Effect of Dresselhaus spin-orbit coupling on spin dephasing in asymmetric and macroscopically symmetric (110)-grown quantum wells

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We develop the microscopic theory of electron spin dephasing in (110)-grown quantum wells where the electron scattering time is comparable to or exceeds the period of spin precession in the effective magnetic field caused by spin-orbit coupling. Structures with homogeneous and fluctuating Rashba field, which triggers the dephasing of electron spins aligned along the growth direction, are analyzed. We show that the Dresselhaus field, which is always present in zinc-blende-type quantum wells, suppresses the spin dephasing enabling very long spin lifetime of conduction electrons. The dependence of the spin lifetime on the electron mobility is found to be nonmonotonic reaching the minimum in structures where the scattering time is comparable to the period of spin precession in the effective magnetic field.

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I. INTRODUCTION

The spin dynamics of free electrons in zinc-blende-type quantum wells (QWs) is largely determined by the electron subband spin-orbit splitting which can be fruitfully treated as a wavevector-dependent effective magnetic field acting upon the spins of individual electrons \( \hbar \)\[^{1–4} \]. The field occurs due to the lack of space inversion symmetry in a particular QW and can originate from bulk, structure, or interface inversion asymmetry, or their joint action. The field direction and magnitude depend on the QW crystallographic orientation and design as well as external gate voltage \( E \)\[^{5–12} \], which opens a way towards the spin manipulation and spin-based semiconductor devices. Of special interest are QWs grown on the (110) substrate. The effective magnetic field caused by bulk inversion asymmetry (Dresselhaus field) in such structures points along the growth direction and, therefore, does not cause the precession of electron spins oriented along the QW normal \( \hbar \|^1 \). This enables the observation of very long electron spin lifetimes up to hundreds of nanoseconds at low temperatures which are unfeasible for two-dimensional structures of other crystallographic symmetries \( \hbar \)^{13–17}. The electron spin lifetime in (110)-grown QWs at low temperature is likely to be limited by the spin precession in the homogeneous Rashba field present in asymmetric QWs or unavoidable fluctuating Rashba field that emerges due to domain structure formation or inhomogeneous distribution of charged impurities in doped structures \( \hbar \)^{18}. The Rashba field lies in the QW plane and starts up the spin dephasing. The previous theoretical studies of electron spin dephasing in (110)-grown QWs with the homogeneous \( \hbar \)^{19,20} or fluctuating \( \hbar \)^{21–23} Rashba field were limited to the collision-dominated regime of spin dynamics, where the scattering time is much smaller than the spin precession period in the effective magnetic field. Such a consideration is inapplicable for QWs with high electron mobility and/or strong spin-orbit coupling, which are now technological available. In this paper, we develop the microscopic theory of electron spin dephasing for arbitrary ratio between the scattering time and the spin precession period. We show that the interplay of Rashba and Dresselhaus fields leads to dynamic coupling of the in-plane and out-of-plane spin components. The dependence of the spin lifetime on the electron mobility is nonmonotonic and reaches the minimum in QW structures where the scattering time is comparable to the period of spin precession in the effective magnetic field. In high-mobility QWs, the Dresselhaus field efficiently suppresses the spin dephasing allowing one to reach very long spin lifetimes.

II. UNIFORM RASHBA FIELD

First, we consider asymmetric (110)-grown QWs with a spatially homogeneous Rashba field. The linear in the wavevector spin-orbit splitting of the electron subband in such QWs (\( C_x \) point group) is generally described by three linearly independent parameters \( \alpha_1 \), \( \alpha_2 \), and \( \beta \) \( \hbar \)^{24} (see also Refs. \( \hbar \)^{25,26}). The corresponding Hamiltonian has the form

\[
H_{so} = \alpha_1 \sigma_y k_y - \alpha_2 \sigma_y k_x + \beta \sigma_z k_x,
\]

where \( \sigma_j \) (\( j = x, y, z \)) are the Pauli matrices, \( k_j \) are the components of the electron wavevector, \( x \parallel [110] \) and \( y \parallel [001] \) are the in-plane axes, and \( z \parallel [110] \) is the QW normal. The parameters \( \alpha_1 \) and \( \alpha_2 \) are nonzero in QWs with structure inversion asymmetry only, while \( \beta \) requires bulk inversion asymmetry and describes the Dresselhaus splitting in (110)-oriented QWs. The difference between \( \alpha_1 \) and \( \alpha_2 \) is usually small because it is caused by an interference effect between cubic structure of the crystal lattice and QW structure inversion asymmetry \( \hbar \)^{24}. Therefore, we assume in what follows that \( \alpha_1 = \alpha_2 = \alpha \). The Hamiltonian \( \hbar \)\[^{1} \) is similar to the Zeeman term with the effective Larmor frequency

\[
\Omega_k = \left( \Omega_R k_y, -\Omega_R k_x, \Omega_F k_x \right),
\]
In the ballistic regime, i.e., the absence of scattering, the electron spins initially oriented along the QW normal. Precussion of electron spins initially aligned along the QW normal with the frequency $\Omega_k$ leads to the emergence of an average spin component along $y$. (a) Dependence of the Larmor frequency $\Omega_k$ corresponding to the effective magnetic field on the wavevector $k$ at the Fermi circle. (b) The $(x\hat{z})$ plane containing $\Omega_k$ for all the wavevectors $k$. Precussion of electron spins initially aligned along the QW normal $z$ with the frequency $\Omega_k$ leads to the emergence of an average spin component along $y$.

where $\Omega_R = 2\alpha k_F / h$ and $\Omega_D = 2\beta k_F / h$ are the frequencies corresponding to Rashba and maximal Dresselhaus splitting at the Fermi level, and $k_F$ is the Fermi wavevector. Following the experiments [7, 8, 15, 17] we consider the degenerate electron gas and study the spin dynamics at the Fermi level.

Figure 1 shows the distribution of the frequency $\Omega_k$ on the wavevector at the Fermi level. The distribution is strongly anisotropic; the frequency $\Omega_k$ contains both the in-plane and out-of-plane components. However, one can notice that the frequencies $\Omega_k$ corresponding to various wavevectors all lie in a certain plane ($x\hat{z}$), see Fig 1b. The $(x\hat{z})$ plane is obtained from $(xz)$ by the frame rotation around the $x$ axis with the angle $\theta = \arctan(\Omega_R / \Omega_D)$.

A. Time dependence

The fact that the plane of the $\Omega_k$ vectors does not contain the growth direction $z$ nor is perpendicular to $z$ is a feature of low-symmetry QWs and leads to the coupling of the in-plane and out-of-plane spin components. This is illustrated in Fig. 1b which shows the precession of electron spins initially oriented along the QW normal. In the ballistic regime, i.e., the absence of scattering, the time dependence of electron spin at each point $k$ on the Fermi circle is given by

$$s_k(t) = \frac{\Omega_k \cdot s(0)}{\Omega_k^2} + \left[ s(0) - \frac{\Omega_k \cdot s(0)}{\Omega_k^2} \Omega_k \right] \cos(\Omega_k t)$$

$$+ \frac{\Omega_k \times s(0)}{\Omega_k} \sin(\Omega_k t),$$

where $s(0)$ is the initial spin. The time evolution of the components of the total electron spin

$$S(t) = \sum_k s_k(t)$$

is obtained by summation $s_k(t)$ over the wavevectors at the Fermi level, which yields

$$S_j(t) = \sum_j \left< \frac{\Omega_{k,j} \Omega_{k,j'}}{\Omega_k^2} \right> S_j'(0)$$

$$+ \sum_j \left< \frac{\Omega_k^2 \delta_{jj'} - \Omega_{k,j} \Omega_{k,j'}}{\Omega_k^2} \cos(\Omega_k t) \right> S_j'(0),$$

where the angle brackets denote averaging over the direction of $k$. The first term on the right-hand side of Eq. 5 presents the averaged projection of individual electron spins onto the effective field direction and is time-independent. In (110) QWs, this contribution is present even for the spin polarized along the QW normal, which is in contrast to (001)-grown structures where the effective field lies in QW plane and the time-independent term vanishes for $S(0) \parallel z$. The second term on the right-hand side of Eq. 5 describes spin oscillations.

Time dependence of the out-of-plane and in-plane components of the total spin oriented along the QW normal at $t = 0$ is shown by dotted curves in Fig. 2a and 2b. One can see that, in asymmetrical (110)-grown QWs, the $S_z$ and $S_y$ projections are dynamically coupled: the dephasing of spins polarized along the QW normal leads to the emergence of the in-plane spin component. Both $S_z(t)$ and $S_y(t)$ contain oscillations which slowly decay in time.
Analytical expressions for \( S_z(t) \) and \( S_y(t) \) at \( t > 1/\Omega_k \) can be derived by the stationary phase method, which gives

\[
\frac{S_z(t)}{S_z(0)} \approx \left(1 - \frac{\Omega_R}{\Omega_\Sigma}\right) + \frac{2\Omega_R}{\pi\Omega_D^2} t \cos(\Omega_R t + \frac{\pi}{4}) + \frac{2\Omega_R^2}{\pi\Omega_D^3 \Omega_\Sigma^2} \cos(\Omega_\Sigma t - \frac{\pi}{4}),
\]

where \( \Omega_\Sigma = \sqrt{\Omega_R^2 + \Omega_D^2} \). The oscillation frequencies are determined by the smallest (\( \Omega_R \)) and the largest (\( \Omega_\Sigma \)) magnitudes of the effective Larmor frequency at the Fermi circle. At \( t \rightarrow \infty \), the oscillations vanish and the spin polarization reaches a stationary value.

The scattering of electrons by static defects or phonons in real structures leads to the isotropization of spin distribution in \( k \)-space which drastically suppresses the spin oscillations. In the presence of scattering, the spin dynamics can be described by the kinetic equation

\[
\frac{\partial s_k}{\partial t} + s_k \times \Omega_k = g - \frac{s_k - \langle s_k \rangle}{\tau},
\]

where \( g \) is the spin generation rate at the Fermi level, e.g., due to resonant optical pumping with circularly polarized light, and \( \tau \) is the scattering time. We assume that \( g \) is independent of the wavevector direction. The last term on the right-hand side of Eq. (4) presents a simple form of the collision integral. Effects of scattering anisotropy on spin dephasing were studied in Ref. [27].

Solid curves in Figs. 2b and 2d show the time evolution of the out-of-plane and in-plane projections of the total electron spin \( S(t) \) oriented at \( t = 0 \) along the QW normal. The curves are calculated numerically after Eq. (7) for different scattering times \( \tau \). As expected, the decrease in the scattering time results in a suppression of spin oscillations. Surprisingly, the effect of electron scattering on spin lifetime is more complicated. Comparison of different curves in Fig. 2b and 2d shows that the shortest spin lifetime \( T \sim \tau \) is achieved at a certain mobility corresponding to \( \tau \sim 1/\Omega_k \). At both higher, \( \tau > 1/\Omega_k \), and lower, \( \tau < 1/\Omega_k \), mobility the spin lifetime increases as \( T \propto \tau \) and \( T \propto 1/(\Omega_k^2 \tau) \), respectively.

In the collision-dominated regime, \( \Omega_R \tau, \Omega_D \tau \ll 1 \), one can analytically solve Eq. (11) and obtain the master equation for the total electron spin [11]

\[
\frac{dS}{dt} = G - \Gamma_{cd} S,
\]

where \( G = \sum_k g_k \) and \( \Gamma_{cd} \) is the D’yakonov-Perel’ spin-relaxation-rate tensor, \( (\Gamma_{cd})_{j,j'} = \tau \Omega_R^2 \delta_{j,j'} - \Omega_k \langle \Omega_{k,j} \Omega_{k,j'} \rangle \). For (110)-grown QWs, the tensor \( \Gamma_{cd} \) has the form [20]

\[
\Gamma_{cd} = \frac{\tau}{2} \begin{bmatrix} \Omega_R^2 + \Omega_D^2 & 0 & 0 \\ 0 & \Omega_R^2 + \Omega_D^2 & \Omega_R \Omega_D \\ 0 & \Omega_R \Omega_D & 2\Omega_R^2 \end{bmatrix}.
\]

Solution of Eq. (3) with \( G = 0 \) and the initial condition \( S(0) \) has the biexponential form

\[
S_z(t) = \frac{\Omega_R^2 \exp(-\gamma y t) + \Omega_D^2 \exp(-\gamma z t)}{\Omega_R^2 + \Omega_D^2} S_z(0), \quad S_y(t) = \frac{\Omega_R \Omega_D \exp(-\gamma y t) - \exp(-\gamma z t)}{\Omega_R^2 + \Omega_D^2} S_z(0),
\]

where \( \gamma y = (2\Omega_R^2 + \Omega_D^2) \tau/2 \) and \( \gamma z = \Omega_R^2 \tau/2 \) are eigen values of the \( \Gamma \) tensor. The functions (11) are plotted in Figs. 2a and 2b by dashed curves and demonstrate the perfect agreement with the results of numerical calculation for the collision-dominated regime.

B. Continuous-wave excitation

Besides experiments with time resolution, spin dephasing is also widely studied at continuous-wave (cw) pumping. In the regime linear in the pump intensity, the general relation between the steady-state spin \( S \) and the spin generation rate \( G \) has the form

\[
S = TG,
\]

where \( T \) can be referred to as the spin lifetime tensor. To calculate \( T \) we solve Eq. (11) for the stationary case following the procedure described in Ref. [22]. We express the spin distribution function \( s_k \) via its average value

\[
s_k = \langle s_k \rangle + g \tau, \quad \text{sum up Eq. (12) over the wavevector and obtain the linear equation set}
\]

\[
\sum_{j'j} \left( \frac{\Omega_R^2 \tau \delta_{j,j'} - \Omega_k \langle \Omega_{k,j} \Omega_{k,j'} \rangle}{1 + \Omega_k^2 \tau^2} \right) (S_{j'} + G_{j'} \tau) = G_j.
\]

Solution of the equation set (13) for the effective Larmor frequency (3) yields

\[
T = \frac{\tau}{\sqrt{1 + \Omega_R^2 \tau^2}} \sqrt{1 + \Omega_\Sigma^2 \tau^2} - 1
\]

\[
\begin{bmatrix}
1 + \Omega_R^2 (1 + \Omega_\Sigma^2 \tau^2) / \Omega_R^2 & 0 & 0 \\
0 & 2 + \Omega_R^2 \tau^2 - \Omega_D (1 + \Omega_R^2 \tau^2) / \Omega_R & 0 \\
0 & -\Omega_D (1 + \Omega_R^2 \tau^2) / \Omega_R & \Omega_R (1 + \Omega_R^2 \tau^2) / \Omega_R^2
\end{bmatrix}
\]

Equation (13) is one of the principal results of the paper. It enables one to calculate the steady-state spin polarization for arbitrary ratio between the scattering time \( \tau \) and the precession periods \( 1/\Omega_R, 1/\Omega_D \) and shows that the Dresselhaus field considerably affects the electron spin dephasing in (110)-grown QWs.
The dependence of the steady-state projections $S_z$ and $S_y$ of the steady-state spin, caused by excitation along the QW normal $z$, on the Dresselhaus field. The curves are calculated after Eqs. (11) and (14) for different $\Omega_D\tau$. Despite the fact that the Dresselhaus field efficiently suppresses the spin dephasing increasing $S_z$ as well as leads to the emergence of $S_y$. The $S_z$ projection (Fig. 3a) is determined by the $T_{zz}$ component of the spin lifetime tensor. The latter equals to $1/(\Omega_D^2\tau)$ at $\Omega_D = 0$ and increases with the Dresselhaus field as $(\Omega_D^2/\Omega_R^2)\sqrt{1 + \Omega_D^2/\Omega_R^2}$ at $\Omega_D\tau \gg 1$. Such an increase in the spin lifetime is caused by stabilizing action of the Dresselhaus field on the normal spin component. The in-plane spin projection (Fig. 3b) is determined by the off-diagonal component $T_{yz}$ which nonmonotonously depends on the Dresselhaus field. This component linearly increases with the Dresselhaus field as $-\Omega_D(1 + \Omega_D^2\tau^2)/(\Omega_R^2\tau)$ at small $\Omega_D$, reaches a peak value of $-\tau(1 + \Omega_D^2\tau^2)^{3/2}/[\Omega_R^2\tau^2(2 + \Omega_D^2\tau^2)]$ at $\Omega_D = \Omega_R\sqrt{(1 + \Omega_D^2\tau^2)(2 + \Omega_D^2\tau^2)}$, and then saturates to $-\Omega_D(1 + \Omega_D^2\tau^2)/(\Omega_R)$ at $\Omega_D\tau \gg 1$.

The dependence of the steady-state projections $S_z$ and $S_y$ on the scattering time $\tau$ is shown Fig. 4. In low mobility structures, the spin polarization decreases with the increase in $\tau$, in accordance with the theory of D’yakonov-Perel spin relaxation in the collision-dominated regime. In particular, the spin lifetime tensor components determining $S_z$ and $S_y$ in this regime are given by $T_{zz} \approx (\Omega_R^2 + \Omega_D^2)/[\Omega_R^2(\Omega_R^2 + \Omega_D^2/2)\tau]$ and $T_{yz} \approx -\Omega_D/[\Omega_R^2(\Omega_R^2 + \Omega_D^2/2)\tau]$. At further increase in the scattering time, the spin polarization reaches a minimum and then rises linearly with $\tau$, see Fig. 4. Such a behavior is caused by the transition from the collision-dominated to oscillatory regime of spin dynamics, see Fig. 2. In the oscillatory regime, the $T_{zz}$ and $T_{yz}$ components are given by $\Omega_D^2\tau/[\Omega_R\sqrt{\Omega_D^2 + \Omega_D^2}]$ and $-\Omega_D\tau/[\sqrt{\Omega_D^2 + \Omega_D^2}]$, respectively. We note that the nonmonotonic dependence $S_z(\tau)$ is a feature of low-symmetry QWs and absent, e.g., in (001) QWs where $S_z \propto 1/\tau$ in both collision-dominated and oscillatory regimes.

**III. FLUCTUATING RASHBA FIELD**

Now, we analyze the spin dephasing in macroscopically symmetric QWs where the Rashba field arises due to QW domain structure or inhomogeneous distribution of charged impurities in the doping layers\[18\] and fluctuates in the QW plane around its zero average value. In such systems, the Larmor frequency $\Omega_k$ corresponding to the total effective field is given by the sum of $\Omega_{D,k} = \Omega_D(0,0,k_x/k_F)$ and position-dependent $\Omega_{R,k}(\rho) = \Omega_R(\rho)(k_y/k_F,-k_z/k_F,0)$ with $\rho$ being the in-plane coordinate. The fluctuating Rashba field is usually weak, therefore, we assume that $|\Omega_R(\rho)| \ll |\Omega_D|, 1/\tau$.

During the spin dephasing, electrons travel in the QW plane and their spins are affected by the Rashba field of different strength and sign. The characteristic time of the Rashba field change for an individual electron $\tau_c$
domain to another more frequently than the rates. The in-plane component of the net electron spin \( \gamma \) is small. The out-of-plane component exhibits exponential decay

\[
S_z(t) = S_z(0) \exp(-\gamma z t),
\]

with the dephasing time

\[
\gamma z = \frac{\Omega_R^2(\rho) \tau}{1 + \sqrt{1 + \Omega_D^2 \tau^2}}.
\]

The net electron spin averaged over all domains is given by

\[
\bar{S}_z(t) = \int_{-\infty}^{\infty} S_z(t, \Omega) d\Omega_R,
\]

where \( w(\Omega_R) \) is the distribution function of the Rashba field in the QW plane. Different dephasing rates \( \gamma_z \) in domains lead to a non-exponential decay of the net electron spin. In particular, at large times \( t \gg 1/\gamma_z \), the time dependence of the net spin has the form

\[
\bar{S}_z(t) = w(0) \left[ \frac{\pi(1 + \sqrt{1 + \Omega_D^2 \tau^2})}{\pi t} \right]^{1/2} S_z(0).
\]

The presence of the Dresselhaus spin-orbit coupling leads to a slowdown of spin dephasing in all domains and, hence, to an increase in the net spin polarization.

(iii) Small domains, \( 1/(\Omega_R^2 \tau) \gg \tau_c \gg 1/(\Omega_D^2 \tau), \tau \). In this regime of spin dephasing, electrons move from one domain to another more frequently than the \( \gamma_z \) and \( \gamma_y \) rates. The in-plane component of the net electron spin does not emerge due to the frequent change in the Rashba field sign. The out-of-plane component exponentially decays at the rate

\[
\gamma_{zz} = \frac{\Omega_R^2(\rho) \tau}{1 + \sqrt{1 + \Omega_D^2 \tau^2}}.
\]

where \( \rho_0 \) is the initial electron coordinate, \( v = v_F(\vec{k}/k) \), and \( v_F \) is the Fermi velocity. Equation (22) can be solved by changing the spin coordinate frame to the one rotating along the \( z \) axis with the angular frequency \( \Omega_{D,R} \). The procedure is equivalent to the substitution \( \vec{s}_k = \vec{R}_k \vec{s}_k^\prime \) and \( \Omega_{R,k} = \vec{R}_k \Omega_{R,k}^\prime \vec{R}_k \), \( \vec{R}_k \) being the rotation matrix,

\[
\frac{d\vec{s}_k}{dt} + \vec{s}_k \times [\Omega_{D,k} + \Omega_{R,k}(\rho_0 + vt)] = 0,
\]

This yields \( d\vec{s}_k/dt = \vec{s}_k \times \vec{R}_{R,k} \vec{s}_k \). Taking into account that the Rashba field lies in the QW plane and \( \Omega_R(\rho) \tau_c \ll 1 \) we obtain the dephasing rate of the \( z \) spin component for electrons with the wavevector \( \vec{k} \)

\[
\Gamma_{zz}(k) = \int_{0}^{\infty} \Omega_{R,k}(\rho_0 + vt) \cdot \Omega_{R,k}(\rho_0) dt
\]

where the overline denotes averaging over the initial electron coordinate.

The spatial correlation function of the Rashba field caused by Coulomb impurities randomly distributed in the \( \delta \)-doping layers positioned at the distance \( z_d \) below and above the QW has the form (see Ref. \[18\])

\[
\Omega_R(\rho) \Omega_R(\rho') = \frac{\Omega_R^2(\rho)(2z_d)^3}{[(\rho - \rho')^2 + (2z_d)^2]^{3/2}}.
\]

Combining Eqs. (23) and (24) one obtains

\[
\Gamma_{zz}(k) = \frac{\Omega_R^2(\rho) \Omega_D \tau_c^2(k_x/k_F) K_1[\Omega_D \tau_c(k_x/k_F)]}{[2\tau_c(k_x/k_F)^2]^{3/2}}.
\]
The theoretical results on spin dephasing rate in quantum wells with the fluctuating Rashba field are summarized in Fig. 5. The strong Dresselhaus field increases the electron spin lifetime in all regimes of spin dynamics.

IV. SUMMARY

We have developed the microscopic theory of electron spin dephasing for asymmetric and macroscopically symmetric (110)-grown quantum wells where the Rashba field is homogeneous or fluctuates in the quantum well plane, respectively. Both low-mobility and high-mobility structures are analyzed. It is shown that the regime and rate of spin dephasing are determined by the strength of Rashba and Dresselhaus fields, the latter being always present in zinc-blende-type structures, as well as the electron mobility. The Dresselhaus field points along the growth direction in (110) quantum wells and increases the lifetime of the out-of-plane spin component.

In asymmetric structures, the interplay of the homogeneous Rashba and Dresselhaus fields leads to a coupling of the in-plane and out-of-plane spin components. The spin lifetime of the out-of-plane component nonmonotonically depends on the electron mobility. It linearly decreases with the scattering time \( \tau \) in low-mobility structures, reaches the minimum at \( \tau \) comparable to the period of spin precession in the effective magnetic field, and then linearly increases with \( \tau \) in high-mobility structures.

In quantum wells with the fluctuating Rashba field, the scenario of spin dephasing is determined by the scale of the field fluctuations compared to other relevant scales. We show that, for all regimes of spin dynamics, the strong Dresselhaus field slows down the dephasing rate of the out-of-plane spin component.

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[1] M. I. D’yakonov and V. Yu. Kachorovskii, “Spin relaxation of conduction electrons in noncentrosymmetric semiconductors,” Fiz. Tekh. Poluprovodn. 20, 178 (1986), [Sov. Phys. Semicond. 20, 110 (1986)].
[2] V. N. Gridnev, “Theory of Faraday rotation in quantum wells with large spin splitting,” Pis’ma Zh. Eksp. Teor. Fiz. 74, 417 (2001).
[3] N. S. Averkiev, L. E. Golub, and M. Willander, “Spin relaxation anisotropy in two-dimensional semiconductor systems,” Journal of Physics: Condensed Matter 14, R271 (2002).
[4] M. W. Wu, J. H. Jiang, and M. Q. Weng, “Spin dynamics in semiconductors,” Phys. Rep. 493, 61 (2010).
[5] Y. Ohno, R. Terauchi, T. Adachi, F. Matsukura, and H. Ohno, “Spin Relaxation in GaAs(110) Quantum Wells,” Phys. Rev. Lett. 83, 4196 (1999).
[6] O. Z. Karimov, G. H. John, R. T. Harley, W. H. Lau, M. E. Flatté, M. Henini, and R. Airey, “High Temperature Gate Control of Quantum Well Spin Memory,” Phys. Rev. Lett. 91, 246601 (2003).
[7] V. V. Bel’kov, P. Olbrich, S. A. Tarasenko, D. Schuh, W. Wegscheider, T. Korn, C. Schüller, D. Weiss, and S. D. Ganichev, “Symmetry and Spin Dephasing in (110)-Grown Quantum Wells,” Phys. Rev. Lett. 100, 176806 (2008).
[8] P. Olbrich, J. Allerdings, V. V. Bel’kov, S. A. Tarasenko, D. Schuh, W. Wegscheider, T. Korn, C. Schüller, D. Weiss, and S. D. Ganichev, “Magnetogyrotropic photogalvanic effect and spin dephasing in (110)-grown GaAs/Al_{x}Ga_{1-x}As quantum well structures,” Phys. Rev. B 79, 245329 (2009).
[9] P. S. Eldridge, J. Hübner, S. Oertel, R. T. Harley, M. Henini, and M. Oestreich, “Spin-orbit fields in...
asymmetric (001)-oriented GaAs/Al\textsubscript{x}Ga\textsubscript{1-x}As quantum wells,” Phys. Rev. B 83, 041301 (2011)

[10] D. J. English, P. G. Lagoudakis, R. T. Harley, P. S. Eldridge, J. Hübner, and M. Oestreich, “Strain-induced spin relaxation anisotropy in symmetric (001)-oriented GaAs quantum wells,” Phys. Rev. B 84, 155323 (2011)

[11] A. Balocchi, Q. H. Duong, P. Renucci, B. L. Liu, C. Fontaine, T. Amand, D. Lagarde, and X. Marie, “Full Electrical Control of the Electron Spin Relaxation in GaAs Quantum Wells,” Phys. Rev. Lett. 107, 136604 (2011)

[12] H. Q. Ye, G. Wang, B. L. Liu, Z. W. Shi, W. X. Wang, C. Fontaine, A. Balocchi, T. Amand, D. Lagarde, P. Renucci, and X. Marie, “Growth direction dependence of the electron spin dynamics in 111 GaAs quantum wells,” Applied Physics Letters 101, 032104 (2012)

[13] S. Döhrmann, D. Hägele, J. Rudolph, M. Bichler, D. Schuh, and M. Oestreich, “Anomalous Spin Dephasing in (110) GaAs Quantum Wells: Anisotropy and Inter-subband Effects,” Phys. Rev. Lett. 93, 147405 (2004)

[14] O. D. D. Couto, F. Ikawa, J. Rudolph, R. Hey, and P. V. Santos, “Anisotropic Spin Transport in (110) GaAs Quantum Wells,” Phys. Rev. Lett. 98, 036603 (2007)

[15] G. M. Müller, M. Römer, D. Schuh, W. Wegscheider, J. Hübner, and M. Oestreich, “Spin Noise Spectroscopy in GaAs (110) Quantum Wells: Access to Intrinsic Spin Lifetimes and Equilibrium Electron Dynamics,” Phys. Rev. Lett. 101, 206601 (2008)

[16] R. Völkli, M. Griesbeck, S. A. Tarasenko, D. Schuh, W. Wegscheider, C. Schüller, and T. Korn, “Spin dephasing and photoinduced spin diffusion in a high-mobility two-dimensional electron system embedded in a GaAs-(Al,Ga)As quantum well grown in the [110] direction,” Phys. Rev. B 83, 241306 (2011)

[17] M. Griesbeck, M. M. Glazov, E. Ya. Sherman, D. Schuh, W. Wegscheider, C. Schüller, and T. Korn, “Strongly anisotropic spin relaxation revealed by resonant spin amplification in (110) GaAs quantum wells,” Phys. Rev. B 85, 085313 (2012)

[18] E. Ya. Sherman, “Random spin–orbit coupling and spin relaxation in symmetric quantum wells,” Applied Physics Letters 82, 209 (2003)

[19] X. Cartoixà, D. Z.-Y. Ting, and Y.-C. Chang, “Suppression of the D’yakonov-Perel’ spin-relaxation mechanism for all spin components in [111] zincblende quantum wells,” Phys. Rev. B 71, 045313 (2005)

[20] S. A. Tarasenko, “Spin relaxation of conduction electrons in (110)-grown quantum wells: A microscopic theory,” Phys. Rev. B 80, 165317 (2009)

[21] M. M. Glazov, M. A. Semina, and E. Ya. Sherman, “Spin relaxation in multiple (110) quantum wells,” Phys. Rev. B 81, 115332 (2010)

[22] Y. Zhou and M. W. Wu, “Spin relaxation due to random Rashba spin-orbit coupling in GaAs (110) quantum wells,” EPL (Europhysics Letters) 89, 57001 (2010)

[23] M.M. Glazov, E.Ya. Sherman, and V.K. Dugaev, “Two-dimensional electron gas with spin-orbit coupling disorder,” Physica E 42, 2157 (2010)

[24] M. O. Nestoklon, S. A. Tarasenko, J.-M. Jancu, and P. Voisin, “Spin splitting of electron states in (110) quantum wells: Symmetry analysis and k·p theory versus microscopic calculations,” Phys. Rev. B 85, 205307 (2012)

[25] X. Cartoixà, L.-W. Wang, D.Z.-Y. Ting, and Y.-C. Chang, “Higher-order contributions to Rashba and Dresselhaus effects,” Phys. Rev. B 73, 205341 (2006)

[26] V. A. Shalygin, H. Diehl, Ch. Hoffmann, S. N. Danilov, T. Herrle, S. A. Tarasenko, D. Schuh, Ch. Geri, W. Wegscheider, W. Prettl, and S. D. Ganichev, “Spin photocurrents and circular photon drag effect in (110)-grown quantum well structures,” Pis’ma Zh. Eksp. Teor. Fiz. 84, 666 (2006), [JETP Lett. 84, 570 (2006)].

[27] A. V. Poshakinskiy and S. A. Tarasenko, “Electron spin dephasing in two-dimensional systems with anisotropic scattering,” Phys. Rev. B 84, 155326 (2011)

[28] S. A. Tarasenko, “Scattering induced spin orientation and spin currents in gyrotropic structures,” Pis’ma Zh. Eksp. Teor. Fiz. 84, 233 (2006), [JETP Lett. 84, 199 (2006)].

[29] E. I. Grusharova and V. I. Perel’, “Spin relaxation in semiconductors caused by electric fields,” Fiz. Tekh. Poluprovodn. 10, 2272 (1976).

[30] L. E. Golub and E. L. Ivchenko, “Spin splitting in symmetrical SiGe quantum wells,” Phys. Rev. B 69, 115333 (2004)