Optoelectric spin injection in semiconductor heterostructures without ferromagnet

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We have shown that electron spin density can be generated by a dc current flowing across a $pn$ junction with an embedded asymmetric quantum well. Spin polarization is created in the quantum well by radiative electron-hole recombination when the conduction electron momentum distribution is shifted with respect to the momentum distribution of holes in the spin split valence subbands. Spin current appears when the spin polarization is injected from the quantum well into the $n$-doped region of the $pn$ junction. The accompanied emission of circularly polarized light from the quantum well can serve as a spin polarization detector.

72.25.-b, 73.40.Kp, 78.60.Fi

One of the most important problems in spintronics is the efficient injection of spin currents into semiconductor structures. A possible way is to inject spin polarization from a ferromagnet by passing a dc current across the interface [1]. If the ferromagnet is metallic, the efficiency of such injection is controversial [2], and the observed weak spin current has been attributed to the large mismatch of spin diffusion constants between the adjacent semiconductor and metal [3]. Spin injection can be enhanced by using an appropriate interface tunneling barrier [4]. On the other hand, rather high degree of spin current polarization has been detected if the spin is injected from magnetic semiconductors [5], although its high efficiency is restricted to low temperatures.

In this Letter we propose a new method to use a dc current to inject spin polarization, not from any ferromagnetic material, but from a quantum well (QW) embedded into a $pn$ junction. The spin polarized current is generated during the radiative electron-hole recombination in the QW, accompanied by the emission of circularly polarized light. Our new injection mechanism can be easily explained. It is well known that in a QW which is asymmetric along the growth direction, at each finite wave vector $k$ the spin degeneracy of hole subbands is removed. The corresponding splitting of hole energies increases with $k$ and can reach a quite high value. For example, in a $p$ inversion layer of a GaAs-AlGaAs heterojunction, the splitting of the topmost heavy-hole subband was calculated [6] to be about 5 meV at $k=2\cdot10^6$ cm$^{-1}$. Each of the split states is a linear combination of four angular momentum eigenstates which are specified by the $z$-component $J_z=\pm3/2$ and $\pm1/2$. The resulting mean spins of hole states are parallel to the QW interfaces. For a given $k$ the spins in each pair of spin-split subbands have opposite directions, and in a given spin-split subband the mean spin at $-k$ is reversed with respect to that at $k$. The spin orientations in the topmost heavy-hole subband are shown in Fig. 1, where the Fermi energy $\mu_e$ (or $\mu_h$) of the quasiequilibrium degenerate electron gas (or hole gas) is indicated. When a $\sigma$-spin electron in the $k$ state of the conduction band makes a radiative transition to the topmost heavy-hole subband, as marked by the downward vertical arrow line in Fig. 1, the probability of this process depends on $\sigma$ because the population of hole at the $k$ state in the split $(+)$-subband may be different from that in the split $(-)$-subband. Hence, at $k$ the conduction band electron spin will be polarized along $\sigma$ direction if an electron with $(-\sigma)$-spin has higher recombination rate. However, since the hole spin orientation at $-k$ is reversed with respect to that at $k$, the spin polarization of the conduction band electron at the $-k$ state will be along $-\sigma$ direction. Consequently, if the momentum distributions of both the electron gas and the hole gas are isotropic, there will be no net spin polarization of the conduction band electrons.

On the other hand, if the momentum distributions of the electron gas and/or the hole gas are anisotropic, the generation of spin polarization becomes possible. We will explain the generating process with the help of Fig. 1, where the anisotropic momentum distribution is indicated by a shift $\delta k$ of the quasi-equilibrium momentum distribution of electrons with respect to that of holes. The state 3 in the $(+)$-subband and the state 4 in the $(-)$-subband are equally occupied by holes with spins in opposite directions. Hence, the total probability of recombination of a conduction electron with the two holes in the states 3 and 4 is independent of the spin of the electron. However, a hole appears in the state 1 in the $(+)$-subband but not in the state 2 in the $(-)$-subband. Hence, the conduction electrons in the shaded region can recombine only with the holes in the $(+)$-subband, and
the corresponding recombination probability is spin dependent. Such processes create a nonequilibrium electron spin polarization in the QW. When this polarization diffuses out of the QW, a spin current is then generated in the n-doped region. One way to realize the situation shown in Fig. 1 is to apply an electric field \( \mathbf{E} \) parallel to interfaces. Then, both electrons and holes gain their respective drift velocities. Let \( \mathbf{v} \) be the relative drift velocity of electrons with respect to drift velocity of holes. If we ignore the drift of the low mobile holes, the resultant band positions are illustrated in Fig. 1. The shift \( \delta \mathbf{k} \) is simply \( m^* \mathbf{v} / \hbar = m^* \mu \mathbf{E} / \hbar \), where \( \mu \) is the electron mobility. The suggested device for spin current injector, a QW embedded into a \( pn \) junction, is illustrated in Fig. 2.

A bias voltage \( V_b \) is applied along the \( z \) axis, and the transverse voltage source \( V_t \) creates an electric field \( \mathbf{E} \) along \( x \) axis. The potential profile along \( z \) axis and the spin current are also shown schematically in Fig. 2.

![Diagram](image)

**FIG. 2.** Suggested device for spin current injector. Upper plot is the corresponding potential profile.

At low temperatures, for degenerate electron and hole gases, besides the band to band electron transitions, excitonic recombination processes must be considered. To avoid the complicated analysis which will not change the essential physics, in this Letter we will neglect the exciton effect on the spin generation.

For the valence band we define \( |J_z\rangle \) as the zone center Bloch state and \( \psi_{k,n,J_z}(z) \) the \( n \)-th confined state in the QW associated to the two-dimensional wave vector \( \mathbf{k} \) and the hole spin projection \( J_z \). Then, the wave functions \( \Psi_{n,k}(z) \) of the \( n \)-th \((\pm)\) split valence subbands can be expressed as the sum of \( \psi_{k,n,J_z}(z)|J_z\rangle \) over \( J_z = \pm \frac{3}{2} \) and \( \pm \frac{1}{2} \). For electrons in the lowest conduction band, we can similarly define \( |\sigma\rangle \) and \( \chi(z) \). If we neglect the small spin-orbit splitting, the electron wave functions \( \Psi_{k}(z) \) can be written as a linear combination of two degenerate states \( C_{\frac{3}{2}} \chi(z)|\frac{1}{2}\rangle + C_{\frac{1}{2}} \chi(z)|-\frac{1}{2}\rangle \), where \( C_{\pm \frac{1}{2}} \) are normalized amplitudes. When an electron and a hole recombine from their respective quantum states to emit a photon of wave vector \( \mathbf{q} \) and polarization vector \( \mathbf{e}_q^\lambda \), to the lowest order, the quantum probability amplitude of this process can be easily derived as

\[
M_{\sigma,n}(\mathbf{k}, \mathbf{q}) = \frac{eA_q}{mc} \sum_{J_z} p_{\sigma,J_z} \cdot \mathbf{e}_q^\lambda \int dz \psi_{k,n,J_z}(z) \chi(z).
\]

In the above equation, \( p_{\sigma,J_z} = \langle J_z | p | \sigma \rangle \) is the interband matrix element of the electron momentum operator \( \mathbf{p} \), and \( A_q = \sqrt{2\pi e/\hbar nq} \).

Knowing the recombination probability of the conduction band electrons with a specific spin, the spin generation during the recombination process can be calculated. The resultant spin density in the conduction band is derived as

\[
G_x = - \sum_{\sigma \sigma'} G_{\sigma \sigma'} s_{\sigma \sigma'},
\]

where \( 2s_{\sigma \sigma'} \) are Pauli matrices. The matrix elements \( G_{\sigma \sigma'} \) have the form

\[
G_{\sigma \sigma'} = A \sum_{\mathbf{k}, \mathbf{q}} M_{\lambda,\nu}^\lambda(\mathbf{k}, \mathbf{q}) M_{\sigma,\sigma'}(\mathbf{k}, \mathbf{q}) f_\sigma f_{\sigma'}(\mathbf{k}),
\]

where \( A \) is a numerical factor, and \( f_\sigma(\mathbf{k}) [ f_{\sigma'}(\mathbf{k}) ] \) is the momentum distribution functions of electrons (or holes). Within the perturbation theory, \( f_\sigma(\mathbf{k}) \) and \( f_{\sigma'}(\mathbf{k}) \) are assumed to be spin independent. A relevant dimensionless parameter which measures the efficiency of spin generation is

\[
P = \frac{G_x}{G}, \quad G = G_{\frac{1}{2},\frac{1}{2}} + G_{-\frac{1}{2},-\frac{1}{2}}
\]

is the number of radiative recombinations per the time unit.

As mentioned earlier, we will consider higher temperatures such that the excitonic effect can be neglected. Hence, each momentum distribution function is the sum of the equilibrium Boltzmann distribution and the nonequilibrium correction due to the electric field \( \mathbf{E} \).

Since the drift of holes can be ignored, we have

\[
f_\nu(\mathbf{k}) = \exp \left[ \frac{\mu_\nu - \epsilon_\nu(\mathbf{k})}{kBT} \right],
\]

\[
f(\mathbf{k}) = \exp \left[ \frac{\mu - \epsilon(\mathbf{k})}{kBT} \right] \left( 1 - \frac{\hbar \mathbf{v} \cdot \mathbf{k}}{kBT} \right),
\]

where \( \epsilon(\mathbf{k}) \) is the lowest conduction subband, and \( \epsilon_\nu(\mathbf{k}) \) \((\pm)\) split valence subbands. The zero reference energy is set at the bottom of the lowest conduction subband.

We need the valence subband wave functions \( \psi_{k,n,J_z}(z) \) for calculating \( M_{\sigma,\sigma'}(\mathbf{k}) \). Let the growth direction \( z \) be along the [001] crystal. The wave functions can be obtained by applying a proper unitary transformation \( \mathbf{B} \) which block diagonalizes the Luttinger Hamiltonian.

They can be expressed in the general form as

\[
\psi_{k,n,\pm \frac{1}{2}}^\pm(z) = \pm \psi_{k,n,\pm \frac{1}{2}}^\pm(z) = \pm i \exp (i\phi_k) \xi_{\pm \frac{1}{2}}(z)/\sqrt{2},
\]

\[
\psi_{k,n,-\frac{1}{2}}^\pm(z) = \pm \psi_{k,n,-\frac{1}{2}}^\pm(z) = \pm i \exp (i\eta_k) \xi_{-\frac{1}{2}}^\pm(\mathbf{k}, \mathbf{q})/\sqrt{2},
\]
where the real functions $\xi^\pm_{n\ell}(z)$ and $\xi^\mp_{n\ell}(z)$ represent the partial amplitudes of heavy and light holes in the (+)- and the (−)-state. The phases are $\phi_k=3n_k=3\phi_k/2$ with $\phi_k=\cos^{-1}(k_x/k)$. To obtain $\xi_k$ we have neglected the band warping. For convenience we set the $x$-axis along the drift velocity direction.

With all above equations (1)-(8) we are ready to calculate the efficiency of spin generation $P=\Omega_s/G$. In terms of the overlap integrals $b_{n\ell}^\pm=\int dz\xi^\pm_{n\ell}(z)\chi_k(z)$ and $b_{n\ell}^z=\int dz\xi^z_{n\ell}(z)\chi_k(z)$, we define $R_{x,nk}=-(b_{n\ell}^z)^2/3$, $R_{y,nk}=-(b_{n\ell}^x)^2/3$, and $R_{z,nk}=0$. Then, $P$ is derived as

$$P = \frac{\hbar v}{2\kappa T} \left( \sum_{k,n,\nu} \nu \mathbf{k} \cdot \mathbf{R}_{nk} \cdot \mathbf{F}_{k,n,\nu} \right)$$

$$\times \left( \sum_{k,n,\nu} \left[ (b_{n\ell}^x)^2 + (b_{n\ell}^y)^2 / 3 \right] \mathbf{F}_{k,n,\nu} \right)^{-1},$$

where $\mathbf{F}_{k,n,\nu} = \exp \{-[\nu^2(k) + (\nu^2(k))]/k_BT\}$. As expected, only the anisotropic part (ν·k term) of the electron momentum distribution function $f(k)$ has contributed to the generated polarization.

Based on the above formula, $P$ can be investigated numerically in a broad range of parameters. However, qualitative results can be derived from (6) analytically for $kd \ll 1$, where $d$ is the typical length of hole confinement in the QW. We notice that in (6), a characteristic wave vector $\xi_k$ can be obtained at $k_BT \gg 1$, where the real functions $\xi_k$ become proportional to $1 - \exp(-\Delta(k)/k_BT)$.

In the range of temperatures where $\Delta(k)/k_BT$ is due to the anisotropic momentum distribution of the electrons. The other factor $\Delta_1(kT)(kTd)^2/k_BT$ is originated from the ratio of the two summations in (6), and increases with $T$ through $kT$ and $\Delta_1(kT)$. For example, in GaAs at room temperature, $kT \simeq 10^6$ cm$^{-1}$, and then $kT_2 \simeq 1$ for $d=100$ Å. Under this situation, $\Delta_1(kT)$ becomes comparable to the quantization energy of hole subbands $\xi_k$, and (6) is no longer valid. However, a simple scaling analysis shows that around $kT_2 \simeq 1$, Eq. (6) can still be used to evaluate $P$ with the factor $(kTd)^2$ replaced by a numerical factor of the order of unity.

We have shown that spin polarization of conduction electrons can be generated in a quantum well due to the radiative electron-hole recombination. Let $|e|I$ be the electric current across the pn junction, and $n=n_r/\tau_e + \tau_r$ is the luminescence quantum efficiency, where $\tau_e$ and $\tau_r$ are the radiative (or nonradiative) electron-hole recombination time. Then, the number of radiative recombinations per time unit is $G=|e|I$. The so generated spin polarization in the conduction band can either diffuse out of the QW into the n-doped region of the pn junction, or relax within the QW with a spin relaxation time $\tau_{sw}$. In III-V semiconductor QWs the dominating process is the D'yakonov-Perel’ spin relaxation, although the Bir-Pikus mechanism can be efficient due to presence of a large number of holes. In the steady state, the spin current which flows out of the QW is given as

$$I_e = \mathcal{P} \eta n - S/\tau - S/\tau_{sw},$$

where $S$ is the spin polarization, and $1/\tau=1/\tau_r+1/\tau_e$. We choose the spin quantization axis parallel to the direction of the vector $\mathbf{P}$ defined in (6). Then, $S$ is simply $(n_h-n_e)/2$, where $n_s$ is sheet electron density in the QW with $\sigma$-spin.

To relate $S$ to $I$, we need to specify the transport process between the QW and the n-doped region, which involves the spin diffusion and relaxation in the n-doped semiconductor. To avoid such complications, we will consider a simple model of thermionic transport over the barrier, which is suitable for not too low temperatures. In this case the $\sigma$-spin current $i_\sigma$ is determined by the balance of Richardson currents emitted from both sides of the potential barrier. Such emission currents depend on the chemical potentials $\mu_{\sigma w}$ in the QW and $\mu_{\sigma h}$ in the n-doped bulk semiconductor. In the linear response regime, $|\mu_{\sigma w}-\mu_{\sigma h}| \ll k_BT$, and we have

$$i_\sigma = \frac{A^* T^2}{|e|} \frac{\mu_{\sigma w} - \mu_{\sigma h}}{k_BT} \exp \left( \frac{\mu_{\sigma h} - eU}{k_BT} \right),$$

where $A^*$ is the Richardson constant, and the barrier height $U$ is indicated in Fig. 2. With this expression, the
spin current $I_s=(\hat{\imath}\hat{x}-\hat{\imath}\hat{y})/2$ is a function of the chemical potentials.

In the $n$-doped bulk semiconductor we assume diffusive motion of electrons with a diffusion constant $D$ and a spin relaxation time $\tau_{sb}$. We should mention that within the framework of linear transport theory the spin current in a bulk material can be driven only by a spin density gradient, and the corresponding characteristic length of the spin density variation has the form $L_s=\sqrt{D\tau_{sb}}$. In an $n$-doped bulk semiconductor $\tau_{sb}$ can be very long [3], and a large $L_s$ will reduce the efficiency of spin injection. This is the same problem encountered in the study of spin injection from ferromagnetic metals into semiconductors [3], which can be overcome [4] with a proper choice of the barrier height $U$. If $U$ is so chosen that $(L_s/l) \exp ([\mu_{sb} - eU]/k_BT) \ll 1$, where $l$ is the electron mean free path, the injected spin current is no longer limited by the low spin relaxation rate of the bulk semiconductor. For a moderately doped III-V semiconductors, if we take from [3] $\tau_{sb} = 100$ ps and the electron mobility $5\times10^4$cm$^2$/Vsec, $L_s/l \geq 10$ at room temperature. Therefore, the above inequality can be easily satisfied even if the barrier height is rather low. In this situation we obtain

$$I_s = I_P \eta \left[ 1 + \frac{\Delta \mu}{k_BT} (1 + \frac{\tau}{\tau_{sw}}) \right]^{-1},$$

(10)

where $\Delta \mu = (\mu_{\hat{\imath} \hat{x}, w} - \mu_{\hat{\imath} \hat{x}, s} - \mu_{\hat{\imath} \hat{y}, s} - \mu_{\hat{\imath} \hat{y}, w})/2$.

Although $\Delta \mu/k_BT$ is assumed to be small in linear response theory, the ratio $\tau/\tau_{sw}$ can be large, and so the second term in the square brackets of (10) can not be neglected. For example, in bulk III-V semiconductors with carrier density in the range between $10^{17}$ cm$^{-3}$ and $10^{18}$ cm$^{-3}$, at room temperatures the ratio $\tau/\tau_{sw}$ is around ten [3], assuming $\tau$ in the range of nanosecond. The similar ratio was also found in QWs [4]. In our system, as illustrated in Fig 2, the electron density and the hole density in the asymmetric QW are spatially shifted with respect to each other. The electron-hole recombination becomes spatially indirect, and therefore $\tau$ increases. On the other hand, $\tau_{sw}$ increases if the spin relaxation is dominated by the Bir-Pikus mechanism, but remains almost the same if the D'yakonov-Perel' mechanism dominates. Consequently, the actual value of $\tau/\tau_{sw}$ varies with the system to be studied. From (10), we see that the upper limit of the spin polarization of the injected current is $I_s/I = P \eta$.

Since the spin polarization in the QW is created by the emission of circularly polarized light, the spin generation can be detected by measuring the polarization of emitted photons. Circular polarization of a photon is represented by the imaginary off-diagonal elements of the photon polarization matrix $\rho$. This Hermitian matrix can be calculated in a way similar to the calculation of spin polarization. It can be shown that the imaginary elements are related to $\rho$ as $\rho_{x} \propto iP_{x}$ and $\rho_{z} \propto -iP_{z}$. Hence, most of the circularly polarized photons are emitted with their wave vectors parallel to $\mathbf{P}$.

If we introduce a ferromagnetic layer to the $n$-doped part of our sample, the optoelectric injection of spins can be investigated by measuring the resistance of the system. The separation of the ferromagnetic layer and the QW must be less than $L_s$. In this case, the measured resistance depends on relative spin polarizations in the paramagnetic and the ferromagnetic materials [3], and can be varied by changing the direction of the applied electric field $E$.

We close this Letter with one remark. Within the framework of linear response theory used in our analysis, the ratio $\hbar \delta \mu/k_BT$ must be small. With a stronger electric field to increase the electron drift velocity $v$, we conjecture an enhancement of the spin generation under an anisotropic momentum distribution of hot electrons.

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