Spintronics [1,2], which explores the application of the spin degree of freedom to electronics, is one of the major subjects in applied physics. As a milestone in spintronics, Datta and Das have proposed a spin field-effect transistor (FET) [3], in which the spin orientation of an electron, in transport through a two-dimensional channel, rotates around the in-plane effective magnetic field \( \mathbf{B}_{\text{eff}} \) induced by the Rashba spin-orbit interaction [4–7]. This effective magnetic field switches the control current on and off each time the change of \( \mathbf{B}_{\text{eff}} \) increases the angle of the spin rotation in the channel by \( \pi \). Therefore the action of this spin FET relies on the control of \( \mathbf{B}_{\text{eff}} \) by the gate voltage, that is, by the external electric field perpendicular to the plane \( E_{z}^{\text{ex}} \), which has been confirmed by experiments in quantum wells [8–10].

Another spin FET [11,12], called the spin-lifetime FET, has been proposed by Hall and others in which the spin relaxation switches the control current by changing the electron distribution from fully spin-polarized to unpolarized. The spin-lifetime FET, in contrast to the Datta-Das spin FET, uses only one component of the spin polarization vector. The action of the spin-lifetime FET relies on the switching of the spin relaxation rate by the change of the gate voltage. As a method to vary the spin-relaxation rate, the original paper [11,12] proposed to use the above-mentioned \( E_{z}^{\text{ex}} \) dependence of \( B_{\text{eff}} \) \( (\propto E_{z}^{\text{ex}}) \) in quantum wells [8–10], which makes the spin-relaxation rate in the Dyakonov-Perel mechanism [13–15] \( \propto (B_{\text{eff}})^{2} \) proportional to \( (E_{z}^{\text{ex}})^{2} \). The decrease of \( E_{z}^{\text{ex}} \) from a higher value \( E_{z}^{\text{H}} \) to a lower one \( E_{z}^{\text{L}} \) reduces the current from \( I_{\text{on}} \propto (E_{z}^{\text{H}})^{2} \) to \( I_{\text{off}} \propto (E_{z}^{\text{L}})^{2} \), leading to the on/off current ratio \( I_{\text{on}}/I_{\text{off}} = (E_{z}^{\text{H}}/E_{z}^{\text{L}})^{2} \). Unfortunately, this method requires a precise control of the gate voltage to achieve a high on/off current ratio because \( E_{z}^{\text{L}} \) must be in the close vicinity of \( E_{z}^{\text{ex}} = 0 \).

We have recently found [16] that \( B_{\text{eff}} \) due to the Rashba spin-orbit interaction and the resulting spin-relaxation rate vary in a wide range by changing the band offsets (of conduction, valence, and split-off bands) between the well and barrier semiconductors in a quantum well. In particular, the coefficient \( \alpha \) of the Rashba spin-orbit interaction (and of \( B_{\text{eff}} \)) can be tuned to vanish by adjusting the band offsets. For example, in a quantum well consisting of Ga0.47In0.53As (well) and AlxGa1−xAs/Sb1−y (barrier), where the band offsets change with the Al fraction \( x \), \( \alpha \) and the associated spin-relaxation rate become zero at an optimum fraction \( x_{0} \).

In this paper we propose and explore the gate-voltage-induced switching of the Rashba coefficient \( \alpha \) by using the above-mentioned band-offset dependence of \( \alpha \). The first key factor of the proposed switching is an asymmetric quantum-well structure with two barrier layers formed by different semiconductors, AlxGa1−xAs/Sb1−y \( (x \approx x_{0}) \) and Al0.48In0.52As. This combination of barrier semiconductors is chosen so that one (the other) interface is that in a symmetric quantum well with a small \( |\alpha| \) (a large \( |\alpha| \)). The second is the wave function deformation due to the external electric field \( E_{z}^{\text{ex}} \), which is produced by the gate voltage. Then the gate voltage switches \( |\alpha| \) on and off as the external electric field moves the wave function to the interface with a large \( |\alpha| \) and to that with a small \( |\alpha| \). We examine the on/off \( |\alpha| \) ratio when the Al fraction \( x \) in AlxGa1−xAs/Sb1−y deviates from the optimum value \( x_{0} \) and investigate the required \( E_{z}^{\text{ex}} \).

We consider an electron in the conduction band of a quantum-well structure which is formed by three different semiconductors with the zinc-blende structure (Fig. 1): \( S_{B}^{L} \) in the left barrier layer \( (z < z_{1}^{L}) \), \( S_{W}^{L} \) in the well layer \( (z_{1}^{L} < z < z_{1}^{B}) \), and \( S_{B}^{R} \) in the right barrier layer \( (z < z_{1}^{B}) \). Due to the translational symmetry along the \( x \) and \( y \) axes, the associated wave numbers, \( k_{x} \) and \( k_{y} \), are conserved. The Schrödinger equation for the ground-subband wave function \( \varphi_{0}(z) \) and the corresponding eigenvalue \( \epsilon_{0} \) is

\[
\hat{p}_{z}^{2} + \frac{\hbar}{2m} V_{W}(z) \varphi_{0}(z) = \epsilon_{0} \varphi_{0}(z),
\]

where \( \hat{p}_{z} = -i \hbar \nabla_{z} = -i \hbar \partial_{z} / \partial z \) and \( m \) is the effective mass of the conduction band. The confining potential \( V_{W}(z) \) is

\[
V_{W}(z) = V_{B}^{c}(z) + V_{\alpha}(z).
\]

Here \( V_{\alpha}(z) \) is the potential due to the conduction-band offset (Fig. 1), the expression of which is given later with those for valence bands. The second term \( V_{\alpha}(z) \), which is the
The electrostatic potential due to the charge distribution in the quantum-well structure satisfies the Poisson equation

\[ \nabla^2 V_{\text{es}} = -A n(z), \]  

where

\[ A = \frac{4\pi e^2}{\varepsilon}, \quad n(z) = N_s [\phi_0(z)]^2, \]

with \( \varepsilon \) the static dielectric constant, \(-e\) the electron charge (\( e > 0 \)), and \( N_s \) the sheet density of electrons. The total electric field is \( E_z = \nabla V_{\text{es}}/e \), while the external electric field \( E_z^{\text{ex}} \) is that induced by ionized donors and charges on the gate electrode. Such charges, which are placed in \( z < z_{\text{SL}} \) and \( z > z_{\text{SR}} \), are assumed to be far away from the quantum well so that \( E_z^{\text{ex}} \) acting on electrons is constant (this assumption can be satisfied by employing the widely used modulation doping where the spacer layer thicker than 10 nm is inserted between the well layer and each doped layer so that the penetration of the wave function into each doped layer is negligible). We introduce the areal external-charge density in the left side of the well \( (z < z_{\text{SL}}) \), \( \sigma_L \), and that in the right side \( (z > z_{\text{SR}}) \), \( \sigma_R \). Then the boundary condition for \( V_{\text{es}} \) at \( z = z_{\text{SL}} \) becomes

\[ \nabla_z V_{\text{es}} = A (\sigma_L/e) \left( z = z_{\text{SL}} \right). \]  

The charge neutrality gives \( e N_s = \sigma_L + \sigma_R \), while \( E_z^{\text{es}} = (2\pi /e) (\sigma_L - \sigma_R) \).

We derive the formula for the spin-orbit interaction induced by the band offsets and the electrostatic potential, which is denoted by \( V_{\text{W}}^{\text{so}} \), as well as the Rashba coefficient \( \alpha \) in the case of an asymmetric quantum well consisting of three different semiconductors \[17\]. In deriving \( V_{\text{W}}^{\text{so}} \) we employ the \( k \cdot p \) theory developed for heterostructures \[18,19\] (it has been shown in a number of papers \[9,10,20–24\] that the Rashba coefficient \( \alpha \) derived by the \( k \cdot p \) theory in various heterostructures agrees well with the experimental value and with that calculated in the tight-binding model). Then we obtain, for an electron with wave numbers \( k_x \) and \( k_z \) in the conduction band,

\[ V_{\text{W}}^{\text{so}} = \frac{p^2}{3} G(z)(\sigma_z k_y - \sigma_y k_z), \]  

where \( \sigma_x \) and \( \sigma_y \) are the Pauli spin matrices and \( P \) is the Kane matrix element \[25\]. In this equation,

\[ G(z) = \nabla_z \left( \frac{1}{E - \tilde{E}_v} - \frac{1}{E - \tilde{E}_s} \right), \]  

where \( E \) is the electron energy measured from the conduction-band bottom of \( S_W \), while \( \tilde{E}_v \) (\( \tilde{E}_s \)) is the diagonal element of the 8 × 8 Kane Hamiltonian, corresponding to the heavy-hole plus light-hole bands (the split-off band). They are given by

\[ \tilde{E}_v = E_{K0} - E_g + V_{\text{bo}}(z) + V_{\text{es}}(z), \]

\[ \tilde{E}_s = E_{K0} - E_g + V_{\text{bo}}(z) + V_{\text{es}}(z). \]  

Here \( E_{K0} = (\hbar^2/2m_0)(k_x^2 + k_y^2 + k_z^2) \) with \( m_0 \) the electron rest mass and \( \tilde{k}_z = -i\nabla_z \), and \( V_{\text{bo}}(z) = E_g + \Delta_{\text{so}} \) with \( E_g \) (\( \Delta_{\text{so}} \)) the band gap (the spin-orbit splitting) of \( S_W \), while \( V_{\text{es}}(z) [V_{\text{bo}}(z)] \) is the potential due to the band offset for an electron in the heavy-hole plus light-hole bands [the split-off band]. These potentials together with that for the conduction band, \( V_{\text{bo}}(z) \) in Eq. (2), are expressed by

\[ V_{\text{bo}}(z) = V_{\text{bo}}^{\ell}(z) + V_{\text{bo}}^{R}(z), \quad V_{\text{bo}}^{\ell}(z) = \Delta E_s^{\ell} h_{\ell}(z), \]

where \( i = c, v, s, \ell = L, R \), and

\[ h_L(z) = \begin{cases} 1 & (z < z_{\text{SL}}^L), \\ 0 & (z > z_{\text{SL}}^L), \end{cases} \quad h_R(z) = \begin{cases} 0 & (z < z_{\text{SR}}^R), \\ 1 & (z > z_{\text{SR}}^R). \end{cases} \]  

while \( \Delta E_s^{\ell} \), \( \Delta E_c^{\ell} \), and \( \Delta E_{\ell}^{\ell} \) are the band offsets of the semiconductor \( S_\ell \) relative to \( S_W \): \( \Delta E_s^{\ell} = E_s^{\ell} - E_c \), \( \Delta E_c^{\ell} = E_c - E_s \), and \( \Delta E_{\ell}^{\ell} = E_{\ell} - E_s \) with \( E_{\ell}(E_c') \) the energy of the conduction-band bottom, \( E_s(E_c') \) that of the valence-band top, \( E_s(E_c') \) that of the split-off-band top in \( S_W \) (in \( S_R \)). We can also express \( E_g \) and \( \Delta_{\text{so}} \) as \( E_g = E_c - E_s \) and \( \Delta_{\text{so}} = E_s - E_c \).

Here we neglect \( E_v \), \( E_{K0} \), and \( V_{\text{es}}(z) \) in \( G(z) \), compared to \( E_g \) and \( E_s^\ell \) while we take into account the contribution from \( \nabla_z V_{\text{bo}}(z) \) to \( G(z) \), that is, \( [(E_s - V_{\text{bo}})^2 - (E_s - V_{\text{es}})^2]\nabla_z V_{\text{bo}} \), in which we neglect \( V_{\text{es}}^\ell \) and \( V_{\text{bo}}^\ell \) since they are nonzero only in the barrier layers where the squared wave function is small (we have confirmed by the numerical calculation that the correction due to \( V_{\text{es}}^\ell \) and \( V_{\text{bo}}^\ell \) in this contribution is less than 6 percent in the cases we considered in this paper). We finally obtain the expression for \( V_{\text{W}}^{\text{so}} \):

\[ V_{\text{W}}^{\text{so}} = a(z)(\sigma_z k_y - \sigma_y k_z), \]  

with

\[ a(z) = \eta \left[ \nabla_z (b_{\text{off}}^{\ell} V_{\text{bo}}^{\ell} + b_{\text{off}}^R V_{\text{bo}}^R + V_{\text{es}}) \right]. \]  

Here \( \eta \) is the effective coupling constant of the spin-orbit interaction for an electron in the conduction band of the semiconductor \( S_W \), given by

\[ \eta = \frac{p^2}{3} \left[ \frac{1}{(E_g)^2} - \frac{1}{(E_s)^2} \right]. \]  

and \( b_{\text{off}}^{\ell} (\ell = L, R) \) is defined by \[16\]

\[ b_{\text{off}}^{\ell} = \Delta E_{\ell}^{\ell} / \left[ E_g (E_g - \Delta E_{\ell}^{\ell}) \right] - \Delta E_{\ell}^{\ell} / \left[ E_g (E_g + \Delta E_{\ell}^{\ell}) \right] \]  

\[ \Delta E_{\ell}^{\ell} = (1/2)(E_{\ell} - E_g)^2 / (1/(E_g)^2). \]  

Equation (11) with (12) shows that the spin-orbit interaction \( V_{\text{W}}^{\text{so}} \) for an electron in the conduction band, due to the band offsets and the electrostatic potential, is not proportional to \( \nabla_z V_{\text{bo}} \) except the case where \( b_{\text{off}}^{\ell} = 1 \) and \( b_{\text{off}}^R = 1 \).

The Rashba coefficient \( \alpha \) is defined by the expectation value of \( a(z) \) with respect to \( \phi_0(z) = \langle z | 0 \rangle \),

\[ \alpha = \langle 0 | a(z) | 0 \rangle. \]
Here we use the equality \( \langle 0 | \nabla_z V_W | 0 \rangle = 0 \), which means that forces on an electron in a bound eigenstate are balanced \[26, 27\] (the equality is derived according to \[28\] in \[29\]). Then we obtain the expression of \( \alpha \) using values of the wave function at interfaces:

\[
\alpha = \eta (b_{\text{off}}^R - 1)(0 | \nabla_z V_{\text{bo}}^{(1)} | 0) + \eta (b_{\text{off}}^R - 1)(0 | \nabla_z V_{\text{bo}}^{(2)} | 0)
\]

\[
= \eta (b_{\text{off}}^R - 1)[\langle \phi_0(z_{1}^R) \rangle^2 - (\Delta E_{1}^R) + \eta (b_{\text{off}}^R - 1)[\langle \phi_0(z_{2}^R) \rangle^2]
\times \Delta E_{2}^R.
\]

First we assume, for simplicity, that the wave function vanishes at the left heterointerface, that is, \( \phi_0(z_{1}^R) = 0 \) [the numerically calculated wave function, shown in Fig. 3(b), is reduced considerably at the left heterointerface by applying a large negative \( E_{z}^{\text{ex}} \) although it does not vanish completely]. In this case, from Eq. (16), we have \( \alpha = \eta (b_{\text{off}}^R - 1)(0 | \nabla_z V_{\text{bo}}^{(1)} | 0) \), which vanishes at \( b_{\text{off}}^R = 1 \). In Fig. 2, we plot \( b_{\text{off}}^R \) for \( S_{\text{W}} = \text{Ga}_{0.47}\text{In}_{0.53}\text{As} \) and \( S_{\text{R}}^\text{Ga} = \text{Al}_{1-x}\text{Ga}_{x}\text{As} \) as a function of the Al fraction \( x \) in \( S_{\text{R}}^\text{Ga} \) where \( y \) is determined so that \( S_{\text{R}}^\text{Ga} \) is lattice-matched to \( S_{\text{W}} \) \[16\]. This figure shows that \( b_{\text{off}}^R = 1 \) at a value \( x_0 \) which is close to \( x = 0.3 \) for the values of band parameters from Ref. \[30\]. Therefore, an asymmetric quantum well consisting of \( S_{\text{W}} = \text{Ga}_{0.47}\text{In}_{0.53}\text{As} \) and \( S_{\text{R}}^\text{Ga} = \text{Al}_{0.43}\text{In}_{0.57}\text{As} \) \( (b_{\text{off}}^R = -0.32) \), and \( S_{\text{R}}^\text{Ga} = \text{Al}_{1-x}\text{Ga}_{x}\text{As} \) with \( x = x_0 \) \( (b_{\text{off}}^R = 1) \) has \( |\alpha| \approx 0 \) for a large negative \( E_{z}^{\text{ex}} \) where \( \phi_0(z_{1}^R) \) is considerably reduced, while it acquires a substantial value of \( |\alpha| \) for a large positive \( E_{z}^{\text{ex}} \) where \( \phi_0(z_{2}^R) \) is increased. This means that \( |\alpha| \) is switched on and off by changing the polarity of \( E_{z}^{\text{ex}} \).

Now we evaluate the on/off ratio of \( |\alpha| \), denoted by \( R_{\alpha} \), when \( x \) slightly deviates from \( x_0 \). We start from the expression of \( \alpha \) in the case of \( \phi_0(z_{1}^R) = 0, \alpha = \eta (b_{\text{off}}^R - 1)(0 | \nabla_z V_{\text{bo}}^{(1)} | 0) \), and \( (0 | \nabla_z V_{\text{bo}}^{(1)} | 0) = 0 \). By using \( \phi_0(z_{1}^R) = 0 \) as well as \( (0 | \nabla_z V_{\text{bo}}^{(1)} | 0) = 0 \) and \( (0 | \nabla_z V_{\text{bo}}^{(2)} | 0) = eE_{z}^{\text{ex}} \), we have \( (0 | \nabla_z V_{\text{bo}}^{(1)} | 0) = 0 \) and \( (0 | \nabla_z V_{\text{bo}}^{(2)} | 0) = -eE_{z}^{\text{ex}} \) and then obtain \( \alpha = \eta (b_{\text{off}}^R - 1)(1 - eE_{z}^{\text{ex}}) \). Similarly we obtain \( \alpha = \eta (b_{\text{off}}^R - 1)(1 - eE_{z}^{\text{ex}}) \) in the case of \( \phi_0(z_{2}^R) = 0 \). Then the on/off ratio \( R_{\alpha} \), when the polarity of \( E_{z}^{\text{ex}} \) is changed with a large enough \( |E_{z}^{\text{ex}}| \), becomes

\[
R_{\alpha} = \frac{|\alpha|_{\text{on}}}{|\alpha|_{\text{off}}} = \frac{|b_{\text{off}}^R - 1|}{|b_{\text{off}}^R - 1|}.
\]

For a deviation of \( x - x_0 \) = 0.01, we have \( |b_{\text{off}}^R - 1| \approx 1/30 \) from Fig. 2. Therefore we obtain \( R_{\alpha} \approx 40 \), which leads to the on/off ratio of the spin-relaxation rate \( (R_{\alpha}) \) exceeding \( 10^3 \) in the Dyakonov-Perel mechanism where \( R_{\alpha} = (R_{\alpha})^2 \).

In order to examine how large \( |E_{z}^{\text{ex}}| \) is necessary to realize the on/off \( |\alpha| \) ratio \( R_{\alpha} \) in Eq. (17), we perform a numerical calculation for \( \psi_0(z) \) by solving the Schrödinger equation Eq. (1) and the Poisson equation Eq. (3) self-consistently, in which we discretize the \( z \) coordinate. The obtained \( \psi_0(z) \) as well as \( \psi_0(z) \) in \( \text{Al}_{0.48}\text{In}_{0.52}\text{As}/\text{Ga}_{0.47}\text{In}_{0.53}\text{As}/\text{Al}_{1-x}\text{Ga}_{x}\text{As}, x = 0.3 \) are plotted in Fig. 3 for five different values of the external force \( F_z = -eE_{z}^{\text{ex}} \) (the energy separation between the ground subband and the first-excited subband increases with \( |E_{z}^{\text{ex}}| \) and the excited subbands are not occupied by electrons in the cases we considered in this paper), while the value of \( \alpha \) as a function of \( F_z \) is presented in Fig. 4 for three different quantum wells. All of the three quantum wells have a common material in the well layer, \( \text{Ga}_{0.47}\text{In}_{0.53}\text{As} \). The difference of the three is in the combination of interfaces, LL, SS, and LS, where L \{S\} is the interface which gives a large \( \delta \{s\} \mid b_{\text{off}}^R = 1, \text{Ga}_{0.47}\text{In}_{0.53}\text{As}(\text{well})/\text{Al}_{0.45}\text{In}_{0.55}\text{As}(\text{barrier}) \) with \( b_{\text{off}}^R = 1 - 1.32 \{\text{Ga}_{0.47}\text{In}_{0.53}\text{As}(\text{well})/\text{Al}_{1-x}\text{Ga}_{x}\text{As}, x = 0.3 \) with \( b_{\text{off}}^R = 1 - 0.05 \).
FIG. 4. Rashba coefficient \( \alpha \) as a function of the external force \( F_z = -eE_{\text{ex}}^z \) for three combinations of interfaces: LL, SS, and LS (see text). (a) The well-width \( (W) \) dependence for LS. (b) The electron-sheet-density \( (N_s) \) dependence for LS. The dimensionless Rashba coefficient is defined by \( \tilde{\alpha} = (\alpha/\eta)/(R\gamma^*/a_B^*) \). For the definition of other quantities, see the caption of Fig. 3.

In the case of \( \hbar^2 a_{\text{eff}}^L = \hbar^2 b_{\text{eff}}^S \) as in LL and SS, Eq. (16) becomes \( \alpha = \eta(b_{\text{eff}} - 1)(0)(\nabla V_{\text{ex}}^L)(0) = -\eta(b_{\text{eff}} - 1)(0)(\nabla V_{\text{ex}}^S)(0) = -\alpha_{\text{off}} \). We have used the equalities \( (0)(\nabla V_{\text{ex}}^L)(0) = 0 \) and \( (0)(\nabla V_{\text{ex}}^S)(0) = eE_{\text{ex}}^z = -F_z \). Therefore \( \alpha \) of LL and that of SS are independent of the well width \( W \) and \( N_s \). They are plotted in Fig. 4 as straight lines with different gradients.

The calculated \( \alpha(F_z) \) of LS, presented in Fig. 4, approaches the steep-slope \( \alpha(F_z) \) of LL in \( F_z < 0 \) and the gentle-slope \( \alpha(F_z) \) of SS in \( F_z > 0 \) as \( |F_z| \) increases, which demonstrates the switching of \( |\alpha| \) by changing the sign of \( F_z \) (\( \alpha \) of LS deviates from zero at \( F_z = 0 \), which is derived in [31]). A rough estimate for the value of \( \alpha \) at \( F_z = 0 \), at which \( \alpha(F_z) \) of LS merges with that of LL or SS, denoted by \( F_z^* \), is given by the potential difference between two interfaces \( eE_{\text{ex}}^z W = |F_z| \) equal to the energy difference between the ground subband and the first excited subband at \( N_s = 0 \) and \( E_{\text{ex}}^z = 0 \) in the infinite-barrier model \( \epsilon_1 - \epsilon_0 = 3\hbar^2\pi^2/(2mW^2) \). This estimate shows a strong \( W \) dependence of \( F_z^*: F_z^* \propto W^{-3} \), which explains the \( W \) dependence (at \( N_s = 0 \)) presented in Fig. 4(a).

Using dimensionless variables, \( F_z^* = F_z^*/(R\gamma^*/a_B^*) \) and \( \tilde{W} = W/a_B^* \) with \( a_B^* = \hbar^2/\epsilon m^2 \) and \( \gamma^* = \hbar^3/(2m^2) \), such an estimate is expressed as \( F_z^* = 3\pi^2\tilde{W}^{-3} \), which gives \( F_z^* = 137 \) for \( \tilde{W} = 0.6 \) and \( F_z^* = 17.1 \) for \( \tilde{W} = 1.2 \). This estimate for \( F_z^* \) is consistent with that extracted from Fig. 4(a). The corresponding value of \( |E_{\text{ex}}^z| \), denoted by \( E_{\text{ex}}^z \), becomes \( E_{\text{ex}}^z = 2.4 \times 10^7 \) V/m \( (W = 10.3 \) nm \) and \( E_{\text{ex}}^z = 3.0 \times 10^7 \) V/m \( (W = 20.5 \) nm \) when \( a_B^* = 17.1 \) nm and \( \gamma^* = 3.04 \) meV for \( Ga_{0.47}In_{0.53}As \) is used. On the other hand, the dependence of \( \alpha \) on the electron sheet density \( N_s \) is weak as shown in Fig. 4(b). We therefore find that \( |E_{\text{ex}}^z| \) above \( E_{\text{ex}}^z = 3\pi^2\tilde{W}^{-3}R\gamma^*/(ea_B^*) \) is required to realize the on/off \( |\alpha| \) ratio \( R_{\alpha} \) given by Eq. (17).

From the numerical result in Fig. 4(a) showing that \( \alpha \) of LS for \( \tilde{W} = 1.2 \) is well approximated by \( \alpha \) of LL \( (F_z < 0) \) and \( \alpha \) of SS \( (F_z > 0) \) at \( |F_z| > 50 \), it is derived that \( R_{\alpha} \) reaches the value in Eq. (17) at \( |E_{\text{ex}}^z| > 10^7 \) V/m for \( W = 20.5 \) nm (\(|E_{\text{ex}}^z| = 3 \times 10^7 \) V/m has been experimentally attained [32]).

In conclusion, we have calculated the coefficient of the Rashba spin-orbit interaction, \( \alpha \), as a function of the external electric field perpendicular to the well, \( E_{\text{ex}}^z \), in an asymmetric quantum well consisting of \( Ga_{0.47}In_{0.53}As \) (well), \( Al_{0.48}In_{0.52}As \) (left barrier), and \( Al_{0.3}Ga_{0.7}As \) (right barrier). We have found that \( |\alpha| \) and the resulting spin-relaxation rate in the Dyakonov-Perel mechanism can be switched on and off by changing the polarity of \( E_{\text{ex}}^z \) when \( |E_{\text{ex}}^z| \) is large enough. The required \( |E_{\text{ex}}^z| \) is proportional to \( W^{-3} \) with \( W \) the well width, and weakly depends on the electron sheet density.

The Dresselhaus spin-orbit interaction [33] also contributes to the spin relaxation in the Dyakonov-Perel mechanism in addition to the Rashba spin-orbit interaction which has been considered in this paper. However, in a quantum well parallel to the (110) plane of the zinc-blende structure, the effective magnetic field induced by the Dresselhaus spin-orbit interaction is perpendicular to the well layer and does not give the spin relaxation in the Dyakonov-Perel mechanism for the spin component perpendicular to the well layer. This has been shown theoretically in [15] and the resulting reduction in the spin relaxation rate has been demonstrated experimentally in [34–37]. Therefore the estimate of the spin relaxation rate, given in this paper, does not change in (110) quantum wells for the perpendicular spin component even when the Dresselhaus spin-orbit interaction is considered in addition to the Rashba spin-orbit interaction.

[17] E. A. de Andrada e Silva, G. C. La Rocca, and F. Bassani, Phys. Rev. B 55, 16293 (1997). This paper describes a theory on the spin-orbit splitting for an asymmetric quantum well Al$_x$Ga$_{1-x}$As/GaAs/Al$_y$Ga$_{1-y}$As ($x \neq y$) in the absence of the electrostatic potential $[V_{es}(z) = 0]$.


[29] From Eq. (1) $\psi_0(z) = \langle z | 0 \rangle$ is the eigenfunction corresponding to the eigenvalue $\epsilon_0$ of the Hamiltonian $\hat{H} = \hat{p}_z^2/(2m) + V_W(z)$ with $[\hat{p}_z, \hat{H}] = -i\hbar(\nabla_z V_W)$. Since $\langle 0 | [\hat{p}_z, \hat{H}] | 0 \rangle = 0$, we obtain $\langle 0 | (\nabla_z V_W) | 0 \rangle = 0$.


[31] $\alpha$ in Eq. (16) can also be written as the sum of a term proportional to $F_z = -eE x z$ and that proportional to $b_{av}^{L,R}$:

$$\alpha = \eta(b_{av}^{L} - 1)F_z + \eta(b_{av}^{L} - b_{av}^{R})\langle 0 | V_{V_{cu}^{L}} - V_{V_{cu}^{R}} | 0 \rangle,$$

with $b_{av}^{L,R} = (b_{av}^{L} + b_{av}^{R})/2$. At $F_z = 0$ the first term is zero, while the second term is nonzero in an asymmetric quantum well with $b_{av}^{L,R} \neq b_{av}^{L,R}$.


