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LETTER TO THE EDITOR

A perfect spin-filter quantum dot system

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Abstract
The discovery of a novel effect in the transport through a QD spin-dependently coupled to magnetic contacts is reported. For a finite range of source–drain voltages the spin projections of the current cancel exactly, resulting in a completely suppressed output current. The spin down current behaves as one normally expects whereas the spin up current becomes negative. As the source–drain voltage is increased the spin up current eventually becomes positive. Thus, tuning the source–drain voltage such that the spin up current vanishes will result in a perfect spin filter.

Many-body effects are known to be of large importance for the transport behaviour of mesoscopic systems. For instance, the Kondo effect provides a zero-bias anomaly [1–3] in the conductance of a quantum dot (QD) weakly coupled to external contacts. Quantum point contacts display the so-called 0.7 structure at the onset of a single-conducting mode [4, 5] and in (quasi-)two-dimensional systems under the influence of magnetic fields one finds the quantum Hall [6] and fractional quantum Hall effects [7–10]. Common to all these effects are many-body interactions occurring in the systems. In this letter we report the discovery of an unexpected effect in the transport through a QD spin-dependently coupled to magnetic contacts (see figure 1), an effect which stems from strong electron correlations in the system. Within a finite range of source–drain (bias) voltages applied over the system the two different spin projections of the current cancel exactly, producing a completely suppressed total output current. Hence the spin down electrons in the device shown in figure 1 obey the well known transport laws, whereas the spin up electrons break the intuitive picture resulting in a *negative* current. The device shown in figure 1 hence very effectively sorts out different spin projections of the current, something which is highly desired in spin electronics [11–15]. The effect is due to a strongly

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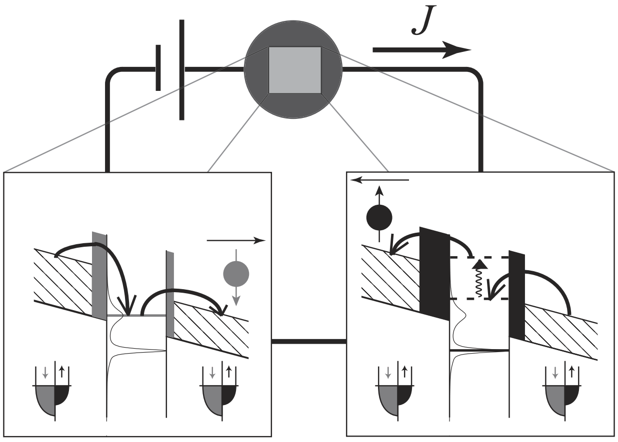

Figure 1. A cartoon of the QD system asymmetrically and spin-dependently coupled to the contacts. The greyscale code in the insets signifies the spin polarization with black and grey representing spin↑ and ↓ channel, respectively. The spin↑ (↓) electrons tunnel through the black (grey) tunnel barriers. An example of the non-equilibrium local density of states in the QD (bias voltage ∼1 mV) is shown between the tunnel barriers (solid curve), clearly indicating a highly occupied spin ↑ (black) state and a less occupied spin↓ state (grey). In the spin↑ channel (right inset), there are two virtual bound states, due to spin-flip scattering, slightly above the chemical potential of the left and right contacts. The spin polarization in the left (right) contact is drawn schematically in the lower left (right) corner.

spin-dependent renormalization of the many-body states in the QD, induced by the spin-dependent couplings to magnetized contacts. We emphasize that the negative spin up current is only obtained if a correct many-body treatment of the occupation numbers of the QD states is considered. Experimentally, our findings should be possible to verify by coupling a single-electron QD to the magnetic contacts via spin-dependent tunnel barriers [16, 17]. Moreover, the coupling to the left and the right contacts should be strongly asymmetric (see figure 1).

Previously, spin-filter systems based on various mechanisms have been proposed [18–23]. In most of these reports, magnetic fields, applied directly over the device, have been used to provide a spin current through the device [18–21]. Another possibility is given by coupling a ferromagnetic metal to a semiconductor via a semiconductor nanocrystal doped by a paramagnetic ion [22]. By exploiting a combination of spin splitting of the resonant level induced by the Rashba effect and the spin blockade phenomenon, it was proposed in [23] that a more than 99.9% pure spin current can be achieved in a triple-barrier resonant tunnelling structure. All these studies, however, were based on the assumption that contributions from spin-flip scattering processes are negligible. Recently it was theoretically discovered that spin current can be driven through a QD in the presence of an oscillating magnetic field [24], although the total charge current is zero. The external magnetic field introduces a spin splitting in the QD and separated spin chemical potentials in the single contact. The spin current driven through the system is then activated by transfer of electrons at the higher spin chemical potential to the QD where the spin-flip processes permits the electrons to proceed into the lower spin chemical potential. In this letter, we consider ferromagnetic contacts, spin-dependent tunnel barriers and include spin-flip scattering processes in the QD. The spin polarization in the contacts and/or the spin-dependent tunnel barriers induces a spin splitting of the localized levels in the QD which allow energy transfer between the spin channels. The effects from the spin-flip scattering in the QD is strongly amplified by the combination of the magnetic contacts and spin-dependent tunnel barriers.

In low density QDs, such as single-level devices, the correlation between the different spin states of the level is very strong. Both theoretically and experimentally it has been shown that the Coulomb repulsion, U, is in the order of or larger than the level separation for QDs in the size range of tens of nanometres. Naturally, then such systems cannot be treated within a non-interacting Landauer–Büttiker picture [25–27]. One rather has to employ more sophisticated tools, such as many-body theory for non-equilibrium Green functions. Being the largest parameter in the present formulation, the intra-dot Coulomb repulsion is taken into account exactly from the very beginning. The QD detached from the contacts is modelled by the Hamiltonian $\mathcal{H}_{\rm QD} = \sum_\sigma \varepsilon_\sigma d_\sigma^\dagger d_\sigma + U n_\uparrow n_\downarrow = \sum_{p=0,\uparrow,\downarrow,2} E_p X^{pp}$, where $d_\sigma^\dagger$ (d_σ) creates (annihilates) a spin σ ($\sigma = \uparrow, \downarrow$) electron in the QD at the energy ε_σ and $n_\sigma = d_\sigma^\dagger d_\sigma$. The term $U n_\uparrow n_\downarrow$ represents the energy of correlation between two electrons. The expression on the right-hand side is a diagonal form of the states in the QD written in terms of the many-body (Hubbard) operators [28, 29] $X^{pp} = |p\rangle\langle p|$ corresponding to the energies E_p of the empty $|0\rangle$, singly occupied $|\sigma\rangle$ and doubly occupied $|2\rangle \equiv |\uparrow\downarrow\rangle$ QD states. Because of the large Coulomb repulsion the doubly occupied state lies far above the range of conduction and can be neglected.

Having the QD expressed in terms of transitions between many-body states we introduce the transition energies $\Delta_{\sigma 0}^0 = E_\sigma - E_0$ corresponding to the transitions $|0\rangle \to |\sigma\rangle$ in the QD. As the contacts are being attached to the QD, the first diagrammatic correction to the transition energies yields the renormalization according to the self-consistent equation [30]

$$\Delta_{\sigma 0} = \Delta_{\sigma 0}^0 + \frac{1}{2\pi} \sum_{k\in {\rm L,R}} |v_{k\bar\sigma}|^2 \int \frac{f(\varepsilon_{k\bar\sigma}) - f(\omega)}{\varepsilon_{k\bar\sigma} - \omega} [-2\,{\rm Im}\, D^{\rm r}_{\bar\sigma\bar\sigma}(\omega)]\,{\rm d}\omega. \qquad (1)$$

In equation (1) $v_{k\bar\sigma}$ represents the hybridization between the QD states and the states in the contacts, $f(x)$ is the Fermi function, whereas $-2\,{\rm Im}\, D^{\rm r}_{\bar\sigma\bar\sigma}(\omega)$ is the spectral density of the transition $|0\rangle \to |\bar\sigma\rangle$, however, without its spectral weight; ω is the energy variable. Throughout the letter, $\bar\sigma$ signifies the spin opposite to σ. The expression given in equation (1) differs from that given in [30], in that we here have included the finite width of the transition $|0\rangle \to |\bar\sigma\rangle$ which yields the integration over the spectral density. It should be noticed that the renormalized transition energy for, say, the spin up ($\sigma = \uparrow$) state strongly depends on the properties of the spin down ($\bar\sigma = \downarrow$) channel in the system. The renormalization depends both on the tunnelling probability for the spin down electrons, reflected in $v_{k\bar\sigma}$, and on the spin polarization, that is the magnetization, of the conduction electrons at the energies $\varepsilon_{k\bar\sigma}$ in the contacts.

The renormalization also strongly depends on the asymmetry of the couplings to the left (L) and the right (R) contacts. An asymmetric coupling of the QD to the left/right contacts will provide a different renormalization of the transition energy in a forward and a backward biased system. As a consequence of the many-body origin of the renormalization, the fillings of the states in the contacts and the QD are reflected in the presence of the Fermi functions $f(\varepsilon_{k\bar\sigma})$ and $f(\Delta_{\bar\sigma 0})$. Therefore, different tunnelling probabilities and/or a spin polarization in the contacts removes the degeneracy of the transition energies to the spin up and spin down states. The widths of the transitions are proportional to the couplings to the contacts and their corresponding *spectral weights*. The latter quantity includes spin-flip scattering processes in the QD. In order to avoid lengthy derivations in this letter, which are given in [31], we only present the final algebraic expression for the first diagrammatic correction to the spectral weight for the transition $|0\rangle \to |\sigma\rangle$, given by

$$\mathbb{P}_{\sigma\sigma}({\rm i}\omega) = P_\sigma - \frac{P_{\bar\sigma} - P_\sigma}{2\pi} \int \sum_{k\in{\rm L,R}} \frac{|v_{k\bar\sigma}|^2}{\varepsilon_{k\bar\sigma} - \omega'} [-2\,{\rm Im}\, D^{\rm r}_{\bar\sigma\bar\sigma}(\omega')] \times \left(\frac{f(\varepsilon_{k\bar\sigma}) - [n_{\rm B}(\Delta_{\sigma\bar\sigma}) + 1]}{{\rm i}\omega - \Delta_{\sigma\bar\sigma} - \varepsilon_{k\bar\sigma}} - \frac{f(\omega') - [n_{\rm B}(\Delta_{\sigma\bar\sigma}) + 1]}{{\rm i}\omega - \Delta_{\sigma\bar\sigma} - \omega'} \right) {\rm d}\omega'. \qquad (2)$$

Here $P_\sigma = N_\sigma + N_0$ is the sum of the population numbers for the states $|\sigma\rangle$ and $|0\rangle$ of the isolated QD, $\Delta_{\sigma\bar{\sigma}} = \Delta_{\sigma 0} - \Delta_{\bar{\sigma}0}$ is the energy of the spin-flip transition $|\bar{\sigma}\rangle \to |\sigma\rangle$ and $n_{\mathrm{B}}(x)$ is the Bose function. It is clear from the expression in equation (2) that whenever the system is in a completely spin-degenerate configuration, the spectral weight $\mathbb{P}_\sigma(\mathrm{i}\omega) = P_\sigma$. For this situation effects from spin-flip scattering events are attributed to the usual Kondo effect [1–3]. However, if the spin degeneracy in the system is broken, the second term in equation (2) becomes non-zero, and the spin flips in the QD open new channels for tunnelling. These channels correspond to the spin-flip energies $\Delta_{\sigma\bar{\sigma}} \neq 0$ around the (quasi-)chemical potential μ_α of each contact $\alpha = \mathrm{L, R}$, through which electrons may pass.

It should be noted that the local Green functions (GFs) of the QD, $G_{\sigma\sigma}(t, t') \equiv (-\mathrm{i})\langle T X^{0\sigma}(t) X^{\sigma}(t')\rangle = \sum_{\sigma'} D_{\sigma\sigma'}(t, t')\mathbb{P}_{\sigma'\sigma}(t)$, are self-consistently solved for each value of the bias voltage. This is clear since the locators and end factors depend on each other and since the bare end factors ($P_\sigma = N_0 + N_\sigma$) are found from $N_\sigma = \mathrm{Im}\int G^{<}_{\sigma\sigma}(\omega)\,\mathrm{d}\omega/(2\pi)$ requiring[6] charge conservation $1 = N_0 + \sum_\sigma N_\sigma$. The same mathematical technique was employed for a theoretical description of asymmetric negative differential conductance [32] and current–voltage asymmetries [33] in double QDs. More details on the diagrammatic technique for Hubbard operator GFs can be found in [30, 31, 34, 35].

An example of the properties of the transport through the device shown in figure 1 is illustrated in figure 2; the currents J_σ are computed at 80 points in the bias voltage interval. We stress that the negative spin↑ current, in the bias voltage between 0 and 1.5 mV in figure 2 (dotted curve), will *not* be obtained without considering the second term in equation (2). This term has not been considered in a non-equilibrium situation before and from a technical, many-body point of view, equation (2) is a novel result. The stationary net current through the system is calculated by means of the conventional formula for the current through a mesoscopic interacting region [36], i.e.

$$J \sim -\,\mathrm{Im}\sum_\sigma \int \{[\Gamma^{\mathrm{L}}_\sigma - \Gamma^{\mathrm{R}}_\sigma] G^{<}_{\sigma\sigma}(\omega) + [f_{\mathrm{L}}(\omega)\Gamma^{\mathrm{L}}_\sigma - f_{\mathrm{R}}(\omega)\Gamma^{\mathrm{R}}_\sigma][G^{\mathrm{r}}_{\sigma\sigma}(\omega) - G^{\mathrm{a}}_{\sigma\sigma}(\omega)]\}\,\mathrm{d}\omega.$$

where $\Gamma^{\mathrm{L/R}}_\sigma = 2\pi\sum_{k\in\mathrm{L/R}} |v_{k\sigma}|^2\delta(\omega - \varepsilon_{k\sigma})$ and $f_{\mathrm{L/R}}(\omega) = f(\omega - \mu_{\mathrm{L/R}})$.

The most spectacular finding from our theoretical analysis, displayed in figure 2, comes from analysing the spin projections of the current. For bias voltages in the range (0, 1) mV, the current $J_\downarrow > 0$ whereas $J_\uparrow < 0$, which results in a vanishing net current $J = J_\uparrow + J_\downarrow = 0$ (in the chosen parameter space [2, 3, 17]). This means that while the spin down electrons are travelling against the direction of the electric field, the spin up electrons are moving along the field. Figure 2 shows that within the interval of bias voltages between 0 and 1 mV, the spin up and spin down currents are equally large but oppositely directed resulting in a zero net current. This opens up the possibility of using the present device as a perfect spin-filter system, with potential applications for the spin-electronics industry [15]. We stress, however, that it is not the vanishing total current that should be of most technological interest. Rather, it is the possibility of tuning the bias voltage such that the spin↑ current vanishes (changes sign) that provides the possibility of a 100% pure spin current.

The effect predicted here cannot be explained intuitively, since the current of the spin up states seems to violate Coulombic forces. However, our finding is the result of a rigorous mathematical treatment of transport through QDs; hence the transport conservation laws for e.g. charge, spin and current are satisfied. Physically, the result may be explained as follows. First of all, we stress that the conservation of charge cannot be separated into conservation of charge in each spin channel. This depends on the spin symmetry being broken due to the

[6] It should be noticed that the lesser DQD GF in the present notation (algebraically) is given by $G^{<}_\sigma = G^{\mathrm{r}}_\sigma V^{<}_\sigma G^{\mathrm{a}}_\sigma + D^{\mathrm{r}}_\sigma \mathbb{P}^{<}_\sigma (1 + V^{\mathrm{a}}_\sigma G^{\mathrm{a}}_\sigma)$, where $V_\sigma(\mathrm{i}\omega) = \sum_{k\in\mathrm{L,R}} |v_{k\sigma}|^2/(\mathrm{i}\omega - \varepsilon_{k\sigma})$.

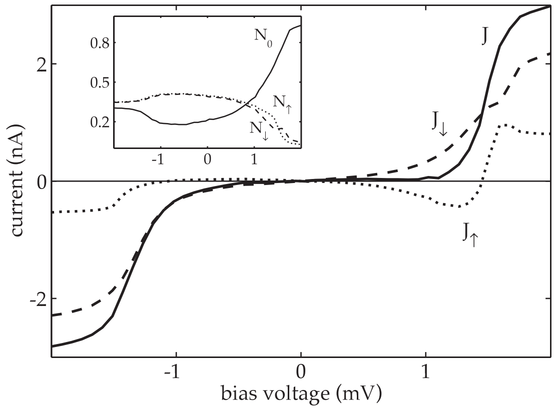

Figure 2. The current–voltage characteristics of the QD spin-dependently coupled to magnetic contacts in an asymmetric fashion with respect to the left and the right couplings. The inset shows the non-equilibrium population numbers N_0, $N_\uparrow$, $N_\downarrow$ of the empty and singly occupied states, $|0\rangle$ and $|\sigma\rangle$, $\sigma = \uparrow, \downarrow$, respectively. Here, the couplings $\{\Gamma^L_\uparrow, \Gamma^R_\uparrow, \Gamma^L_\downarrow, \Gamma^R_\downarrow\} \sim 5.6 \times \{1, 11.5, 4, 46\}$, the bare energy $\Delta^0_{\sigma 0} \sim -500$, the lower cut-offs of the conduction bands (bandwidths $\sim 2 \times 10^6$) in the contacts $W^{L,R}_\uparrow \sim -10^5$, $W^{L,R}_\downarrow \sim -1.9 \times 10^6$ and $k_B T \sim 10^2$ (units: μeV).

strong coupling and processes of spin flip between the spin channels, which is illustrated by equations (1) and (2). For the spin↑ projection, the QD state lies far below the chemical potentials of both contacts. In addition, the population number for that state is sufficiently high to prevent a significant tunnelling through this state. In the bias voltage range between 0 and 1.5 mV, spin↑ electrons in the left contact are hence reflected at the tunnelling interface to the left of the QD. However, due to spin-flip processes, a virtual level at $\varepsilon_{k\uparrow} + \Delta_{\downarrow\uparrow}$ is found to lie above μ_α ($k \in \alpha = \mathrm{L, R}$) and it opens for conduction. For spin↓ electrons the corresponding level lies at $\varepsilon_{k\downarrow} + \Delta_{\uparrow\downarrow} < \mu_\alpha$. The spin↑ electrons travelling along the field in the right contact couple sufficiently strongly to the QD to enable tunnelling into the levels $\varepsilon_{k\uparrow} + \Delta_{\downarrow\uparrow}$, due to energy transfer between the spin channels. This mechanism causes the spin↑ current to become negative. This possibility is, however, unavailable for the spin↓ electrons since energy is transferred from the spin↓ channel to spin↑ and not vice versa. On the other hand, the spin↓ state in the QD has a sufficiently low occupation number to assist regular tunnelling through this state. Therefore the spin↓ current is positive.

Transport in conventional materials is often viewed as being due to electrons that move back and forth in the system, sometimes in the direction of an applied field and sometimes against the direction of the field, but with an average drift velocity in the direction against the field which, by definition, is a positive current. As is well known, this view has been emphasized when transport in nanomaterials is considered, as reflected by, e.g., non-Ohmic behaviour [1–10, 30, 36, 37]. In the currently suggested device the situation is even more complex—due to many-body effects in the QD, the net current vanishes as a result of the compensation of two non-zero currents with opposite spin projections and opposite directions. However, this does not mean that a negative spin↑ current would be measured by blocking the transport of the spin↓ states in the contacts, since it is the interactions between electrons with opposite spins that cause the effect discussed. Nevertheless, a pure spin↓ current should be measurable when the bias voltage is tuned to the value where the spin↑ current vanishes; see figure 2.

In conclusion, we report the discovery of an unexpected effect in the transport behaviour of a QD spin-dependently coupled to magnetic contacts, arising due to strong electron correlations in the QD. The spin dependence of the contacts and the tunnel barriers induces a spin splitting of the localized level in the QD. Spin-flip scattering processes within the QD create an energy transfer between the two spin channels. Hence, within a finite range of bias voltages applied over the system the two spin projections of the current cancel exactly, thus providing a vanishing total current. The spin down electrons in the system obey the well known transport laws, whereas the spin up electrons break the intuitive picture resulting in a *negative* current. For increasing bias voltages the spin↑ current eventually becomes positive. The system thus constructed effectively sorts out the spin projections and tuning the system to a certain finite bias voltage, such that the spin↑ current vanishes, will result in a perfect spin-filter device. Further (experimental) investigations are highly desired.

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