Application of split-gate structures as tunable spin filters

M. J. Gilbert and J. P. Bird[a)]
Center for Solid State Electronics Research & Department of Electrical Engineering, Arizona State University, Tempe, Arizona 85287-5706

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We describe an electron filter that exploits the known magnetotransport properties of quantum point contacts to provide *local* and *tunable* control of the spin polarization in a semiconductor. When properly configured, we show that the conductance of this device gives a direct measure of the spin polarization of carriers transmitted through it. By modeling the transport through a potential barrier with experimentally realistic parameters, we discuss the factors which must be satisfied in order to successfully implement such a device.
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Exploitation of the spin degree of freedom offers much potential for application in new electronic devices, and is an area of research interest that is currently attracting increasing attention. Motivated by recent experiments,[1,2] which have demonstrated the existence of long relaxation times for spin-polarized distributions of carriers, a number of theoretical reports have explored the possibility of exploiting the spin states of the electron as the basis of a quantum logic scheme.[3,4] For applications such as quantum computing, an *electrical* means of generating spin-polarized distributions of carriers is required.[5,6] Furthermore, it is desirable to be able to induce spin polarization at *local* points within a circuit, and to *vary* the degree of the polarization *in situ*. In this letter, we describe the design of such a tunable *spin filter*, that exploits the known magnetotransport properties of quantum point contacts. We also explore the different requirements that must be satisfied in order to successfully realize such a device.

The device we consider is based upon the split-gate quantum point contact, which is known from experiment to provide a spin-dependent barrier to electron motion in the presence of an external magnetic field.[7,8] A schematic illustration of the spin filter is shown in Fig. 1, in which the lower set of gates is used to define a quantum point contact in the high-mobility electron gas of a suitable heterostructure. The split gates are separated from an upper *continuous* gate by the inclusion of a thin insulating layer, which may be formed by oxide deposition or by hardening of electron-beam resist.[9] In basic operation, the lower gates are configured to set the quantum point contact close to its threshold for conduction. A current is then driven through the upper gate, inducing a local magnetic field that lifts the spin degeneracy of the carriers and causes preferential transmission of one spin species through the quantum-point-contact barrier. By measuring the conductance of the point contact as the current in the upper gate is varied, a direct measure of the spin polarization of the transmitted carriers can then be obtained.

The self-consistent potential profile of a quantum point contact exhibits a saddle-like form, which, in the region near the saddle minimum, can be well approximated as[10,11]

$$V(x,y)=V_o-\tfrac{1}{2}m^*\omega_x^2x^2+\tfrac{1}{2}m^*\omega_y^2y^2. \quad (1)$$

In this equation, x is the distance measured *along* the axis of the quantum point contact, while y is the *transverse* displacement with respect to its center (see Fig. 1). V_o is the height of the saddle barrier, m^* is the electron effective mass, and ω_x and ω_y are characteristic oscillator frequencies. ω_y determines the energy splitting of the one-dimensional subbands in the quantum point contact, while ω_x dictates how sharply its transmission drops to zero when the Fermi level falls below the saddle minimum. With a magnetic field applied perpendicular to the electron gas, the transmission probability of the different subbands may be written as[11,12]

$$T_n=\frac{1}{1+\exp(-\pi\varepsilon_n)}. \quad (2)$$

Here, n is the usual harmonic-oscillator index ($n=0,1,2,3,...$) and ε_n is defined as

$$\varepsilon_n=\frac{E-E_2(n+1/2)-V_o}{E_1}, \quad (3)$$

where E is the energy of the incoming carriers and the parameters E_1 and E_2 are defined via

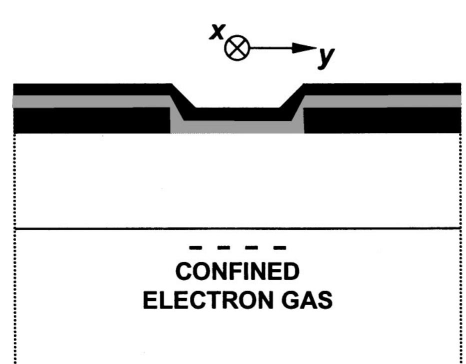

FIG. 1. Schematic illustration of a possible realization of the spin filter. The device is presumed to be fabricated on top of a suitable heterostructure and the black regions correspond to metallic gates, while the gray shading denotes a thin insulating layer. The coordinate system is also indicated for reference.

[a)]Electronic mail: bird@asu.edu

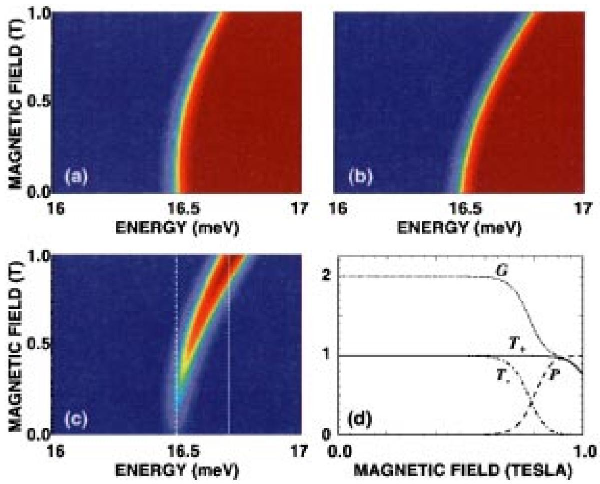

FIG. 2. (Color) (a) T_+ vs energy and magnetic field. (b) T_- vs energy and magnetic field. (c) $T_+ - T_-$ vs energy and magnetic field. (d) T_+, T_-, P, and G (in units of e^2/h) vs magnetic field, for an initial energy of $E = 16.7$ meV. In all of these calculations we have assumed $V_0 = 15$ meV, $\hbar\omega_y = 3$ meV, and $\hbar\omega_x = 0.1$ meV. In the contour plots, the color scale ranges from blue to red, indicating a variation of transmission between zero and one.

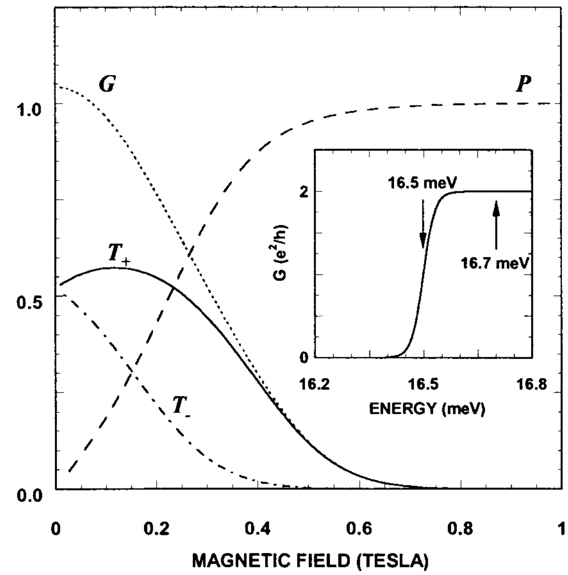

FIG. 3. Main panel: T_+, T_-, P, and G vs magnetic field, for an initial energy of 16.5 meV. We have once again assumed $V_0 = 15$ meV, $\hbar\omega_y = 3$ meV, and $\hbar\omega_x = 0.1$ meV. Inset: the variation of G (in units of e^2/h) as a function of energy at zero magnetic field, calculated using the Landauer formula and the results of Eqs. (1)–(5).

$$E_1 = \frac{\hbar}{2\sqrt{2}}([\Omega^4 + 4\omega_x^2\omega_y^2]^{0.5} - \Omega^2)^{0.5}, \tag{4}$$

$$E_2 = \frac{\hbar}{\sqrt{2}}([\Omega^4 + 4\omega_x^2\omega_y^2]^{0.5} + \Omega^2)^{0.5}. \tag{5}$$

In these latter equations, Ω is a magnetic-field (B) dependent oscillator frequency, $\Omega^2 = \omega_c^2 - \omega_x^2 + \omega_y^2$, and the cyclotron frequency $\omega_c = eB/m^*$. By introducing Ω in this manner, it is possible to describe the transmission through the saddle barrier *continuously*, as one transitions from zero magnetic field to the high-field quantum-Hall regime. Typically, however, we will be interested in magnetic field ranges for which the electrostatic confinement in the quantum wire remains at least comparable to the magnetic force (i.e., $\hbar\omega_c \leq \hbar\omega_y$).

The spin filter we discuss exploits the spin-dependent tunneling probabilities that exist for a potential barrier in the presence of a magnetic field. This magnetic field induces a Zeeman splitting $\pm\frac{1}{2}g\mu_B B$ of the electron energy, where g is the effective g factor, and μ_B is the Bohr magneton. For an electron with energy E at zero magnetic field, we may use Eqs. (2)–(5) to define *separate* transmission probabilities for its spin components in the presence of a magnetic field

$$T_\pm = T(E \pm \tfrac{1}{2}g\mu_B B). \tag{6}$$

To illustrate this concept, we model the response of a device whose barrier parameters are chosen to correspond to those of a typical point contact biased close to pinch off.[13] Since the effective g factor can exhibit a magnetic field dependence, and is also a materials-dependent parameter, for ease of implementation here we initially assume the free-electron value ($g=2$). (Further below, however, we discuss how the choice of substrate materials with large g factors can improve the operation of the spin filter.) In Figs. 2(a) and 2(b), we plot the variation of the spin-dependent transmission probabilities, T_+ and T_-, as a function of energy and magnetic field. The energy range considered here is chosen to correspond to that over which the conductance of the point contact is dominated by the first subband ($n=0$) alone. In Fig. 2(c), $T_+ - T_-$ is plotted for the same energy range and the bright region indicates the range of parameter space where the transmission of one spin species is high while that of the other is low. Starting at an initial energy of 16.7 meV [dotted line in Fig. 2(c)], for example, strong spin polarization of the transmitted carriers can be obtained by generating a magnetic field of 0.9 T. This is shown more clearly in Fig. 2(d), where we plot T_+ and T_- as a function of magnetic field, for this initial energy. We also show the variation of the spin polarization, which we define as $P = (T_+ - T_-)/(T_+ + T_-)$. The conductance of the barrier may be determined using the Landauer formula, $G = (e^2/h)(T_+ + T_-)$, and provides a direct measure of the spin polarization. Note how in Fig. 2(d) the conductance initially equals $2e^2/h$, since the lowest subband is spin degenerate and is fully transmitted at zero magnetic field (Fig. 3, inset). As the magnetic field is increased, however, the conductance decreases to e^2/h as the spin polarization reaches unity.

Proper configuration of the quantum point contact is important for operation of the spin filter. In Fig. 2, for example, by setting the initial energy such that the lowest subband was fully transmitted at zero magnetic field, we obtained a clear indication of the spin polarization in the conductance. In Fig. 3, however, we show the behavior obtained for an initial energy of 16.5 meV [dash-dotted line in Fig. 2(c)], for which value the conductance is less than $2e^2/h$ at zero magnetic field (Fig. 3, inset). In this case, the polarization rises to unity at a significantly lower magnetic field than that obtained in Fig. 2, but there is now no obvious indication in the conductance that polarization has been achieved. Furthermore, the polarization is induced at the expense of the transmission of *both* spin species. For successful application of this device, we not only require a high degree of polarization, but we *also* desire that the transmission of the preferred spin species be close to unity. (In actual experiments, of course, the initial conductance will be configured by using the gate voltage to adjust the height of the saddle barrier relative to the *fixed* Fermi level in the reservoirs.)

The preceding analysis suggests that locally induced magnetic fields of order a Tesla are required to generate a sufficiently strong energy splitting of the two spin components. For a typical separation of 100 nm between the upper gate and the electron gas, this requires a drive current of roughly 500 mA if the gate is modeled as a microstrip 50 nm wide.[14] While this current can be lowered by reducing the separation between the electron gas and the upper gate, a more fruitful approach should be to exploit heterostructures in which the g factor is strongly enhanced above the free electron value. In InAs, for example, the value of the g factor is -15, while in InSb it is as large as -51;[15] using heterostructures based on these materials,[16,17] it should be possible to achieve spin polarization with magnetic fields as small as a few mTesla. By fabricating the upper gate from a ferromagnetic material, such magnetic fields can easily be generated by using short current pulses to magnetize and demagnetize this gate. This approach obviates the need for a continuous drive current in the upper gate, and so should help in reducing power dissipation. High temperature operation of the device is more problematic, however. For the highest magnetic fields shown in Figs. 2 and 3, the magnitude of the spin splitting is not much more than a degree Kelvin. Using materials with large g factors, however, it should be possible to increase this value by more than an order of magnitude. Nonetheless, we view the immediate usefulness of this device as a means for generating spin-polarized carriers in fundamental spin-transport studies. By configuring pairs of such spin filters in series with each other, for example, it should be possible to study the basic details of spin relaxation in semiconductors.

In conclusion, we have described a tunable spin-filter that exploits the well-known transmission properties of quantum point contacts in a magnetic field. This device should allow for local control of spin polarization in an electrical circuit and may prove useful for a basic demonstration of spin-based quantum computing, where an important requirement will be the ability to prepare nanostructures in a well-defined spin state.

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