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2007 J. Phys.: Condens. Matter 19 076205

(http://iopscience.iop.org/0953-8984/19/7/076205)

IOP Publishing — Journal of Physics: Condensed Matter
J. Phys.: Condens. Matter **19** (2007) 076205 (7pp) doi:10.1088/0953-8984/19/7/076205

The effect of magnetic field on a nonballistic spin field effect transistor

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Received 11 October 2006, in final form 7 December 2006
Published 2 February 2007
Online at stacks.iop.org/JPhysCM/19/076205

Abstract
A spin field effect transistor (FET) made of a nonballistic quantum wire with a single transport channel is considered in the presence of a magnetic field. The magnetic field includes either the externally applied field or the stray field due to ferromagnetic contacts used as injector and collector. When a magnetic field is applied the conductance fluctuations alter the spin precession and moreover spin flip occurs if the magnetic field is perpendicular to the Rashba field. Necessary conditions for a successful spin FET operation is obtained in the presence of a magnetic field.

1. Introduction

Spin-polarized electron transport has been a subject of persistent interest [1, 2]. When successfully combined with semiconductor functionalities [3–5] spin-based electronics or spintronics may have considerable impact on future electronic device applications. The goal of spintronics is to develop devices by employing the electron's spin degree of freedom. One of the representative spintronic devices is the so-called spin field effect transistor (spin FET) proposed by Datta and Das [6]. The core idea of this device is to induce spin precession by the Rashba spin–orbit interaction [7] in a two-dimensional electron gas (2DEG) and to use spin-dependent materials, such as ferromagnets, for electron injectors and collectors so that they sense the spin precession and the conductance of the device varies sinusoidally with the spin precession angle.

While the initial proposal of Datta and Das assumes ballistic transport, in the absence of magnetic field, it is still essential [8–13] to have a good understanding of how sensitive the spin FET is to scattering events, and to external magnetic field. For example a sample may not be as ideal as desired, and unintended impurities in the 2DEG may cause elastic scattering. The scattering may also be caused by tunnelling barriers [14] introduced at the 2DEG–injector (collector) interface to enhance the spin injection (detection) efficiency. In addition the spin

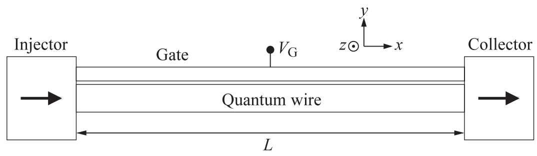

Figure 1. Schematic diagram of a Datta–Das spin field effect transistor made of a nonballistic single-channel quantum wire.

FET can be very sensitive to the magnetic field which may be intentionally applied or arise from the stray field induced by the ferromagnetic electrodes [15].

Elastic scattering and external magnetic field can induce spin relaxation when combined with spin–orbit interaction (see for instance [16]). Semiclassical Monte Carlo calculations showed that the spin relaxation can be suppressed [8, 9, 17, 18] and the ideal sinusoidal variation of the spin FET signal can be achieved. As the channel width w of the 2DEG reduces, quantum mechanical effects become more important, which may go beyond the semiclassical treatment. Here we address the effect of impurity scattering in the vanishing width ($w \to 0$) limit by using a full quantum mechanical analysis. When $w \ll \hbar^2/m^*\alpha$ (α is the Rashba coefficient and m^* is the effective mass) the intersubband mixing between quantized subbands (each of which provides a transport channel), emerging as a result of the quantization in the transverse direction, can be neglected [6]. Moreover the number of available channels reduces to two, including the spin degree of freedom, for sufficiently small w. It was suggested [19] that a spin FET with only one transverse mode is desired to achieve large current modulation and low power consumption, and it is found that it exhibits interesting mesoscopic phenomena [20].

In [21] we showed the necessary condition for a successful nonballistic spin FET operation in the absence of magnetic field. However, for realistic applications it is also necessary to treat the influence of the magnetic field which may arise as a result of ferromagnetic metals being used as the injector and collector [15]. Hence the present work will be an extension of our previous paper [21], namely we investigate a spin FET made of a nonballistic quantum wire ($w \to 0$) with one transport channel, in the presence of a magnetic field (figure 1). Although the magnetic field and the elastic scattering obviously affect the signal of the spin FET, what is important is the issue of its sinusoidal modulation by the Rashba interaction. We address this issue by employing a full quantum mechanical analysis (in the single-particle level) [22, 23] of nonmagnetic scattering effects on such a one-dimensional (1D) spin FET. We find that with a magnetic field parallel to the Rashba field, even though there is no spin relaxation, the spin precession shows sample to sample fluctuations and that with the magnetic field perpendicular to the Rashba field spin flip occurs. We found the necessary conditions for successful operation of a nonballistic 1D spin FET influenced by an external magnetic field.

2. Zero magnetic field

For a nonballistic 1D spin FET in the absence of external magnetic field, the scattering effects within the two spin channels can be expressed by the following effective mass single-particle Hamiltonian:

$$H = \frac{p_x^2}{2m^*} + V(x) + \alpha\sigma_z \frac{p_x}{\hbar} \tag{1}$$

where the first term is the kinetic energy, $V(x)$ is the nonmagnetic scattering potential causing spin-conserved scattering (for $w \ll \hbar^2/m^*\alpha$), and the last term represents the Rashba interaction. α is the standard parameter measuring the strength of the Rashba interaction. Its value can be controlled [24] by the gate electrode in figure 1. The Rashba term in the

Hamiltonian is formally the same as the Zeeman term $-g\mu_B \boldsymbol{\sigma} \cdot \boldsymbol{B}_R$ with an effective magnetic field $\boldsymbol{B}_R = -\hat{z}(\alpha/g\mu_B\hbar)p_x$ which, unlike a real field, depends on the electron momentum. In order to focus on the scattering effects within the quantum wire of length L, it is assumed that $V(x)$ is nonzero only in $0 \leqslant x \leqslant L$. Hence the segment with length L describes the nonballistic quantum wire while the regions $x < 0$ and $x > L$ correspond to (fictitious) leads which are free from any scattering (see figure 1). Though the introduction of them is rather arbitrary, scattering effects within the quantum wire are still correctly described by this approach [25]. It is also supposed that the injector and the detector are ideal, both of which are 100% spin polarized, and that the injection–detection efficiency is perfect[1].

Following the scattering matrix approach [25], consider a left-incoming scattering state $\psi = c_+\psi_+ + c_-\psi_-$, a superposition of the spin-up scattering state $\psi_+ = \phi_+\chi_+$ and the spin-down scattering state $\psi_- = \phi_-\chi_-$. Here $\phi_{+(-)} = \exp[\mathrm{i}k_{+(-)}x] + r_{+(-)}\exp[-\mathrm{i}k_{-(+)}x]$ for $x < 0$ and $\phi_{+(-)} = t_{+(-)}\exp[\mathrm{i}k_{+(-)}x]$ for $x > L$, where $E = \hbar^2k^2_{+,-}/2m^* \pm \alpha k_{+,-}$. The spinors are described by $\chi_+ = (1, 0)^{\mathrm{T}}$ and $\chi_- = (0, 1)^{\mathrm{T}}$ and the conservation of σ_z, $[H, \sigma_z] = 0$, is employed. Let θ_0 and φ_0 (θ_L and φ_L) denote the polar and azimuthal angles of the spin direction at $x = 0_\leqslant$ ($x = L_\geqslant$) with respect to the z axis. Just before the injection into the quantum wire, $x = 0_\leqslant$, $\tan(\theta_0/2) = |c_-/c_+|$ and $\exp[\mathrm{i}\varphi_0] = (c_-/c_+)/|c_-/c_+|$. Right after the transmission, $x = L_\geqslant$,

$$\tan\frac{\theta_L}{2} = \left|\frac{c_-}{c_+}\frac{t_-}{t_+}\right|, \qquad \mathrm{e}^{\mathrm{i}\varphi_L} = \frac{(c_-/c_+)}{|c_-/c_+|}\frac{(t_-/t_+)}{|t_-/t_+|}\mathrm{e}^{\mathrm{i}(k_- - k_+)L}. \tag{2}$$

For the sake of simplicity it is supposed that injected electrons are polarized along the $+\hat{x}$ direction (figure 1) and that the ratio c_-/c_+ is taken to be unity. In the equation above the crucial term is the ratio t_-/t_+ determining the effects of the scattering on the spin precession. For ballistic case $t_- = t_+ = 1$, thus $\theta_{\mathrm{prec}} \equiv \theta_L - \theta_0 = 0$ and $\varphi_{\mathrm{prec}} \equiv \varphi_L - \varphi_0 = 2m^*\alpha L/\hbar^2$. Then as α varies the conductance, $G \propto \cos^2(\varphi_{\mathrm{prec}}/2)$, of the ballistic spin FET exhibits a sinusoidal variation.

To address nonballistic situations (mean free path $l \leqslant L$) the following gauge transformation[2] is utilized,

$$\tilde{\psi} = \mathrm{e}^{\mathrm{i}(m^*\alpha x/\hbar^2)\sigma_z}\,\psi. \tag{3}$$

In the absence of magnetic field, for a nonballistic spin FET in [21] it has been obtained that the spin precession was inert to the scattering and that for the successful sinusoidal variation of the spin FET signal a necessary condition, (7), should be fulfilled.

3. Effect of magnetic field

The external magnetic field $\boldsymbol{B}_{\mathrm{ext}}$ may be intentionally applied or it may arise from a stray magnetic field [15] of ferromagnetic metals used as the electron injector/collector. We will examine two cases: (i) when $\boldsymbol{B}_{\mathrm{ext}}$ is parallel to $\boldsymbol{B}_R$ which is oriented along the $+\hat{z}$ direction, $\boldsymbol{B}_{\mathrm{ext}} \parallel \boldsymbol{B}_R$, and next (ii) when $\boldsymbol{B}_{\mathrm{ext}}$ is perpendicular to $\boldsymbol{B}_R$, $\boldsymbol{B}_{\mathrm{ext}} \perp \boldsymbol{B}_R$.

3.1. Parallel magnetic field

Consider the parallel case first and take $\boldsymbol{B}_{\mathrm{ext}} = B_\parallel\hat{z}$. The Hamiltonian becomes

$$H = \frac{p_x^2}{2m^*} + V(x) + \alpha\sigma_z\frac{p_x}{\hbar} - g\mu_B\sigma_z B_\parallel \tag{4}$$

[1] We also ignore Fabry–Perot-type coherent multiple scattering effects between the injector and the detector discussed by Mireles and Kirzcenow [4, 26].

[2] A similar transformation has been exploited in [12].

which commutes with σ_z and there are no fluctuations of the spin quantization axis. Without a magnetic field the spin precession is inert to scattering [21]; with $B_\parallel$, however, it is not inert to scattering any more, as demonstrated below. For systematic analysis, we can perform the gauge transformation in (3). Upon the transformation, the Schrödinger equations for $\tilde{\psi}_+$ and $\tilde{\psi}_-$ become $\tilde{h}\tilde{\psi}_{+,-} = \tilde{E}_{+,-}\tilde{\psi}_{+,-}$, where $\tilde{h} = p_x^2/2m^* + V(x)$ and $\tilde{E}_{+,-} = \tilde{E} \pm g\mu_B B_\parallel$ with $\tilde{E} = E + m^*\alpha^2/2\hbar^2$. Note that the spin-up and spin-down components now obey a spinless problem [12] with different effective energy. Taking into account this difference, the transmission amplitudes and wavenumbers before and after the gauge transformation are related as follows:

$$t_{+,-}(E) = \tilde{t}(\tilde{E}_{+,-}) \tag{5a}$$
$$k_{+,-}(E) = \tilde{k}(\tilde{E}_{+,-}) \mp m^*\alpha/\hbar^2. \tag{5b}$$

Since $\tilde{h}$ is independent of α and $B_\parallel$, information on spin precession can be extracted from (5a) and (5b). For the ballistic case, $\tilde{t} = 1$ regardless of energy, and thus $\theta_{\rm prec} = 0$ from (2). For $\varphi_{\rm prec} = 0$, taking into account (5b) results in $\varphi_{\rm prec} = 2m^*\alpha L/\hbar^2 + \delta\varphi_{\rm prec}^{B_\parallel}$, where $\delta\varphi_{\rm prec}^{B_\parallel} \equiv [(2m^*\tilde{E}_-/\hbar^2)^{1/2} - (2m^*\tilde{E}_+/\hbar^2)^{1/2}]L$ represents the additional spin precession caused by $B_\parallel$. For usual situations with $g\mu_B B_\parallel, m^*\alpha^2/2\hbar^2 \ll E$ its α-dependent part is smaller than $2m^*\alpha L/\hbar^2$ by a factor $(-1/4)(g\mu_B B_\parallel/E)(m^*\alpha^2/2E\hbar^2)^{1/2} \ll 1$ and thus $\delta\varphi_{\rm prec}^{B_\parallel}$ is essentially independent of α.

For the nonballistic case, on the other hand, $\tilde{t}$ depends on energy. Then the ratio $t_-(E)/t_+(E) = \tilde{t}(\tilde{E}_-)/\tilde{t}(\tilde{E}_+)$ deviates from unity and depends on details of $V(x)$. Thus with nonzero $B_\parallel$, the spin precession angles, (2), depends on $V(x)$, i.e., the spin precession angles are not inert to scattering any more. Note that since $|t_+(E)/t_-(E)| \neq 1$, not only $\varphi_{\rm prec}$ but also $\theta_{\rm prec}$ is affected by scattering. In comparison, in semiclassical treatments that ignore the coherence in the orbital part, $\theta_{\rm prec}$ should be strictly zero regardless of $V(x)$ since the total effective magnetic field $\boldsymbol{B}_R + B_\parallel\hat{z}$ is always along the z-axis. This result of the nonzero $\theta_{\rm prec}$ thus goes beyond the semiclassical treatments and indicates the importance of equal-footing treatment of orbital and spin parts.

To estimate $\theta_{\rm prec}$, one can use the energy scale (Thouless correlation energy [27]) $E_c \sim (\hbar v_F/L)(l/L)$, where v_F is the Fermi velocity[3], over which $|\tilde{t}|$ is correlated [25]. Then the deviation of $|t_+(E)/t_-(E)|$ from unity depends on the ratio $g\mu_B B_\parallel/E_c$. For $g\mu_B B_\parallel \ll E_c$, $\theta_{\rm prec} = c_1 g\mu_B B_\parallel/E_c$ (in radians), where c_1 is of order one. Its precise value shows sample-to-sample fluctuations and, for a given sample, depends on $E + m^*\alpha^2/2\hbar^2$, with the correlation energy scale given again by E_c. For $g\mu_B B_\parallel \gg E_c$, on the other hand, $\theta_{\rm prec} = c_2$ (radians), where c_2 is of order one. The value of c_2 again shows sample-to-sample fluctuations and, within a single sample, depends on $E + m^*\alpha^2/2\hbar^2$ and $g\mu_B B_\parallel$, with the correlation scale given by E_c.

We then address the variation of the conductance G as a function of α. For the ballistic case, G shows the ideal sinusoidal variation with full magnitude, though $\varphi_{\rm prec}$ is shifted by the essentially α-independent contribution $\delta\varphi_{\rm prec}^{B_\parallel}$. For the nonballistic case, the sinusoidal variation of G requires both $t_-(E_F) = \tilde{t}(E_F + m^*\alpha^2/2\hbar^2 - g\mu_B B_\parallel)$ and $t_+(E_F) = \tilde{t}(E_F + m^*\alpha^2/2\hbar^2 + g\mu_B B_\parallel)$ to be essentially independent of α in the range $\alpha_{\rm av} - \Delta\alpha/2 < \alpha < \alpha_{\rm av} + \Delta\alpha/2$, which is possible only when

$$\Delta\left(\frac{m^*\alpha^2}{2\hbar^2}\right) + 2g\mu_B B_\parallel = \frac{m^*\alpha_{\rm av}\Delta\alpha}{\hbar^2} + 2g\mu_B B_\parallel \ll E_c. \tag{6}$$

[3] To be precise, v_F is the Fermi velocity in the gauge-transformed system. However, when v_F is sufficiently larger than $\alpha/\hbar$, which is usually valid, the Fermi velocity in the original system is again comparable to v_F and we may not distinguish the two.

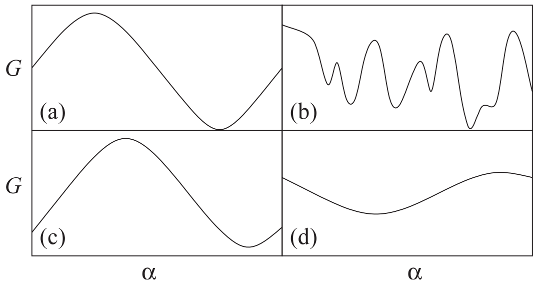

Figure 2. Schematic plots of the conductance G, varying between zero and one in units of e^2/h, of a 1D Datta–Das spin FET as a function of the Rashba coefficient α. (a) Ideal sinusoidal variation. (b) Sample-specific random signal due to conductance fluctuations. (c) Phase shifted by $B_\parallel$. (d) Phase shifted by a sample-specific random number and amplitude suppressed by a sample-specific random factor.

In the presence of $B_\parallel$, (6) is the generalization of the condition, below, found in [21] without taking into account the magnetic field:

$$\Delta\left(\frac{m^*\alpha^2}{2\hbar^2}\right) = \frac{m^*\alpha_{\mathrm{av}}\Delta\alpha}{\hbar^2} \ll E_{\mathrm{c}}. \tag{7}$$

For the case $\Delta\alpha = \pi\hbar^2/m^*L$, the value for $\Delta\varphi_{\mathrm{prec}} = 2\pi$ in the ballistic case, (7) reduces to $L \ll l(\hbar v_{\mathrm{F}}/\pi\alpha_{\mathrm{av}})$, which can be satisfied simultaneously with the nonballisticity condition $l \leqslant L$ since $\alpha_{\mathrm{av}}/\hbar$ is usually smaller than v_{F}.[4] Figures 2(a) and (b) show schematic plots of G as a function of α in two situations: (a) when (7) is satisfied and (b) when it is severely violated. Note that, for given $\Delta\alpha$, smaller α_{av} is preferred by (7) for minimal damage by the conductance fluctuations and that the amplitude of the G modulation is subject to sample-to-sample fluctuations even when (7) is satisfied, since $|t_{+,-}|^2$ depends on details of the nonballistic samples.

$B_\parallel$ makes (7) more restrictive. When (6) is satisfied, θ_{prec} remains close to zero and φ_{prec} grows linearly with α, $\varphi_{\mathrm{prec}} \approx 2m^*\alpha L/\hbar^2 + \delta\varphi_{\mathrm{prec}}^{B_\parallel}$, similar to the ballistic case. Figure 2(c) shows a schematic plot of G as a function of α. Note the phase shift by $\delta\varphi_{\mathrm{prec}}^{B_\parallel}$, compared to 2(a). An insight may be gained by studying properties of the spin FET signal in situations where (6) is not satisfied. We consider three such situations. First, when $\Delta(m^*\alpha^2/2\hbar^2) \gg E_{\mathrm{c}}$ but $2g\mu_{\mathrm{B}}B_\parallel \ll E_{\mathrm{c}}$, the situation is essentially the same as the case without $B_\parallel$ [21] and G as a function of α shows sample-specific random variation (figure 2(b)) due to the random fluctuation of $|t_+(E_{\mathrm{F}})|^2 \approx |t_-(E_{\mathrm{F}})|^2$. Second, when $\Delta(m^*\alpha^2/2\hbar^2) \ll E_{\mathrm{c}}$ but $2g\mu_{\mathrm{B}}B_\parallel \gg E_{\mathrm{c}}$, the two transmission amplitudes $t_+(E_{\mathrm{F}})$ and $t_-(E_{\mathrm{F}})$ are not correlated, though both of them are essentially independent of α. The absence of the correlation modifies both φ_{prec} and θ_{prec}; φ_{prec} acquires an additional contribution $\delta\varphi_{\mathrm{prec}}^{\mathrm{NB}}$ to become $2m^*\alpha L/\hbar^2 + \delta\varphi_{\mathrm{prec}}^{B_\parallel} + \delta\varphi_{\mathrm{prec}}^{\mathrm{NB}}$, where $\delta\varphi_{\mathrm{prec}}^{\mathrm{NB}} = \arg[t_-(E_{\mathrm{F}})/t_+(E_{\mathrm{F}})]$ is an essentially α-independent but sample-specific random constant of order one (in radians). Similarly, θ_{prec} acquires a nonzero contribution, via (2), which is an essentially α-independent but sample-specific random constant of order one (in radians). In this situation, the spin FET still shows the sinusoidal modulation (see figure 2(d)), though the modulation phase is shifted by a sample-specific random number due to the random shift of φ_{prec} by $\delta\varphi_{\mathrm{prec}}^{\mathrm{NB}}$, and the modulation amplitude is suppressed by the sample-specific random factor $\cos\theta_{\mathrm{prec}}$ compared to the case without $B_\parallel$ [21]. Third, when $\Delta(m^*\alpha^2/2\hbar^2) \gg E_{\mathrm{c}}$ and $2g\mu_{\mathrm{B}}B_\parallel \gg E_{\mathrm{c}}$, the amplitudes $t_+(E_{\mathrm{F}})$ and $t_-(E_{\mathrm{F}})$ are uncorrelated, and moreover vary

[4] For example, $\alpha_{\mathrm{av}}/\hbar v_{\mathrm{F}} \sim 0.1$ for $In_{1-x}Al_xAs/In_{1-x}Ga_x$.

independently with α. Then both θ_{prec} and φ_{prec} show sample-specific random fluctuations of order one (in radians) as α varies, and the ideal sinusoidal modulation of the spin FET signal is replaced by sample-specific random fluctuations; see figure 2(b).

3.2. Perpendicular magnetic field

When the external magnetic field is perpendicular to the Rashba field, say $\boldsymbol{B}_{\mathrm{ext}} = B_{\perp}\hat{x}$, the Schrödinger equation becomes

$$\left(\frac{p_x^2}{2m^*} + V(x) + \alpha\sigma_z\frac{p_x}{\hbar} - g\mu_{\mathrm{B}}B_{\perp}\sigma_x\right)\psi = \mathrm{i}\hbar\partial_t\psi. \tag{8}$$

Unlike in the parallel magnetic field case, σ_z is not conserved any more, $[H, \sigma_z] \neq 0$, since $B_{\perp}$ flips the spin and leads to the spin relaxation in the nonballistic environment.

For successful operation of the spin FET, the spin flipping probability should be small. Below we derive the condition for negligible spin flip. First we perform the gauge transformation

$$\tilde{\psi} = \mathrm{e}^{\mathrm{i}(g\mu_{\mathrm{B}}B_{\perp}t/\hbar)\sigma_x}\psi. \tag{9}$$

In terms of the new wavefunction $\tilde{\psi}$, the Schrödinger equation takes the form

$$\left(\frac{p_x^2}{2m^*} + V(x) + \alpha\sigma\cdot\hat{n}(t)\frac{p_x}{\hbar}\right)\tilde{\psi} = \mathrm{i}\hbar\frac{\partial}{\partial t}\tilde{\psi} \tag{10}$$

where $\hat{n}(t) = \hat{z}\cos(2g\mu_{\mathrm{B}}B_{\perp}t/\hbar) - \hat{y}\sin(2g\mu_{\mathrm{B}}B_{\perp}t/\hbar)$. Note that the transformation removes the Zeeman term at the expense of making the direction $\hat{n}$ of the Rashba field time-dependent. Note also that the absence of the spin flipping in the original gauge (ψ) is equivalent to the adiabatic evolution of the spin in the transformed gauge ($\tilde{\psi}$). Such adiabatic spin evolution is possible [28] if the spin precession angle $2m^*\alpha L/\hbar^2$ by the Rashba interaction alone is much larger than that by $B_{\perp}$ alone. For the α-modulation in the range $\alpha_{\mathrm{av}} - \Delta\alpha/2 < \alpha < \alpha_{\mathrm{av}} + \Delta\alpha/2$ with $\Delta\alpha = \pi\hbar^2/m^*L$, this consideration results in the condition

$$g\mu_{\mathrm{B}}B_{\perp} \ll E_{\mathrm{c}}\frac{2m^*\alpha_{\mathrm{av}}L}{\hbar^2}. \tag{11}$$

(11), together with the constraint for weak conductance fluctuations, (7), identifies the regime for successful spin FET operation in the presence of $B_{\perp}$. An applied magnetic field perpendicular to the 2DEG was shown to have little effect on the spin precession length [29]. Note that for given $B_{\perp}$, (11) prefers a large α_{av}, while smaller α_{av} is preferred by (7) for the minimal damage by the conductance fluctuations. Thus a careful tuning is necessary.

In the spin FET configuration originally proposed by Datta and Das, $B_{\perp}$ may have a nonvanishing value even when no field is applied externally, since the stray field generated by the injector and collector is perpendicular to $\boldsymbol{B}_{\mathrm{R}}$ and can be as high as 1 T [15]. Thus careful adjustment may be necessary to satisfy (11). One possible adjustment method is to apply an external magnetic field, which is antiparallel to the stray field, so that the total $B_{\perp}$ satisfies (11). Of course the strength of this counter field should be weaker than the coercive field of the injector and collector. Otherwise it flips the magnetization direction of the injector and collector, and the stray field becomes parallel to the external field.

4. Conclusion

In summary, we have studied a 1D Datta–Das spin FET made of a nonballistic quantum wire and identified the necessary conditions, (6), (7) and (11), for the sinusoidal modulation of the spin FET signal in nonballistic environments. We suggest that the effects of impurity scattering and magnetic field, which can be very harmful for the spin FET operation, can be avoided by tuning the parameters of a spin FET to satisfy the inequalities (6), (7) and (11). Using the obtained values in a recent experimental study on InGaAs/InP quantum wires [30], $[(m^*\alpha_{av}\Delta\alpha)/\hbar^2)]/E_c \sim 0.01$ and $(g\mu_B B_{ext})/[E_c(2m^*\alpha_{av}L/\hbar^2)] \sim 10^{-5}$, confirming that our conditions can be fulfilled.

Acknowledgments

We are greatly indebted to Professor H Woo Lee for stimulating discussions and comments. We thank Pohang University of Science and Technology, where most of this work was performed.

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