Giant edge spin accumulation in a symmetric quantum well with two subbands

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Abstract – We have studied the edge spin accumulation due to an electric current in a high mobility two-dimensional electron gas formed in a symmetric well with two subbands. This study is strongly motivated by the recent experiment of Hernandez et al. [Phys. Rev. B, 88 (2013) 161305(R)] who demonstrated the spin accumulation near the edges of a symmetric bilayer GaAs structure in contrast to no effect in a single-layer configuration. The intrinsic mechanism of the spin-orbit (SO) interaction we consider arises from the coupling between two subband states of opposite parities. We obtain a parametrically large magnitude of the edge spin density for a two-subband well as compared to the usual single-subband structure. We show that the presence of a gap in the system, i.e., the energy separation Δ between the two subband bottoms, changes drastically the picture of the edge spin accumulation. The gap value governs the effective strength of the inter-subband SO interaction which provides a controllable crossover from the regime of weak spin accumulation to the regime of the strong one by varying the Fermi energy (electron density) and/or Δ. We estimate that by changing the gap Δ from zero up to 1–2 K, the magnitude of the effect changes by three orders of magnitude. This opens up the possibility for the design of new spintronic devices.

Introduction. – Spin current and spin accumulation [1,2] which appear due to the spin-orbit (SO) coupling in the presence of electric currents are topics of great current interest which are important for the future of spin electronics [3]. There are two distinct SO mechanisms, the extrinsic one due to the Mott asymmetry in the electron scattering off impurities [4–7], and the intrinsic one [8,9] due to SO-induced splitting of the electron spectrum. The edge spin-density accumulation, related to either the Mott asymmetry by impurities [10] (2D electrons) or the intrinsic mechanism (2D holes) [11,12], has been experimentally observed.

It is known [2,13–16] that in the diffusive regime (and when the spin diffusion length is much larger than the mean free path) the edge spin density is entirely due to the spin flux coming from the bulk. In contrast, the physics of the edge spin-density accumulation for the intrinsic mechanism in the opposite case of strong SO splitting [17] only recently has been understood [18–22]. This includes both the case of a ballistic structure in which the mean free path is the largest lengthscale in the problem, and the experimentally important case of a diffusive sample when the mean free path being much smaller than the sample size is much larger than the spin-precession length. We term the latter case the quasi-ballistic regime. In particular, it has been shown [21,22] that in this regime there is no relation between edge spin density and bulk spin current. In other words, the edge spin density is mostly created by the carriers scattering off the boundary itself (the boundary mechanism). For example, in the case of 2D holes in the quasi-ballistic regime the edge spin density, which is due to the spin current from the bulk, is parametrically smaller than the density generated upon the boundary scattering [22].

Recently, using Kerr rotation spectroscopy, Hernandez et al. [23] demonstrated electric-current–induced spin accumulation near the edges of a high-mobility two-dimensional electron gas in a symmetric two-subband...
GaAs structure in contrast to no accumulation in a single-subband configuration. This result is interesting and intriguing in many aspects. The observed effect is quite large despite the fact that the electric field in the high-mobility channel is 300–400 times smaller than that in the experiment by Kato et al. [10], where for a GaAs sample the result was explained by the extrinsic interaction with impurities. Note that the structure studied in [23] has inversion symmetry and, therefore, the usual Rashba term [24] is absent. On the other hand, the linear-in-momentum term [25] originating from a cubic Dresselhaus term is known not to lead to a spin current in the bulk. In addition, it cannot create a smooth, within the scale of the spin-precession length, edge spin density upon boundary scattering for any reasonably short-ranged impurity scattering in the bulk, as is shown in ref. [22], see eq. (1). A significant difference between the observed edge spin density in the two-subband vs. the one-subband cases has motivated us to look for the explanation of this phenomenon using the inter-subband Rashba-like Hamiltonian arising in two-subband wells [26,27]. We will further disregard the intra-band linear Dresselhaus term [25].

Here we follow the method proposed in [22] to calculate the edge spin density which appears due to boundary scattering (fig. 1(a)) in the quasi-ballistic regime for a Rashba-like Hamiltonian [26,27] describing the two-subband well (fig. 1(b)). In this quasi-ballistic regime the characteristic length of the spin accumulation near the boundary is smaller than the mean free path. Since the latter is around 30 μm [23], it indeed exceeds all the characteristic lengths of our theory. We have explained the experimental results, in particular, the large magnitude of the edge spin density for the two-subband sample compared to the usual single-band structure with either the Rashba or Dresselhaus interactions.

Two bands vs. one band case. – Interestingly, we have found that despite the problem in question resembling very much the usual Rashba problem (there are two copies of them because each state is doubly degenerate), the presence of the gap Δ between two sub-band edges (fig. 1(b)) changes the physics of the edge spin accumulation completely. This happens because the gap magnitude governs the effective strength of the inter-subband SO interaction leading to different solutions compared to the one-band Rashba case for the occupation numbers of the incoming states participating in the boundary scattering (fig. 1(a)).

The physics is now determined by the value of the parameter \( \xi = 2\eta p_F/\Delta \equiv L_\Delta/L_s \) (fig. 2(a)), where \( L_\Delta = \hbar v_F/\Delta \), \( L_s = \hbar^2/(2m\eta) \) are the coherence and the spin-precession lengths. Here \( p_F = \hbar k_F = n v_F \) is the Fermi momentum, \( \eta \) is the inter-subband SO coupling constant [26,27]. The parameter \( \xi \) can have an arbitrary value even for a small \( \eta \) since the gap \( \Delta \) can be made much smaller than the Fermi energy.

The eigenvectors of the Hamiltonian (2), see eqs. (S2), (S3) in the Supplemental Material (SM), are determined now by two angles. Angle \( \varphi \) is the polar angle of vector \( p \) within the 2D plane. The angle \( \theta(p) \) which depends on the modulus of the momentum and is defined with respect to the \( z \)-axis (normal to the 2D plane) determines the effective strength of the SO coupling. This effective strength is given by the function \( \sin \theta = (2\eta p/\hbar)\sqrt{(2\eta p/\hbar)^2 + \Delta^2} \) . When \( \Delta \) is much bigger than the SO energy \( 2\eta p/\hbar \), then the \( z \)-axis can be chosen as a “true” spin quantization axis. The usual Rashba single-band model (in our case we have two copies of it) corresponds to a limit of strong SO coupling when \( \Delta \to 0 \) and the SO energy dominates \( \theta \to \pi/2 \), \( \sin \theta = 1 \).

As has been shown in [21,22] (see also eq. (7) below), because of the unitarity of the boundary scattering the magnitude of the edge spin accumulation is proportional to the difference \( f_1(\varepsilon F, k_y) - f_2(\varepsilon F, k_y) \) between the distribution functions of the incoming electron states belonging to the subbands 1 and 2 for a given Fermi energy and a given wave vector \( k_y \) along the boundary, see fig. 1. (In the case of the one-band Rashba model the incoming states belong to the branches of opposite helicities.) These distribution functions are found from the solution of the kinetic equation for the spin-density matrix in the bulk (2D) of the sample in the presence of impurity scattering and an electric field [28,29] and used as the input parameters for the boundary scattering problem.

The important point is that in the one-band Rashba case the difference of the distribution functions in question is of the third order with respect to \( p_1 - p_2 = \hbar/L_s \) for any reasonably short-ranged impurity potential in the bulk.
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(when the correlation radius of the impurity potential \(d\) is much smaller than \(L_s\)). This has been rigorously proven in ref. [22] (see eqs. (14)–(16) in [22]). Thus, in the leading order (i.e., \(\propto 1/L_s\)) the above-mentioned distribution functions take the form which they have in the absence of any spin-orbit coupling. The first-order effect appears only for very smooth impurity potential when \(d > L_s\), and the magnitude of the edge spin density for the one-band Rashba case is given by the expressions [22]

\[
\langle S_z \rangle \simeq \frac{k_F d^2}{L_s^3} \quad \text{at} \quad d \ll L_s; \quad \langle S_z \rangle \simeq \frac{k_F}{L_s} \quad \text{at} \quad d > L_s.
\]

(1)

Here \(k_F = eE\tau_F/\hbar\), \(e\) is the modulus of the electron charge, \(E\) is the magnitude of the in-plane (driving) electric field directed along the \(y\)-axis and \(\tau_F\) is the transport scattering time due to the impurities in the bulk of a sample. Note that we consider in our paper the case \(d \ll L_s\), which is only the realistic one.

Note that the cancellation of the effect for the Rashba model in the leading order (i.e., \(\propto 1/L_s\)) occurs when the transition rates between branches of opposite helicities have the same strength as the transition rates within the branch of the same helicity (one of the necessary conditions). That is why for smooth impurity potential \((d > L_s)\) which cannot support the transitions between branches of opposite helicities, one observes the recovering of the first-order effect, see eq. (1).

From the above considerations we can immediately understand the role of the gap \(\Delta\) in the two-subband model considered in this work. As explained above, the effective SO interaction decreases with increasing \(\Delta\). This causes suppression of the inter-subband transition rates, since these transitions are accompanied by spin flip, as compared to the intra-subband ones, which do not need spin flip. The suppression factor is \(\sin^2 \theta\), which is the probability of spin-flip (see also SM, sect. III). Since the inter-band and intra-band rates are different now, this prevents the complete cancellation that occurs for the quantity \(f_1(\varepsilon_F, k_y) - f_2(\varepsilon_F, k_y)\) in the one-band Rashba model, and leads to the recovery of the first-order effect with respect to small splitting \(p_1 - p_2 \ll p_F\) even for a short-ranged \((d \ll L_s)\) impurity potential.

As we discuss later on, for a GaAs structure similar to that used in ref. [23], it is enough to change the gap \(\Delta\) from zero up to about 1–2 K in order to increase the magnitude of the effect by the three orders ("giant effect"). The most pronounced change occurs at \(L_\Delta \simeq 0.5L_s\), where the edge spin density \(S_z(x)\) is maximized and its order of magnitude is given by \(k_F/L_s\), which is parametrically larger than in the single-band case (fig. 2(a)).

We consider specular scattering (i.e., a straight boundary for simplicity, fig. 1(a), assuming also zero rate of spin relaxation at the edge) and a Fermi energy much larger than the gap \(\Delta\) between the subbands, i.e., \(\varepsilon_F \gg \Delta\) (see footnote 2). Moreover, the SO interaction is weak \((\eta_F \ll \varepsilon_F)\) and therefore "coherent" and spin-precession length scales are large compared to the Fermi wavelength, \(L_\Delta, L_s \gg \lambda_F = 2\pi/k_F\). The ratio \(L_\Delta/L_s\) can be arbitrary. Our calculation shows that the characteristic spatial scale of the edge spin density is \(\Lambda = L_\Delta L_s/\sqrt{L_\Delta^2 + L_s^2}\).

Model Hamiltonian. – The Hamiltonian of a symmetric quantum well with two subbands and inter-subband–induced SO interaction resembles that of the ordinary Rashba model. In contrast to the latter, the intersubband SO interaction is nonzero even in symmetric structures with the 4 × 4 Hamiltonian being [26,27]

\[
H = \left(\frac{p^2}{2m} + \varepsilon_+\right) \mathbf{1} - \varepsilon_\Delta \sigma_z \otimes \mathbf{1} + \left(\frac{\eta}{\hbar}\right) \sigma_x \otimes (p_x \sigma_y - p_y \sigma_x).
\]

(2)

Here \(\sigma_x, \sigma_y, \sigma_z\) denote the Pauli matrices describing the subband (or pseudospin) degree of freedom, and \(\sigma_{x,y,z}\) are Pauli matrices referring to the electron spin. The intersubband SO coupling \(\eta\) (which has the dimensionality of square of charge) is expressed [26] in terms of the gradients

\[^2\text{For the two-subband wide-well sample of ref. [23] the total electron density is } n = 9.2 \times 10^{11} \text{ cm}^{-2} \text{ and the Fermi wave vectors are } k_{F,1} = k_{F,2} = k_F = \sqrt{2m/\varepsilon_F} = 1.7 \times 10^6 \text{ cm}^{-1}. \text{ The Fermi energy is } \varepsilon_F = \hbar^2 k_F^2/2m = 16.4 \text{ meV} \text{ (assuming } m = 0.067m_0 \text{ for a GaAs well). Note that } \varepsilon_F \gg \Delta = 1.4 \text{ meV in ref. [29].}\]
of the Hartree-type contribution to the electron potential, the external gate and doping potentials, and the structural quantum-well potential profile. Note that the gap is \( \Delta = \varepsilon_0 - \varepsilon_c = 2\varepsilon_c - \varepsilon_c \).

**Theoretical approach.** – To calculate the edge spin density in the quasi-ballistic regime we follow the method developed in refs. [21,22] for the case of the single-subband Rashba Hamiltonian. Assuming that the spatial scale of the edge spin accumulation \( \Lambda \) is much smaller than the mean free path \( l \), we solve the edge spin problem by the method of scattering states, i.e., we find the exact quantum-mechanical solution of the electron scattering by an impenetrable straight boundary (fig. 1(a)) at a given Fermi energy. These solutions are then used in the calculation of the (mean) spin density profile. The populations of the incoming states are found from the solution of the kinetic equation for the spin-density matrix in the bulk (2D) of the sample in the presence of electric field, see SM.

The Hamiltonian (2) has 4 eigensolutions \( \Psi_{1,i}, \Psi_{1,\uparrow}, \Psi_{2,\uparrow}, \Psi_{2,i} \) (see SM for their explicit form) with the corresponding energy spectrum

\[
\varepsilon_{1,2}(p) = \frac{p_i^2}{2m} + \varepsilon_s + \sqrt{\varepsilon_c^2 + \eta^2 p^2 / \hbar^2},
\]

where the subscript \( i = 1,2 \) corresponds to the higher (lower) energy subband. Each subband is doubly degenerate with respect to the “spin direction” \( s = \uparrow, \downarrow \) (Kramers pairs). Upon scattering by the straight boundary where energy and momentum \( p_y \) along the boundary are conserved (and assuming zero rate of spin relaxation at the edge), the states in the pair \( \Psi_{1,\uparrow}(\varphi_1, \theta_1), \Psi_{2,\uparrow}(\varphi_2, \theta_2) \) mix up and form two scattering states, eqs. (4), (5) (similarly for the pair \( \Psi_{1,\downarrow}(\varphi_1, \theta_1), \Psi_{2,\downarrow}(\varphi_2, \theta_2) \)). For this pair of scattering states, we have the following boundary condition for the scattering by a hard wall located at \( x = 0 \) (fig. 1(a)):

\[
\begin{align*}
\tilde{\Psi}_{1,\uparrow}(x,y)_{\mid x=0} &= e^{ik_y y} \left[ \Psi_{1,\uparrow}(\pi - \varphi_1, \theta_1) e^{-ik_1 x} + F_{1,\uparrow}^{1,1} \Psi_{1,\uparrow}(\varphi_1, \theta_1) e^{ik_1 x} + F_{1,\downarrow}^{2,2} \Psi_{2,\downarrow}(\varphi_2, \theta_2) e^{ik_2 x} \right]_{\mid x=0} = 0, \\
\tilde{\Psi}_{2,\uparrow}(x,y)_{\mid x=0} &= e^{ik_y y} \left[ \Psi_{2,\downarrow}(\pi - \varphi_2, \theta_2) e^{-ik_2 x} + F_{2,\uparrow}^{1,1} \Psi_{1,\uparrow}(\varphi_1, \theta_1) e^{ik_1 x} + F_{2,\downarrow}^{2,2} \Psi_{2,\downarrow}(\varphi_2, \theta_2) e^{ik_2 x} \right]_{\mid x=0} = 0,
\end{align*}
\]

where \( p_1^2 = \hbar^2 (k_1^2 + k_2^2), p_2^2 = \hbar^2 (k_1^2 + k_2^2), \varepsilon_1(p_1) = \varepsilon_2(p_2) = \varepsilon_c. \) The momenta \( p_1, p_2 \) describe states belonging to subbands 1 and 2 for a given energy \( \varepsilon \), see fig. 1(b). The angles \( \varphi_1, \varphi_2 \) (between the corresponding momenta and the positive direction of the \( x \)-axis) are expressed as \( \sin(\varphi_1) = k_{1y} / p_1 \) and \( \sin(\varphi_2) = k_{2y} / p_2 \). The angles \( \theta_1, \theta_2 \) are defined via \( \cos(\theta_1, \theta_2) = 1/\sqrt{1 + 2p_{1,2}/\hbar^2} \). The expressions for the scattering amplitudes (\( F_{1,\uparrow}^{1,1}, \text{ etc.} \)) and the corresponding components of the unitary scattering matrix \( \hat{S} \) are presented in the SM. Similar equations can be written for the pair \( \Psi_{1,\uparrow}(\varphi_1), \Psi_{2,\uparrow}(\varphi_2) \), and the corresponding scattering matrix elements are also determined.

The expectation value of the \( z \)-component of the spin as a function of the coordinates is given by the following expression:

\[
\langle S_z(x) \rangle = \sum_{i,s} \int \frac{dk_y}{2\pi} \frac{d\varepsilon}{2\varepsilon_c} f_s(\varepsilon, k_y) \langle \hat{S}_{i,s} | \hat{S}_{i,s}(x) \rangle.
\]

Here \( f_s(\varepsilon, k_y) \) is the distribution function of the electron state in the subband \( i \) for a given energy and given wave vector \( k_y \) along the boundary and the group velocity \( v_{x,i} = \partial \varepsilon_i / \partial p_x \).

We can then calculate the most important part of the edge spin density which is smooth on the scale of the Fermi wavelength\(^a\) and involves the interference of the outgoing waves (two last terms in eqs. (4) and (5)). The corresponding formula for \( \langle S_z(x) \rangle \) valid for general values of the parameters (in the case in which both subbands are occupied) is presented in the SM. In the most important case \( p_1 - p_2 < p_F \), when the energy separation between two subbands \( \sqrt{\Delta^2 + 4\eta^2 k_F^2} \) is much smaller than the Fermi energy, the edge spin density for arbitrary values of the parameter \( \xi = L_\Delta / L_s \) reads (here we disregard the intra-band linear Dresselhaus term [25]):

\[
\langle S_z(x) \rangle = -\sin^2 \theta \int \frac{dk_y}{2\pi} \frac{d\varepsilon}{2\varepsilon_c} \sin \left( \frac{x}{\Lambda \sqrt{1 - (k_y/k_F)^2}} \right) \langle f_s(\varepsilon, k_y) - f_\downarrow(\varepsilon, k_y) \rangle.
\]

Here \( \varepsilon = p_F^2 / 2m \) is the Fermi energy. While deriving eq. (7), we used that \( \theta_1 - \theta_2 \ll \theta_{1,2} \), and \( \varphi_2 - \varphi_1 \ll \varphi_{1,2} \). The difference of the distribution functions entering eq. (7) is calculated in the SM assuming the set of inequalities \( k_1^2 < d \ll L_s \), where \( d \) is the correlation radius of the impurity potential in the bulk of the structure. The first condition means that the scattering in the bulk is of the small-angle type. Both conditions are fulfilled for a high-mobility GaAs structure. The final result derived from eq. (7) reads

\[
\langle S_z(x) \rangle = \frac{3k_F}{T_s} \Phi(\xi) J(x/\Lambda); \quad \Phi(\xi) = \frac{\xi}{(2\xi^2 + 1)^2 + 1}\frac{\xi}{(2\xi^2 + 1)^2}.
\]

with the spatial dependence given by the integral

\[
J(x/\Lambda) = \int_0^1 \frac{dz}{\pi z^2} \frac{x}{z(\lambda \sqrt{1 - z^2})}, \quad \Lambda = \frac{L_\Delta L_s}{\sqrt{L_\Delta^2 + L_s^2}}.
\]

We recall that \( \Lambda \) is the characteristic spatial scale of the edge spin density. For \( x \ll \Lambda \) we have \( J(x) \propto x/\Lambda \). In the opposite limit \( x \gg \Lambda \), we obtain \( J(x) \propto (\Lambda / x)^{3/2} \cos[(x/\Lambda) + \pi/4] \).

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\(^a\)A fast contribution to the edge spin density which oscillates as a function of \( x \) with \( 2k_F \) wave vector, gives parametrically smaller contribution to the total spin \( \int_0^\infty dx (S_z(x)) \), and we omit it [22].
Weak SO coupling: \( L_s \gg L_\Delta \). – To contrast our results with the usual one-band Rashba case it is instructive to consider the limit where a weak SO coupling \( 2\pi k_F \Delta \ll \Delta \). In this limit we can calculate the difference of the distribution functions entering eq. (7) using their standard expressions at \( \eta = 0 \) (see also SM), i.e.,

\[
f_{1,2} = (\varepsilon \varepsilon h k_0 / m) \tau_{\varepsilon} (p_{1,2}) \partial f_0 / \partial \varepsilon,
\]

where \( f_0 \) is the Fermi function, the electric field \( E \) is directed along the \( y \)-axis, and \( \tau_{\varepsilon} (p) \) is the momentum-dependent transport scattering time calculated within the Born approximation due to impurity scattering in the bulk. The values of \( p_{1,2} \) are related through \( \varepsilon_1 (p_{1}) = \varepsilon_2 (p_{2}) = \varepsilon = \varepsilon_F \), fig. 1(b). Using the condition \( k_F d > 1 \) (small-angle scattering in the bulk), we obtain \( \tau_{\varepsilon} (p_{1}) \approx 3 \varepsilon_1 (p_{1}) / 2 \approx 3 \varepsilon_2 (p_{2}) / 2F \approx (3 / k_F L_\Delta) \) (see footnote \( 4 \)). We note that compared to the usual Rashba one-band case the difference of the distribution functions considered here is finite at \( \eta = 0 \), and is of the first order in \( p_{1} - p_{2} = h / L_\Delta \). Since the SO coupling is weak, the probability of the spin flip is small which shows up as the small factor \( \sin^2 \theta \approx L_\Delta ^2 / L_s ^2 \ll 1 \) in eq. (7), and finally we obtain

\[
\langle S_z (x) \rangle = 3 k_F L_\Delta / L_s J (x / \Lambda),
\]

which coincides with the result which follows from eq. (8) in the limit \( \varepsilon \rightarrow 0 \).

The calculated edge spin density, eq. (8), is maximal at \( L_\Delta \approx L_s \) when it is of the order of \( k_F L_s \). By decreasing the gap \( (L_s < L_\Delta ) \) the spectrum approaches the usual Rashba model type (doubly degenerate), and because of the specific cancellation inherent in that model \( \langle S_z (x) \rangle \) decreases in magnitude as \( k_F L_s / L_\Delta ^2 \) (see fig. 2(a)), finally approaching the limit calculated in ref. [22] given by \( \approx (k_F L_s / L_\Delta ^2) (d^2 / L_s ^2) \) (see also eq. (1)). Thus, for a given strength of the SO interaction, the magnitude of the edge spin density has non-monotonic dependence as a function of the \( L_s \) (or \( \Delta \)), fig. 2(a). We note that if one takes for the ratio \( d / L_s = 0.1 \), then the edge spin density obtained in ref. [22] for the usual Rashba system with one subband equals in magnitude the density which follows from eq. (8) at \( \xi \approx 35 \), where the latter is three orders of magnitude smaller than its maximal value at \( \xi = 0.56 \).

Comparison with the experiment. – The experimental estimate of \( L_\Delta \) is \( \approx 1.4 \times 10^{-5} \) cm. For \( L_s \) we take two characteristic lengths \( 1.1 \) \( \mu \)m and \( 2.8 \) \( \mu \)m. Note that the corresponding values of \( \eta \) are consistent with the ones obtained from the theoretical calculations [30] for structures similar to that used in the experiment of ref. [23]. Thus, the value \( \xi = 0.1 \) will reasonably correspond the above chosen lengths. Calculating \( \tau_{\varepsilon} \) from the mobility \( 1.9 \times 10^{6} \) \( \)\( \mu \)s/\( \)cm\( \mu \)V, and using \( E = 0.05 \) mV/\( \mu \)m for the electric field in the quasi-ballistic region of the sample (both the mobility and \( E \) are taken from ref. [23]), we plot \( (S_z (x)) \), fig. 2(b). The exact experimental value of the edge spin density is not known; the authors of ref. [23] have estimated the threshold minimal value compatible with their observation as \( 3 \times 10^{6} \) \( \)\( \mu \)s\( \mu \)m\( \mu \)V. Hence this number is consistent with our calculation. In addition, we stress that the procedure just described, i.e., comparison of our theoretical predictions for the edge spin density with the experimental value of this quantity, allows one to extract the value of \( \eta \).

In conclusion, using a Rashba-like SO interaction arising from the coupling between two subband states of opposite parities in a symmetric two-subband quantum well, we have explained the great difference between the edge spin density in a bilayer structure as compared to the one in a single-layer configuration observed in the experiment of ref. [23]. The presence of the gap between the two subbands governs the effective strength of the inter-subband SO interaction and changes drastically the picture of the edge spin accumulation. Thus, by varying the gap value one can easily proceed from the regime of strong spin accumulation to the regime of weak spin accumulation. This opens up the possibility for the design of new spintronic devices.

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