated through carrier motion in a magnetic field gradient are also available [35, 36].

A further challenge toward spin-based information processing schemes is the realization of scalable systems consisting of several spin subsystems. These processes normally require the transport of spins, which cannot be easily combined with 3D-confinement. In this contribution, we explore the use of moving piezoelectric potentials created by surface acoustic waves (SAWs) for spin transport and manipulation. These moving fields provide a way of combining the advantage of confinement with the long range transport required to couple remote spin systems.

In this work, we review recent results on the control of the SO-interaction and acoustic transport in GaAs(110) and (111) QWs. We start in Section 2 with a theoretical description of SO-effects arising from the symmetry of these QWs, and how they can be controlled using spatial confinement or external electric and strain fields. In (110) QWs, the lifetime enhancement associated with the crystallographic symmetry is restricted to spins aligned with the growth direction [23, 37]. Higher lifetimes for other orientations can be achieved by lateral confinement. (111) QWs are particularly interesting because SO-effects can be suppressed for all spin orientations via the application of a vertical (i.e., perpendicular to the QW plane) electric field [38, 39]. Experimental exploitation of the non-conventional QW orientations requires the development of epitaxial growth procedures to create (Al,Ga)As structures with quality similar to the conventional (001) ones. Recent results about the growth via MBE of high quality QWs as well as QWs embedded in microcavities are presented in Section 3.

The experimental studies of the spin dynamics in the QWs were carried out by combining spectroscopic techniques with spin transport by SAWs. The use of SAWs for the transport of carriers and spins in (Al,Ga)As(001) structures is reviewed in Section 4. The application of SAWs for the transport of spins in (110) QWs is discussed in Section 5. Emphasis is placed on the combination of lateral confinement with acoustic transport in order to achieve efficient spin transport at temperatures above 100 K. Section 6 addresses the spin dynamics in GaAs(111) QWs. Here, we demonstrated that an electric field applied across the QW can suppress SO-effects, leading to lifetimes exceeding 100 ns. Approaches for acoustic spin transport are discussed in Section 7, which also summarizes the main conclusions of this work.

2 Symmetry effects on the spin–orbit interaction

2.1 The Dresselhaus contribution The SO-interaction in III–V QWs is governed by two major contributions. The first one is of intrinsic nature and associated with the bulk inversion asymmetry (BIA) of the III–V zinc-blende lattice. The effective SO-magnetic field $\mathbf{B}_{\text{SO}}(\mathbf{k}) = \hbar \Omega_{\text{SO}} / (g \mu_B)$ associated with this contribution, which is normally denoted as the Dresselhaus term, can be expressed in...
induced by the bulk inversion asymmetry (Dresselhaus contribution [43], \(\hbar \Omega_D\)) and by structural inversion anisotropies created by the application of a vertical electric field \(E_z\) (Rashba contribution [44], \(\hbar \Omega_R\)). Also listed are the SIA contributions due to the biaxial strain induced by lattice mismatch (\(\hbar \Omega_{d}\)) and to the strain field of a Rayleigh SAW (\(\hbar \Omega_{SAW}\)). The strain-induced contribution of BIA-type [41], which is considered to be small, is not taken into account. \(d_{\text{eff}}\) denotes the effective QW thickness, while \(\gamma, r_{41}\), and \(C_3\) are constants describing the dependence of the spin splittings due to the Dresselhaus, Rashba, and strain contributions. The parameter \(\gamma\) relates the splitting between the energetic spin eigenstates in the bulk III–V material to the electron wave vector (cf. Ref. [45]). The parameter \(r_{41}\) for the Rashba term follows the convention in Winkler [46]. For the biaxial strain, the non-vanishing strain components are assumed to be \(u_{xz}, u_{yz} = u_{xx}\), and \(u_{zz}\). The SAW calculations assume that electrons propagate with a wave vector \(k_s = m_s v_{\text{SAW}} / \hbar\) (where \(m_s\) and \(v_{\text{SAW}}\) denote the electron mass and the SAW propagation velocity, respectively) under the strain field of a Rayleigh SAW, which have non-vanishing strain components \(u_{xz}, u_{zz}, u_{zz}\) (see Section 4 for details).

| structure | (001) QW | (110) QW | (111) QW |
|-----------|----------|----------|----------|
| x-axis    | [110]    | [001]    | [112]    |
| y-axis    | [110]    | [110]    | [110]    |
| z-axis    | [001]    | [110]    | [111]    |

| Dresselhaus (\(\hbar \Omega_D\)) | \(\gamma \left(\frac{\hbar}{2d_{\text{eff}}}\right)^2\) \begin{bmatrix} k_y & k_z \\ k_x & \end{bmatrix} | \(\frac{1}{2} \gamma \left(\frac{\hbar}{2d_{\text{eff}}}\right)^2\) \begin{bmatrix} 0 \\ k_y \end{bmatrix} | \(\frac{1}{2} \gamma \left(\frac{\hbar}{2d_{\text{eff}}}\right)^2\) \begin{bmatrix} k_x & \end{bmatrix} |
| Rashba (\(\hbar \Omega_R\)) | \(2E_z r_{41}\) \begin{bmatrix} k_y \\ -k_x \end{bmatrix} | \(2E_z r_{41}\) \begin{bmatrix} k_y \\ -k_x \end{bmatrix} | \(2E_z r_{41}\) \begin{bmatrix} k_x \end{bmatrix} |
| biaxial\(^a\) strain\(^b\) (\(\hbar \Omega_{B}\)) | \begin{bmatrix} 0 \\ 0 \end{bmatrix} | \(\frac{1}{2} C_3 (u_{zz} - u_{xx}) \begin{bmatrix} 0 \\ k_y \end{bmatrix} \end{bmatrix} | \(\frac{1}{2} C_3 (u_{zz} - u_{xx}) \begin{bmatrix} -k_y \end{bmatrix} \end{bmatrix} |
| SAW\(^c\) strain\(^b\) (\(\hbar \Omega_{SAW}\)) | \(\frac{1}{2} C_3 u_{xx} \begin{bmatrix} -k_x \end{bmatrix} \end{bmatrix} | \(\frac{1}{2} C_3 u_{xz} \begin{bmatrix} 2k_x \\ 0 \end{bmatrix} \end{bmatrix} | \(\frac{1}{2} C_3 (u_{zz} - 2u_{xz}) \begin{bmatrix} 0 \\ k_y \end{bmatrix} \end{bmatrix} |

\(^a\) Calculated for a biaxial strain with in-plane component \(u_{yz} = u_{xy}\) and out-of-plane component \(u_{zz}\);

\(^b\) Calculated including only the SIA-type strain contribution [41, 42];

\(^c\) The SAW propagates along the \(x\) direction (cf. Section 4).

Terms of a wave vector-dependent spin-splitting energy \(\hbar \Omega_D\). Here, \(\mu_B\) is the Bohr magneton and \(g\) the electron g-factor. Expressions for \(\Omega_D(k)\) are summarized for QWs with different growth directions in Table 1. \(\Omega_D(k)\) depends on the spatial extent \(d_{\text{eff}}\) of the electronic wave function along the growth direction, as well as on the spin-splitting parameter \(\gamma\) for the QW material. The expressions listed in Table 1 assume the axes listed in the first row of the table (where \(x\) and \(y\) lie in the QW plane and \(z\) is along the growth direction) and apply for small \(k\)-vectors (i.e., \(k_x, k_y \ll \pi/d_{\text{eff}}\)).

The first row of plots in Fig. 1 sketches the orientation of \(B_0(k)\) for QWs grown along different crystallographic directions. \(B_0(k)\) lies in the QW plane for (001) and (111) QWs, its amplitude and orientation depending on \(k\). For (110) QWs, in contrast, \(B_0(k)\) has a single component along \(z\) that only depends on \(k_z\). As a result, DP relaxation does not affect \(z\)-oriented spins, thus leading to long spin lifetimes [23, 37, 40]. The \(\hbar \Omega_D\) component along \(z\) leads, however, to short relaxation times for spins oriented in the \(xy\) plane, thus preventing efficient spin manipulation. The lifetime of these spins can be enhanced by using lateral confinement to force the carriers to move along a single direction. Note, in particular, that \(B_{SO}\) only depends on the \(\hat{y}\) component of the carrier momentum and vanishes for \(k_z = 0\). As a result, the SO-field vanishes for transport channels along the \(\hat{z}\) direction. This effect will be explored for spin transport at high temperature to be described in Section 5.2.

![Figure 1](https://www.pss-b.com)
2.2 The Rashba contribution

The second important contribution to the SO-interaction arises from structural inversion asymmetries (SIA) in the QW potential introduced by an external field along z. In most cases, the SIA is generated by an electric field \( E_z \) perpendicular to the QW plane, leading to the so-called Rashba SO-contribution [44]. Expressions for this contribution to the SO-precession frequency, \( \Omega_R(k) \), in QWs of different symmetry are listed in the 3rd row of Table 1. \( \Omega_R(k) \) lies always in the QW plane and perpendicular to \( \mathbf{k} \). Since its strength can be electrically controlled, the Rashba effect provides a powerful approach for the dynamic manipulation of moving spins [1].

In (110) QWs, the Rashba effect can be used to rotate spins around axes perpendicular to \( z \) [47]. For the other orientations, the Rashba field can also compensate the BIA contributions. In (001) QWs, \( \Omega_D(k) + \Omega_R(k) \) vanishes for spins propagating along the \( x \)-direction ([110] in cartesian coordinates) provided that [48]

\[
E_z = \frac{1}{2} \frac{\gamma}{r_{41}} \left( \frac{\pi}{d_{\text{eff}}} \right)^2, \quad \text{(001) QWs. (1)}
\]

A similar expression, but with a negative sign, applies for \( k_x \)[[110]. Under compensation, long-living spin eigenstates in the form of a persistent spin helix exist, as recently verified by different groups [49–51].

(111) QWs are particularly interesting for the investigation of SO-effects because the Rashba and Dresselhaus contributions have the same symmetry (cf. Fig. 1). As a result, the total SO-interaction precession field \( \Omega_0(k) + \Omega_R(k) \) vanishes for all (small) \( k \)-vectors by selecting an electric field given by

\[
E_x = -\frac{1}{\sqrt{3}} \frac{\gamma}{r_{41}} \left( \frac{\pi}{d_{\text{eff}}} \right)^2, \quad \text{(111) QWs. (2)}
\]

This compensation mechanism was originally proposed in the theoretical works by Cartoixa et al. [38] and Vurgaftman and Meyer [39]. Numerical calculations have also been carried out for doped QWs [52]. Recent experimental verification has been provided by Ballocci et al. [53] and Hernández-Mínguez et al. [54] (see also Refs. [55–57]). These experiments will be reviewed in Section 6.

2.3 Strain effects

In addition to electric fields [1, 7], strain fields also affect the QW symmetry, thereby introducing a SIA contribution to the SO-coupling.\(^1\) Epitaxial layered structures are normally subjected to biaxial strain induced by the mismatch between the lattice parameters of the individual layers. Theoretical predictions for the dynamics of spins in strained QWs have been presented in Refs. [42, 60] – experimental results demonstrating spin precession induced by a strain field in (001) QWs are given in Refs. [8, 59, 61].

Expressions for biaxial strain contribution (\( \hbar \Omega_B \)) in QWs with different orientations are summarized in the 4th row of Table 1. As in the case of the Dresselhaus and Rashba mechanisms, the expressions are only valid for small electron wave vectors \( \mathbf{k} \). In this approximation, \( \Omega_B \) vanishes in (001) QWs. In (110) and (111) QWs, in contrast, \( \Omega_B \) has the same symmetry as the Dresselhaus contribution [42]. Lattice strain induced effects are expected to play a prominent role in (In,Ga)As structures grown on GaAs substrates.

2.4 Higher order terms in \( \mathbf{k} \)

The expressions for the Dresselhaus contribution in Table 1 are valid in the low temperature regime, where the thermal expectation values of \( \langle k_x^2 \rangle, \langle k_y^2 \rangle \ll \langle k_z^2 \rangle = (\pi/d_{\text{ew}})^2 \). Outside this regime, one has to take into account contributions of higher order in \( \mathbf{k} \) in the Dresselhaus term. In (110) QWs and in the absence of many-body effects, these contributions do not change the symmetry properties of the SO-interaction [62]. We will examine here only the case of (111) QWs, where the following correction up to third order terms in \( \mathbf{k} \) has to be added to \( \Omega_0 \) [38]:

\[
\Omega_D^{(3)}(k) = \frac{\gamma}{2\hbar \sqrt{3}} \left( \begin{array}{c} -k_x^2 k_y^2 \\ k_x^2 \end{array} \right), \quad \text{(3)}
\]

being \( k_i = (k_x, k_y) \).

An additional SO-contribution arises from the fact that the GaAs QWs are normally grown on surfaces slightly tilted away from the the [111] direction (see Section 3.1). By calculating the Dresselhaus contribution for this new orientation, one can show that the tilting, \( \delta \theta \), introduces an extra SO-term given by:

\[
\Omega_D^{(\delta \theta)}(k) = \frac{\gamma \delta \theta}{2\hbar \sqrt{3}} \left( \begin{array}{c} \sqrt{2} k_x \left( k_x^2 + k_y^2 - 2 \langle k_z^2 \rangle \right) \\ k_y \left( k_y^2 + 4 \langle k_z^2 \rangle - 4 \langle k_x^2 \rangle \right) \end{array} \right). \quad \text{(4)}
\]

Note that this contribution lifts the degeneracy between the \( x \) and \( y \) components of the precession frequency given by the other Dresselhaus contributions.

The lifetimes \( \tau_i \) for spins oriented along the axis \( i \) (\( i = x, y, z \)) can be estimated from the total precession frequency \( \Omega_{SO}(k) = \Omega_0(k) + \Omega_D^{(3)}(k) + \Omega_D^{(\delta \theta)}(k) \) according to:

\[
\frac{1}{\tau_i} = \tau_p^\ast \left( \langle \Omega_{SO,x}^2(k) \rangle + \langle \Omega_{SO,y}^2(k) \rangle \right) \quad \text{with} \quad i \neq j \neq l = x, y, z.
\]

Here \( \tau_p^\ast \) depends on the momentum scattering time. The thermal averages indicated by \( \langle \ldots \rangle \) can be calculated by taking

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into account that \( k_x^2 = \langle k_x^2 \rangle = \frac{1}{2} k_\perp^2 \), where \( k_\perp^2 = 2m^*kT/\hbar^2 \).
However, a better approximation is obtained integrating the previous equation for \( k_\perp(E) = \left( \frac{2m^*E}{\hbar^2} \right)^{1/2} \), where the electron energy \( E \) is assumed to follow a Boltzmann distribution.

Equation 5 yields the lifetime of spins oriented along a certain crystallographic direction. For spins precessing around an in-plane magnetic field with a short precession period compared with the spin lifetime, the average lifetime \( \tau_R \) becomes [23]

\[
\frac{1}{\tau_R} = \frac{1}{2} \left( \frac{1}{\tau_x} + \frac{1}{\tau_z} \right),
\]

where the magnetic field is oriented along the \( y \) direction. Since \( \tau_z \approx \tau_x \), the same expression also applies for spins precessing around \( x \).

3 Epitaxial growth of GaAs(110) and GaAs(111)B quantum wells As was pointed out in the introduction, QWs grown in other than the most commonly used [001] direction are predicted to be advantageous for electron spin control. From the growth point of view, molecular beam epitaxy (MBE) on GaAs(110) and GaAs(111) appears very interesting and different from growth on GaAs(001). Already a comparison of the surface energies for different substrate orientations suggests that epitaxial growth on (111)B and (110) surfaces is distinctly different from epitaxy on common (001) substrates. The surface energy of the (001) surface can be estimated to be 2.9 J m\(^{-2}\) and is thus much larger than the surface energies derived for the (110) and (111) surfaces, which amount to 2.1 and 1.7 J m\(^{-2}\), respectively [63]. The surface energy can be assumed to be inversely proportional to the lifetime of Ga adatoms on the growing GaAs surface [64]. Therefore, Ga adatoms migrating on the (001) surface are expected to be much more efficiently incorporated into the crystal than on (110) and (111) surfaces.

The previous growth model ignores the distinct surface reconstructions occurring during growth on different GaAs surfaces. Growth on GaAs(001) is usually accomplished under As-rich condition on a 2 \( \times \) 4 reconstructed surface, where As dimers are typically arranged in parallel rows along [\( \overline{1}01 \)], separated by left-out trenches in between [65, 66]. The nucleation process starts on a ridge of As dimers when a trapped Ga atom bonds to another As dimer. This nucleus grows then as a two-dimensional island by capturing further Ga atoms and As dimers, thus filling also the trenches. When this island is big enough, the uppermost As dimers split to form the stable 2 \( \times \) 4 reconstructed surface once again [67]. Contrary, the nucleation behavior on (110) and (111) surfaces can be described by conventional two-dimensional growing islands without the need to consider any site specificity [67]. Thus, the widespread growth on GaAs(001) has to be considered as a well-studied but rather unique scenario, which does not necessarily apply to other substrate orientations.

3.1 MBE growth of (Al,In,Ga)As(110) QWs The GaAs(110) surface is very unique in the sense, that it is the only stable surface of GaAs that does not show any surface reconstruction. It is further characterized by the same number of Ga and As atoms in each monolayer. The GaAs bonds run symmetrically to the (110) plane leading to the non-polar nature of this surface. When homo-epitaxial growth is accomplished on GaAs(110) under the conditions typically used for growth on GaAs(001), strong facetting is observed leading primarily to elongated triangular features with [100] side-facets [68].

During the hetero-epitaxial growth required for the fabrication of QW structures on GaAs(110), effective strain relaxation by dislocations moving on slip planes is restricted to the [001] direction. In contrast, this mechanism is possible along both orthogonal <110> surface directions of hetero-epitaxial layers grown on GaAs(001) [69]. Furthermore, the elastic properties of GaAs(110) and GaAs(001) are different, resulting in a strongly reduced critical thickness for hetero-epitaxial growth on GaAs(110) compared to growth on GaAs(001) [70].

The (110) samples used in this study were grown at a relatively low temperature of 490 °C under high As\(_4\) pressure in order to suppress facetting and to ensure a mirror-like surface. This is illustrated in Fig. 2, where three Nomarski micrographs of the surfaces of nominally identical (Al,Ga)As(110) microcavity structures are shown. The only varied parameter was the As\(_4\) background pressure \( (p_{As_4}) \), which was increased from (a) \( p_{As_4} = 6.0 \times 10^{-7} \) mbar to (c) \( p_{As_4} = 1.4 \times 10^{-6} \) mbar. Only the sample grown under the highest \( p_{As_4} \) shows a mirror-like surface. The reduced mobility of the group-III adatoms due to the low growth temperature and the high As flux seems to promote smooth surfaces on GaAs(110) based heterostructures. Under these conditions, QWs with good optical qualities can be grown. Figure 3 shows a photoluminescence (PL) spectrum taken at \( T = 5 \) K on a sample containing one GaAs QW and three (In,Ga)As/(Al,Ga)As QWs. The thickness of each QW is 20 nm (the In contents \( x \) and the full width at half maximum, FWHM, of the lines are indicated in the plot). The PL line of the GaAs QW is characterized by a narrow linewidth of only 0.5 meV, which is not much broader than the best values of

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about 0.16 meV (cf. Fig. 6) of comparable QWs grown on GaAs(001).

To increase the efficiency of photon to electron–hole pair interconversion during optical probing of the spin dynamics, (110) single QWs were grown embedded in a microcavity structure (cf. Section 5.2). In such cases, all λ/4 layers of thick mirror layers were grown as short-period-super-lattices (SPSLs), as it has been shown that the numerous interfaces of SPSLs hinders the propagation of misfit dislocation through the epilayers [70]. The matching of microcavity resonance wavelength and QW emission wavelength was ensured by applying in situ continuous spectral reflectivity measurements [71].

### 3.2 MBE growth of (Al,Ga)As(111) QWs

The GaAs(111) surface is characterized by its threefold symmetry with alternating (211) and (110) directions separated by 30°. Unfortunately, this threefold symmetry leads to the formation of pyramidal structures on the surface under usual growth conditions [72]. However, mirror-like morphologies without any pyramidal features can be achieved when substrates with a slight off-orientation (≤ 5°) are used [73]. Epilayers with surfaces free of pyramids can also be deposited on exactly oriented GaAs(111)B(= (111)) substrates if the V/IIII flux ratio and growth temperature are accurately chosen [74]. Under these conditions, a static $\sqrt{3} \times \sqrt{3}$ R 23° surface reconstruction can be observed [72, 74–77]. Correspondingly, the (111)B samples used in this study were grown at a temperature of 600 °C on $\sqrt{3} \times \sqrt{3}$ R 23° reconstructed surfaces of slightly off-oriented (1–3°) GaAs(111)B substrates to ensure mirror-like surfaces without pyramidal features.

Due to the threefold symmetry of the (111) surface, opposing directions are not equivalent, an important point to take into account when choosing the direction of the off-orientation. The enormous impact of the direction of the off-orientation on the epitaxial growth is illustrated in Fig. 4. Nomarski micrographs display the morphologies of two nominally identical (Al,Ga)As microcavity structures grown on (111)B substrates with a 3° off-orientation toward (a) (211) and (b) (211). Whereas in the latter case a mirror-like surface is obtained (panel b), an off-orientation toward the opposing direction leads to zig-zag shaped growth steps, which evolve to extended triangular features (as depicted in panel a) when thick structures are grown.

Even when the growth parameters are chosen to yield mirror-like surfaces on slightly off-oriented (111)B substrates, the next challenge is then to adjust the level of step bunching according to the desired application. Although Ren and Nishinaga [78] demonstrated a clear correlation of step bunching with growth temperature, the controlled tuning of step bunching level is experimentally difficult to realize. This is illustrated in Fig. 5, which displays three atomic force micrographs taken at different positions of a 2° substrate. In our growth setup, there is a slight lateral temperature gradient from the centre of the wafer (highest temperature) to the edge of the wafer (lowest temperature). Thus, for the highest temperature (panel a) we observe a rather regular step bunching with step heights of about eight monolayers (MLs) and terrace widths on the order of 70–80 nm. In the area midway between centre and edge (medium temperature, panel b) we see some long range step bunching with step heights on the order of about 30 MLs and large terrace widths of about 300 nm. Finally, in the area near to the wafer edge (coldest region, panel c) the surface is characterized by double ML steps, thus indicating that step bunching was strongly suppressed in this area. Depending on the intended application, step bunching might or might not be favorable and, correspondingly, a tight control of growth temperature can play an important role.

Although there is a good understanding of the epitaxial growth processes on GaAs(111)B substrates, the optical quality of QWs are still inferior compared to QWs grown on the GaAs(001) surface. This is illustrated in Fig. 6, where...
are accompanied by a piezoelectric potential 

we compare our best PL spectra obtained from GaAs(001) and GaAs(111)B QWs. Both spectra were recorded at a temperature of \( T = 5 \) K under comparable conditions. Whereas the PL signal of the 20 nm thick GaAs(001) QW (black, full line) appears as a very narrow line with a FWHM of only 0.16 meV, the PL line of the 25 nm thick GaAs(111)B sample (red, broken line) is much broader and reveals a FWHM of 2.1 meV. The spectrum of the (111)B SQW sample seems to consist of two lines and might indicate the appearance of neutral and charged excitons. However, this assumption has to be clarified in further experiments.

4 Acoustic transport of carriers and spins

This section reviews recent results on carrier and spin transport by SAWs in III–V semiconductor nanostructures. SAWs are elastic vibrations propagating along a surface. In a piezoelectric material (such as the III–V compounds), these waves are accompanied by a piezoelectric potential \( \Phi_{\text{SAW}} \), which enables their electric generation by interdigital transducers (IDTs) placed on the surface of the material (cf. Fig. 7). As illustrated in the inset of the figure, \( \Phi_{\text{SAW}} \) creates a moving modulation of the conduction and valence band edges, which can capture electrons and holes and transport them with the acoustic velocity.

The initial investigations of the acoustic carrier transport in semiconductors date back to the seventies, when acoustically induced electron transport [79,80], as well as the role of photogenerated carriers in the SAW attenuation was established [81]. By reducing the lateral dimensions of the potentials, the acoustic transport of single-electrons [82–84] has been demonstrated. The type II potential modulation can trap and transport photoexcited electrons and holes over hundreds of \( \mu \text{m} \) [85]. This approach has also been used to create acoustically pumped single-photon sources [86].

The mobile character of the piezoelectric potential is especially suitable for the transport and manipulation of photoexcited spins using the scheme illustrated in Fig. 7. Here, optically oriented electron spins are captured by \( \Phi_{\text{SAW}} \) and transported along the SAW propagation direction. Magnetic fields or electric gates based on the SO-interaction can then be used to control the spin orientation during transport. The spin state can be probed during transport via PL spectrometry [87,88] or Kerr reflectometry [89,90]. After transport, the electrons and holes can be forced to recombine, leading to the emission of circularly polarized light. The acoustic potentials have the following favorable properties:

(i) the spatial separation of electrons and holes by the type II potential increases the recombination lifetime and suppresses spin dephasing via the electron–hole exchange interaction (the excitonic or BAP mechanism) [5,87,88];

(ii) motional narrowing effects induced by mesoscopic (\( \mu \text{m-sized} \)) confinement potentials reduce spin dephasing within an electron spin ensemble. Very long spin lifetimes (> 25 ns) and coherent transport lengths (> 100 \( \mu \text{m} \)) have been measured during acoustic transport by mobile potential dots (dynamic quantum dots, DQDs) produced by acoustic fields in GaAs (001) QWs [88,91];

(iii) the spins are transported with a well-defined average momentum \( h k_{\text{SAW}} = m^* v_{\text{SAW}} \) along the SAW propagation direction, where \( m^* \) is the electron effective mass and \( v_{\text{SAW}} \) the acoustic propagation velocity. This is specially interesting for studies of SO-effects in view of the dependence of \( B_{\text{SO}} \) on electron momentum. The controlled precession of spins during acoustic transport has been demonstrated in Ref. [88]. The dependence of the precession frequency on QW thickness has been used for the direct determination of the SO-splitting constant in GaAs QWs (see Section 2) [88,91,92]. Finally, Sanada et al. [93] have recently demonstrated the full control of the spin vector via the SO-interaction by tailoring the shape of the acoustic transport channel;
(iv) the dynamic strain field of a SAW can also be used to induce a SO-field for spin control [90, 94, 95]. We briefly consider here the effects of the strain field of a Rayleigh SAW. For a Rayleigh SAW propagating along $x$, the strain has three non-vanishing components with amplitudes $u_{xx}$, $u_{yx}$, $u_{zx}$, which vary in time and space with the acoustic frequency ($f_{\text{SAW}}$) and wavelength ($\lambda_{\text{SAW}}$), respectively. If the carriers are transported close to the minima of the electronic piezoelectric energy, they experience a constant strain field during transport. The last row of Table 1 summarizes the $k$-dependence of the SAW strain contribution $\hbar\Omega_S$ for QWs with different orientations. For a SAW along $x$ (i.e., for $k_x = 0$), $(\hbar\Omega_S)$ is proportional to the Rashba contribution and can, therefore, be compensated by an applied electric field. This becomes particularly interesting for the acoustic transport of long-living $z$-oriented spins in (110) QWs, where precession due to the strain field can be avoided by the application of an electric field.

Most of the previously mentioned investigations have been carried out in (Al,Ga)As(001) QW structures. In the following sections, we review recent results on spin transport in (110) QWs as well as future perspectives for the spin transport and manipulation in (111) structures.

5 Acoustic spin transport in GaAs(110) QWs

This section reviews results on the acoustic transport of electron spins in GaAs(110) structures. We first analyze spin transport by SAWs along a single QW at low temperature and compare the transport dynamics to the predictions of Section 2. The long lifetimes of spins in these QWs make them good candidates for spintronic devices also at higher temperatures, where spin dephasing is more pronounced. In the second part of this section, we address acoustic transport at higher temperatures in piezoelectrically defined channels created in GaAs(110) microcavity structures. In these channels, the electron propagation is confined within mesoscopic 1D channels, which, according to the discussion in Section 2.1, provides further suppression of the spin dephasing mechanisms. This, together with the enhanced optical generation and detection of electron spin polarization in QWs embedded in microcavities, allows for spin transport over long distances at temperatures exceeding $T = 100$ K.

5.1 Low temperature spin transport

The studies were carried out on a 20 nm-thick undoped GaAs(110) QW with $\text{Al}_{0.15}\text{Ga}_{0.85}\text{As}$ barriers located 400 nm below the surface. In order to enhance the SAW piezoelectric field, the samples were coated with a 424 nm thick piezoelectric ZnO film. SAWs propagating along the $x$-direction [001] surface direction were generated by IDTs for an acoustic wavelength $\lambda_{\text{SAW}} = 5.6$ $\mu$m fabricated by optical lithography [40].

The experiments for the optical detection of spin transport were carried at 20 K in a cold finger cryostat with an optical window and coaxial connections for the application of radio-frequency (rf) signals to the IDTs. An external coil applies in-plane magnetic fields, $B_{ext}$, of up to 150 mT. The acoustic transport of spins was studied in this sample by spatial- and time-resolved PL as well as magneto-optic Kerr reflectometry [96, 97]. This last technique has been successfully applied for spin transport studies in GaAs QWs involving a single acoustic wave [89], as well as moving potential dots generated by the interference of two perpendicular SAW waves [90]. Here, a circularly polarized pump laser pulse focused on the SAW path (15 $\mu$m diameter spot) optically generates out-of-plane polarized spins. They are then probed during transport by measuring the rotation of the polarization angle, $\delta\theta$, of a weaker, linearly polarized probe pulse reflected on the sample (spot of typically 10 $\mu$m diameter). $\delta\theta$ is proportional to the projection of the average spin vector along the direction perpendicular to the QW plane (the $z$ direction). The probe pulse can be displaced by $\Delta x$ along the SAW path and time-delayed, $\Delta t$, with respect to the pump pulse. Both pump and probe energies are tuned to the electron heavy-hole exciton energy on the QW ($\lambda_{hh} \simeq 813.1$ nm). Further details about the experimental setup can be found at Ref. [89].

The SAWs were generated by applying a radio-frequency signal of $f_{\text{SAW}} = 524$ MHz with a nominal power $P_r = 20$ dBm to the IDT. Although the cryostat nominal temperature was 20 K, additional heating induced by the rf power increases it to $\approx 53$ K, as estimated from the energetic shifts of the electron heavy-hole emission line.

The two-dimensional plots of Fig. 8 compare the spatial-temporal evolution of the out-of-plane spin polarization generated by the pump in the absence and presence of a SAW (left and right panels) under in-plane magnetic fields.
dependence. Fits to Eq. (7) show that the dependence of the spin polarization is mapped into a spatial confinement. By fitting the data along the line $\Delta t = 0$ (dashed line) in Fig. 8a to Eq. (7), we obtained $T_2^∗$ = 1.15 ns for non-precessing spins. As the direction of the optically excited spins is parallel to the effective magnetic field associated with the SO-interaction, DP spin dephasing is not expected in this case. The spin lifetime is mainly limited by the BAP recombination time [23]. By fitting the data along the line $\Delta t = 0$ (dashed line) in Fig. 8a to Eq. (7), we obtained $T_2^∗$ = 1.15 ns for non-precessing spins. As the direction of the optically excited spins is parallel to the effective magnetic field associated with the SO-interaction, DP spin dephasing is not expected in this case. The spin lifetime is mainly limited by the BAP recombination time [23]. By fitting the data along the line $\Delta t = 0$ (dashed line) in Fig. 8a to Eq. (7), we obtained $T_2^∗$ = 1.15 ns for non-precessing spins. As the direction of the optically excited spins is parallel to the effective magnetic field associated with the SO-interaction, DP spin dephasing is not expected in this case. The spin lifetime is mainly limited by the BAP recombination time [23].

In our undoped samples, $T_2^∗$ is related both to the intrinsic spin lifetime $\tau_s$ and to the carrier recombination time, $\tau_r$, by the expression

$$\frac{1}{T_2^∗} = \frac{1}{\tau_s} + \frac{1}{\tau_r}. \tag{9}$$

We have estimated $\tau_r$, from the time dependence of the probe reflectance on the sample surface, which yields 1.84 ns in the absence of a SAW. The situation changes dramatically when SAWs are applied to the QW. In this case, the photo-excited electrons and holes are stored and transported at different phases of the SAW field (cf. upper inset of Fig. 7), thus reducing the overlap of their wave functions and increasing $\tau_s$ to values above 40 ns. The electrons maintain the spin polarization over transport distances of more than 15 $\mu$m, as illustrated in Fig. 8c and d. The oscillations in $\delta\theta_k$ along the dot-dashed line in panel (d) are attributed to the precession around $B_{ext}$ of the photoexcited electron spins moving with $v_{SAW}$. Due to the constant transport velocity, the time-dependence of the spin polarization is mapped into a spatial dependence. Fits to Eq. (7) show that $T_2^∗$ is enhanced to 3.78 and 2.15 ns for $B_{ext} = 0$ and 65 mT, respectively. The longer recombination lifetime, together with the reduction of BAP lifetimes for the present experimental conditions.

5.2 High temperature transport in GaAs (110) microcavity structures

The efficiency of the optical processes for the generation and detection of spins, as well as the spin lifetimes, normally reduce at high temperatures. The reduction in optical efficiency can be partially compensated by embedding the QWs within optical microcavities. Spin scattering can be minimized by combining the special properties of (110) QWs with lateral confinement.

We argued in Section 2.1 that electron transport along 1D channels fixes the rotation axis of the SO-field. If it is further assumed that these fields depend only linearly on $k$, spin dephasing by means of the DP mechanism becomes totally suppressed [28, 99, 100]. The enhancement of the spin lifetimes of electrons in etched (In,Ga)As channels [on GaAs(001)] with decreasing channel widths has been experimentally observed in experiments at $T = 5$ K by Holleitner et al. [29]. Here, we use a similar technique to transport spins in (110) QWs embedded in microcavities at temperatures above 100 K.

Instead of defining the channels by chemical etching, which might induce spin dephasing by electron scattering at the rough sidewalls (i.e., via the Elliot-Yafet spin dephasing mechanism [22]), our approach employs piezoelectrically defined acoustic transport channels. The experimental setup is depicted in Fig. 9a: the transport channel is defined by a thin metal layer containing a narrow slit (3–5 $\mu$m-wide) oriented along the [001] surface direction of a (110) QW structure grown on an (Al,Ga)As Bragg mirror stack (Fig. 9b). The whole structure is then coated by a piezoelectric ZnO$_x$/SiO$_{2-x}$ layer stack, which simultaneously acts as the upper Bragg mirror of an optical microcavity. A narrow (approx. 30 $\mu$m wide) SAW beam with wavelength $\lambda_{SAW} = 5.6$ $\mu$m (corresponding to a frequency $f_{SAW} = 536$ MHz) is then launched along the slit direction by a focusing IDT. The metal layer screens the piezoelectric potential generated by the SAW everywhere in the QW plane except underneath the slit. The latter then defines a narrow channel for the ambipolar acoustic transport of electrons and holes. The transport is blocked at the end of the slit, where the quenching of $\Phi_{SAW}$ by the metal forces the recombination of the carriers.

The lower Bragg mirror and the cavity layers (including the 20 nm-thick GaAs QW) in Fig. 9b were grown by MBE. The 12 nm thick Ti layer was evaporated and photolithographically patterned to form slits with nominal widths ($L_w$) of 3 and 4 $\mu$m. The upper ZnO/SiO$_2$ Bragg mirror was rf-sputtered on top of the patterned Ti layer. The vertical distance between metallic layer and QW is only 110 nm,
Figure 9 (a) Hybrid microcavity with an acoustic transport channel formed by the structured semitransparent Ti layer. The dashed line illustrates the lateral profile of the piezoelectric energy $-e\Phi_{\text{SAW}}$. (b) Detailed layer structure of the hybrid microcavity. The optical layer thicknesses are given in units of the cavity resonance wavelength $\lambda = 820$ nm at $T \approx 100$ K. Thus enhancing the intensity of the stray fields created by the piezoelectric layers.

The composition and thicknesses of the layers were optimized to increase the SAW field close to the QW structure. SAWs result from the confinement of the acoustic fields due to the lower acoustic velocity near the surface. This waveguiding effect reduces if a GaAs substrate is overgrown with an Al$_x$Ga$_{1-x}$As layer, since the average acoustic velocity near the surface increases with the Al content $x$. The latter has been minimized in Fig. 9b by reducing the average Al content ($\leq 56\%$) of the Bragg mirror layers and by the insertion of 3/4 $\lambda$ layers with low $x$ in the lower Bragg mirror. For the same reason, the upper Bragg mirror consists of only two mirror layer pairs on top of a $\lambda/2$ ZnO layer, leading to a cavity quality factor $(Q)$ of only about 100. A significantly enhanced QW PL has nevertheless been observed when the temperature dependent cavity resonance wavelength matches the QW heavy hole exciton wavelength at $T = 105$ K (not shown here).

Figure 10 shows a calculated snapshot of the moving profile of the piezoelectric energy $-e\Phi_{\text{SAW}}$ at the depth of the QW in the hybrid cavity structure depicted in Fig. 9. Calculations were carried out for a nominal channel width ($L_W$) of 3 $\mu$m, an acoustic wavelength $\lambda_{\text{SAW}} = 5.6 \mu$m and an acoustic linear power density $P_{\text{SAW}} = 200$ W m$^{-1}$. The spatially separated electrons and holes are symbolized by red and blue dots, respectively. The orange lines mark the borders of the semitransparent Ti layer.

Figure 11 (a) Schematic illustration of the scanning method used to monitor the acoustic transport of electron spins. (b) PL intensity images taken at $T = 105$ K corresponding to acoustic transport distances of 96, 58, and 26 $\mu$m (top to bottom).
applied acoustic power (acoustically in a GaAs(110) QW at T width. The carriers are then transported by the SAW toward the end of the channel, where the screening of $\Phi_{\text{SAW}}$ terminates the transport and provokes radiative recombination. The spin polarization ($\rho_s$) is determined from the PL intensities with left ($I_-$) and right ($I_+$) circular polarization emitted at the end of the slit according to

$$\rho_s = \frac{I_- - I_0}{I_+ + I_0}. \quad (10)$$

The two-dimensional intensity images of the total PL ($= I_+ + I_-$) for different $\Delta x$ are depicted in Fig. 11b. At the generation point, the high concentration of photo-generated charge carriers partially screens $\Phi_{\text{SAW}}$, thus reducing the transport efficiency. As a result, some of the carriers recombine radiatively giving rise to a weak PL signal. Away from the generation point, the electrons and holes are efficiently captured and transported by the SAW field, thereby making the recombination (e.g., via trapping sites) negligible. In contrast, at the end of the channel, where $\Phi_{\text{SAW}}$ is screened by the Ti layer, considerable PL intensities are recorded and the electron spin polarization can be analyzed. By scanning $\Delta x$ in small steps (approx. 2 $\mu$m) one can determine the electron spin polarization as a function of the transport length [$\rho_s(\Delta x)$]. Taking into account the well-defined SAW velocity $v_{\text{SAW}}$, the electron spin lifetime $\tau$ can be determined from the decay $\rho_s(t) = \rho_s(t = 0) \cdot \exp(-\Delta x/v_{\text{SAW}}\tau_s)$. The propagation delay $\tau$ and distance $\Delta x$ are related to each other by $\tau = \Delta x/v_{\text{SAW}}$.

Figure 12 displays $\tau_s$ measured in samples with nominal channel widths of 3 and 4 $\mu$m under varying acoustic power levels. Considering the large range of applied acoustic power values, $\tau_s$ can be regarded as power independent. However, electron spin lifetimes measured in the narrower channel ($\tau_s \approx 10–11$ ns) are clearly larger than the ones deduced for the wider channel ($\tau_s \approx 7–8$ ns). This enhancement of electron spin lifetimes is expected as the channel widths of the investigated samples are comparable to the SO precession length, which is of approx. 4 $\mu$m for a 20 nm thick GaAs QW.

Figure 13 displays the out-of-plane spin lifetimes ($\tau_s$) of electrons transported acoustically in a GaAs(110) QW at $T = 105$ K as a function of the applied acoustic power ($P_{\text{SAW}}$). $L_W$ indicates the nominal channel width.

### 6 Suppression of the spin–orbit interaction in GaAs(111)B quantum wells

The special characteristics of the SO-fields in GaAs(111) QWs makes them very good candidates for spintronics. As mentioned in Section 2, an electric field applied across a (111) QW can suppress spin dephasing mechanisms associated with the SO-interaction. The studies of the spin dynamics in GaAs(111) QWs as a function of electric and magnetic fields, as well as temperature presented in this section provide experimental evidence for the electric suppression of spin dephasing.

#### 6.1 Electric control of the spin lifetime

The experiments were carried out in QWs embedded in the intrinsic region of a n–i–p structure (cf. inset of Fig. 13). A bias voltage ($V_b$) applied between the top n-doped layer and the p-doped substrate generates the vertical electric field ($E_z$) necessary for SO-compensation. Two kinds of samples were studied: a GaAs single quantum well (SQW) with a thickness of 25 nm and a multiple quantum well (MQW) composed of 20 QWs, each 25 nm-thick, separated by (Al,Ga)As barriers. The position of the QWs between the n-doped layer and the p-doped substrate is the same in both samples. To confine the applied electric field along the z-direction, the n-doped layer and part of the top (Al,Ga)As spacer at the intrinsic region were chemically etched into mesa structures with a diameter of 300 $\mu$m.

We determined the field applied to the QWs by measuring the PL under different voltage biases (cf. Fig. 13 for the SQW sample, note that forward and reverse biases correspond to $V_b > 0$ and $V_b < 0$, respectively). The quantum confined Stark effect (QCSE) induced by $V_b$ modifies the energy, line width, and intensity of the PL line. By comparing the energy shift induced by $V_b$ under reverse bias with numer-
ical simulations of the field distribution, we determined the relationship between $V_b$ and $E_z$ indicated in the upper and lower horizontal axis of the figure.

The n–i–p structures were designed to reach SO-compensation under a reverse bias given by Eq. (2). For low reverse bias ($V_b \geq -0.9$ V), the PL intensity is relatively high (cf. Fig. 13) and the spin polarization $\rho_s$ can be determined from PL circular polarization using Eq. (10). The dependence of $\rho_s$ on reverse bias obtained in the MQW sample using this approach is illustrated by the profiles in Fig. 14a, while the corresponding spin lifetimes are plotted as black dots in Fig. 14c. The spin lifetime increases with the amplitude of the reverse bias from barely 1 ns for $V_b = 0.3$ to up to 60 ns for $V_b = -0.9$ V. Control experiments carried out in a p–i–n structure on GaAs(111)B, where the reverse electric field has the opposite orientation relative to the [111] axis, have shown the opposite behavior [53, 57]. Since in both cases the bias leads to the spatial separation of the carriers along the growth direction, the bias dependence of $\rho_s$ cannot be accounted for by the BAP mechanism.

For reverse voltages $V_b < -1$ V, the overlap between the electron and hole wave functions in the QW as well as the PL reduce significantly, thereby hindering the detection of the spin dynamics from the PL polarization [57]. We overcome this limitation by carrying out experiments under pulsed reverse bias [54]. Here, laser and bias pulses with the same repetition frequency are synchronized with each other so that the laser pulse hits the sample at a time instant $t = 0$ shortly after the application of a reverse bias pulse of amplitude $V_b$. The photoexcited carriers are then driven toward opposite interfaces of the QW, where they remain stored until the pulse is removed. The stored carriers quickly recombine when the bias pulse is removed giving rise to a short PL pulse.

Figure 15 shows time-resolved PL traces of the MQW sample submitted to $t_p = 40$ ns long bias pulses with a repetition period of 80 ns. For $V_b < -1$ V, the QCSE prevents the recombination of the photoexcited carriers during the bias pulse. Electrons and holes remain then stored in the QWs during a time $t_p$. Information about the stored carrier density and spin polarization is extracted at the end of the bias pulse, when a small forward bias ($V_b = 1$ V) is applied to induce carrier recombination. As shown in Fig. 15, the amplitude of the PL pulses after the bias pulse increases for $V_b < -1$ V, since for these biases the carrier lifetime exceeds the pulse width. The PL rise time is determined by the falling time of the bias pulses of 2 ns. The time-integrated PL after bias pulses with $V_b < -2$ V corresponds to 60% of the PL emitted right after the laser pulse under a bias of 0 V (dark blue curve in Fig. 15). This means that the QWSs can efficiently store a high density of carriers over long times. The reduction of the retrieved PL intensities for pulse voltages $V_b < -3.5$ V is attributed to field-induced carrier extraction to the contacts during the storage time, as this bias range also coincides with the onset of the breakdown current of the structures.

The spin dynamics for large reverse biases ($V_b < -1$ V) can be obtained from the circular polarization of the PL pulses emitted at the end of the bias pulses. Figure 14b shows $I_{\uparrow}(t)$ and $I_{\downarrow}(t)$ for bias pulses of $V_b = -2.1$ V and different lengths. The remarkable difference in intensity of $I_{\uparrow}$ and $I_{\downarrow}$ after a time delay of more than 20 ns attests to the conservation of the spin polarization during the long charge storage.

Figure 14 (a) Time-resolved spin polarization during reverse bias of different amplitudes $V_b$. (b) $I_{\uparrow}(t)$ and $I_{\downarrow}(t)$ traces for bias pulses with $V_b = -2.1$ V and durations of 20 ns (black), 30 ns (red), 40 ns (green), 50 ns (blue), and 60 ns (cyan). (c) Spin lifetimes as a function of $V_b$ obtained from the measurements on panels (a) [black dots] and (b) [open squares].
Figure 15 Time-resolved PL of the MQW sample under $t_p = 40$ ns bias pulses with a repetition period of 80 ns. The value of $V_b$ during the pulse goes from a forward bias of 0.6 to a $-3.6$ V reverse bias. Between pulses, $V_b = 1$ V, so that recombination of the carriers stored during the pulse is allowed.

We have determined the out-of-plane spin lifetimes $\tau_z$ for each $V_b$ by measuring $\rho_s$ at the end of voltage pulses of different durations. The results for the MQW sample, indicated by the open squares in Fig. 14c, cover the range of large reverse bias and are thus complementary to those obtained for lower fields from Fig. 14a [black dots]. The spin lifetime increases with reverse bias until it reaches a maximum of about 100 ns for a bias $V_b \approx 2.1$ V, which corresponds to a vertical electric field of about $E_z = 15$ kV cm$^{-1}$ (the compensation field). Beyond this value the spin lifetime decays. The observation of a maximum in Fig. 14 unambiguously establishes the BIA/SIA compensation as the mechanism for bias-induced spin lifetime enhancement in (111) QWs. Furthermore, the peak spin lifetimes exceeding 100 ns are among the longest values reported for undoped GaAs structures. Finally, it is important to note that the Rashba contribution in the n–i–p structures can be electrically increased to values up to approximately 1.5–2 times the Dresselhaus contribution. This range is limited by the onset of field-induced carrier extraction from the QW for $V_b < -3.6$ V.

6.2 Spin lifetime under an external magnetic field The reverse field is also expected to increase the lifetime $\tau_R$ of spins precessing around an externally applied in-plane magnetic field $B_{\text{ext}} \parallel y$. In order to avoid perturbations related to the non-uniformity of the vertical field distribution across the QWs in MQW samples, the investigations under $B_{\text{ext}}$ were performed on the SQW sample. Figure 16a shows $I_{\parallel}(t)$ and $I_{\perp}(t)$ under $V_b = -1.2$ V and $B_{\text{ext}} = 60$ mT for this sample. Quantum beats accompanying the PL intensity decay are clearly observed, which are attributed to the Larmor precession of the electron spins around the external magnetic field [101]. Figure 16b shows the spin polarization, $\rho_s$, obtained from Eq. (10). Due to the fact that the temporal width of the laser pulse (600 ps) is comparable to the spin precession period (3 ns), the initial spin polarization is smaller than the theoretically expected value of approx. 0.25. The $g$-factor obtained from a fitting to an exponentially decaying cosine function (pink broken line) is $|g| \approx 0.42$, in good agreement with the expected value for such a wide QW. As the bias voltage increases, $|g|$ shifts toward smaller values at a rate $\approx 0.013$ V$^{-1}$.

Figure 16 (a) $I_{\parallel}(t)$ and $I_{\perp}(t)$ traces measured at $V_b = -1.2$ V and an in-plane $B_{\text{ext}} = 60$ mT. (b) Time-dependent spin polarization extracted from traces in (a). The pink broken line shows the result of the fitting. (c) Spin lifetime as a function of $E_z$ for different in-plane magnetic fields. The solid lines are guide to the eye, while the broken lines correspond to the values extracted from the model described in the text. The inset shows the spin lifetime as a function of the in-plane magnetic field for a fixed $V_b = -1$ V. All measurements are performed at $T = 39$ K.
Figure 16c compares the electric field dependence of the spin lifetime for out-of-plane (black squares) and precessing (red circles and green triangles) spins in the SQW. The electric field dependence of $\tau_z$ and $\tau_{x,y}$ can be divided into two regions: for electric fields away from $E_c$, $\tau_{x,y}$ is longer than $\tau_z$. This behavior is due to the fact that, at low temperatures (that is, for small $k$-vectors), $\Omega_{SO,c}(k_z) \ll \Omega_{SO,x}(k_z) \approx \Omega_{SO,y}(k_z)$ (cf. Section 2). Under this condition, spins along one of the in-plane directions (e.g., $x$) only feel the dephasing SO-field along the orthogonal in-plane direction (i.e., $\Omega_{SO,y}$). Spins along $z$, in contrast, will see two dephasing fields (associated with $\Omega_{SO,x}$ and $\Omega_{SO,y}$) and decay faster. Therefore, $\tau_z \approx 2 \tau_{x,y}$, and, according to Eq. (6), $\tau_R \approx 4 \tau_z/3$.

The situation reverses close to $E_c$: here, the in-plane linear Dresselhaus contribution can be compensated by the Rashba term (cf. Table 1), but the $z$-component $\Omega_{SO,z}$ (cf. Eqs. 3 and 4) not, thus leading to $\tau_R < \tau_z$. A reduced lifetime for precessing spins (relative to $z$-oriented ones) has been observed over a wide range of magnetic fields (cf. inset of Fig. 16c, with data measured for $V_s = -1$ V). Finally, at the intersection of the two regions all three components of $\Omega_{SO}$ are equal, and the spin lifetime becomes isotropic.

The measured spin lifetimes $\tau_z$ and $\tau_{x,y}$ can be compared to the expected values from the model explained in Section 2. According to the model, the ratio between the Dresselhaus and Rashba constants, $\gamma/r_{d1}$, and the momentum scattering time $\tau_{d1}^*$ determine the compensation field $E_c$ (cf. Eq. 2) and the spin lifetime at compensation (cf. Eq. 5). These two parameters were varied in Eq. (5) to fit the measured results. The broken lines in Fig. 16c show the result of our calculations for $\gamma/r_{d1} = 2.86$ VÅ, and $\tau_{d1}^* = 5$ ps. The model predicts a much stronger field dependence than the measured one. The reason for these discrepancies will be further discussed below. The qualitative features, however, are very well reproduced, including the reduction of the spin lifetime around $E_c$ for precessing spins as well as opposite behavior away from the compensation field.

6.3 Spin lifetime dependence on temperature

According to the theoretical model, the Rashba term compensates the Dresselhaus one when the following condition is fulfilled:

$$E_c = -\frac{\gamma}{\sqrt{3}r_{d1}} \left[ \left( \frac{\pi}{d_{eff}} \right)^2 - \frac{1}{4} \frac{\hbar^2}{k_z^2} \right].$$  \hspace{1cm} (11)

At very low temperatures, the $k_z^2$ term, which comes from nonlinear terms of the Dresselhaus contribution (cf. Eq. 3), can be neglected. In this case, compensation takes place at a vertical electric field fully determined by $\gamma$ and $r_{d1}$ (cf. Eq. 2). When the temperature increases, however, a larger range of electron $k$-space states above the conduction band minimum are occupied, and $k_z$ in Eq. (11) cannot be neglected any more. As a consequence, SO-compensation does not occur simultaneously for all wave vectors at a fixed external electric field. The inclusion of high order terms in $k$ leads to a temperature dependence of the maximum spin lifetime, as well as of $E_c$.

Figure 17 shows the dependence of $\tau_z$ on $E_z$ and temperature, obtained from the decay of the spin polarization in a SQW after optical excitation with a laser pulse. The solid lines are a guide to the eye. The dashed lines are spin lifetimes calculated at the same temperatures using the model described in the text. The inset shows the value of the compensation point, $E_c$, extracted from the experimental values for each temperature.
which is much smaller than the width of the lifetime peak in Figs. 16 and 17.

An alternative mechanism to account for the discrepancy between calculation and measurements are SO-contributions induced by strain fields (cf. Section 2.3). The biaxial strain due to the difference between the elastic constant of GaAs and (Al,Ga)As (which was estimated to have levels lower than $3 \times 10^{-6}$), is expected to have negligible effects on the spins. In this way, it was shown that a SO-contribution in GaAs(111) QWs with the same symmetry as the Dresselhaus term in Table 1, and thus can be compensated by the applied bias. Another possibility are uniaxial strain fields induced, for instance, by sample mounting.

7 Conclusions and future perspectives The special symmetry of the SO-interaction in III–V semiconductor QWs grown along the non-conventional crystallographic directions [110] and [111] allows the detection of electron spins with long decoherence times. In (110) QWs, the combination of such SO-symmetries with the special properties of SAWs leads to acoustic transport of electron spins with out-of-plane polarization along distances of several micrometers at low temperatures. Acoustic spin transport also reduces BAP spin dephasing via the spatial separation of electrons and holes by the type II modulation potential. Furthermore, by embedding these QWs within optical microcavities compatible with SAW generation, we have demonstrated acoustic spin transport within narrow piezoelectric channels over distances of several tenths of $\mu$m at temperatures above 100 K. Further developments of these structures should enable spin transport also at higher temperatures.

In the case of GaAs(111) QWs, the application of vertical electric fields enables the efficient suppression of the SO-interaction at low temperatures. As a result, spin lifetimes up to 100 ns have been reached. These lifetimes arise from the compensation of terms of the SO-interactions that are linear to $1/\hbar k$. At higher temperatures, higher orders terms in $k$ become important and prevent the compensation effect that leads to such long lifetimes. According to the results obtained from their (110) counterparts, acoustic spin transport within narrow channels appear to be a promising candidate to overcome this limitation: the well-defined velocity and direction of the acoustic wave allows the spatial confinement of spins and their transport along special crystallographic directions, for which the high order terms of the SO-fields in (111) QWs are minimized. Figure 18 shows a device proposal combining SAWs with vertical electric fields. This type of devices was designed based on our experience with (110) QW structures for acoustic transport (see, for instance, Fig. 9). The QW is embedded within a piezoelectric ZnO/SiO$_2$ microcavity for operation at relatively high temperatures. The gate placed along the SAW path allows the application of vertical electric fields for the suppression of the SO-interaction in the QW. In addition, the narrow channel defined by the gate enhances the piezoelectric potential $\Phi_{\text{SAW}}$ underneath it, thereby laterally confining the electrons and holes moving in the SAW. This confinement effect, which is discussed in detail in Ref. [102] and illustrated by the calculated depth profiles for $\Phi_{\text{SAW}}$ in Fig. 18c, arises from the fact that $\Phi_{\text{SAW}}$ near the interface between the ZnO/SiO$_2$ stack and the GaAs structure increases when the surface of the samples is coated with a conductive layer. The metal structure placed underneath the ZnO/SiO$_2$ DBR at the end of the transport path (named “stopper” in Fig. 18) quenches the SAW piezoelectric field at the end of the transport path, thereby forcing the carrier recombination required for the optical retrieval of the spin information. The realization of a device of such characteristics is expected to be one of our main activities during the next years.

Acknowledgements We thank F. Iikawa for discussions, as well as M. Hörnicke, W. Seidel, B. Drescher, S. Rauwerdink, and A. Tahraoui for sample processing. Financial support by the DFG (Priority Program SPP1285) is thankfully acknowledged.

Note from the authors After submission of this manuscript two new articles appeared in the literature regarding spin dynamics in GaAs(111) QWs. The first one deals with the estimation of the Dresselhaus and Rashba coefficients by spin grating experiments [103], while the second one shows a model for the DP mechanism which takes the collision processes explicitly into account and allows $\tau_p^*$ to change with temperature [104].

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