Back-action-driven electron spin excitation in a single quantum dot

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New Journal of Physics 15 (2013) 023021 (10pp)
Received 22 October 2012
Published 13 February 2013
Online at http://www.njp.org/
doi:10.1088/1367-2630/15/2/023021

Abstract. We perform real-time charge counting with a quantum point contact (QPC) for the last six electrons in a single quantum dot. At zero magnetic field, the charge-counting statistics show distinctive non-thermal-equilibrium effects for the even and odd electron numbers. A detailed study relates this difference to the excitation from the spin singlet state to triplet states driven by QPC back-action. At a finite magnetic field, spin excitations to different triplet states and Zeeman states are also observed. A master-equation model is developed to quantitatively characterize the back-action-driven spin excitation rate.

Online supplementary data available from stacks.iop.org/NJP/15/023021/mmedia

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1. Introduction

Individual electron spins in semiconductor quantum dots (QDs) are a prospect for implementing solid-state quantum computers [1, 2]. To read out the state of a spin-based qubit, a commonly used scheme is to monitor the status of charge occupation on a QD with a nearby quantum point contact (QPC). The QPC measurement has an inevitable back-action on the QD in various forms as shown by extensive theoretical and experimental research [3–7]. The QPC bias voltage could cause direct back-action through the photon-mediated QD–QPC Coulomb coupling [3, 8]. The QPC current could also cause indirect back-action if the QD electrons re-absorb the high-frequency quasi-particles such as acoustic phonons emitted by the heated electron reservoir [9, 10]. Experimentally, the assisted tunneling through higher energy levels has been widely used as a tool to study the QPC back-action [8–10].

In this work, particular attention is paid to the effect of back-action on the electron spin states in a single QD. The back-action strength is tunable through the QD–QPC coupling or QPC bias, and is configured at a high level in this experiment. We perform real-time charge counting for the last six electrons and observe dramatic signatures of excitation for the even electron numbers at zero magnetic field. We show that this effect is related to the back-action-driven excitation between the spin singlet–triplet (S–T) states. A finite magnetic field enables us to study the effect of back-action on individual triplet states for even electron numbers and on Zeeman states for odd numbers. By solving density matrix rate equations, we quantitatively determine the back-action-assisted spin excitation rates.

2. Experiment

We measure a GaAs/AlGaAs single QD with an integrated QPC in a helium-3 refrigerator at the base temperature of 240 mK. A scanning electron microscopy (SEM) image of the sample is shown in figure 1(a). All the ohmic contacts are grounded except that one side of the QPC is biased at a small dc voltage $V_{dc}$ to generate a current $I_{QPC}$. A tiny gap between the left top gate (LT) and the right top gate (RT) enhances the QPC sensitivity. We always close the left barrier of the QD and let the electrons tunnel through the right barrier. Figure 1(b) is a gray-scale plot of the numerical derivative of $I_{QPC}$, from which we can tell the tunneling of each individual electron. In this paper, we focus on the last six electrons. Figure 1(c) shows a random telegraph signal (RTS) in the real-time trace of $I_{QPC}$, reflecting the tunneling of the last electron on/off.
Figure 1. (a) An SEM image of the QD–QPC structure. (b) The gray-scale plot of the numerical derivative of $I_{QPC}$ with respect to $V_P$. The other gate voltages are $V_T ≡ V_{LT} = V_{RT} = −1.20\, V$, $V_{LB} = −1.40\, V$ and $V_Q = −1.25\, V$, $V_{dc} = 1\, mV$. (c) A typical trace of RTS when $V_P = −1.15\, V$ and $V_{RB} = −0.95\, V$, corresponding to the $0e ⇔ 1e$ tunneling.

the QD. The tunneling rate is conveniently controlled by the right barrier gate (RB) to around a few kHz. The detection bandwidth is limited to 30 kHz.

RTS is a thermally activated two-level switching behavior and is known to follow thermal-equilibrium Fermi statistics [11, 12]. There have been occasional reports about non-equilibrium effects in the RTS statistics [13–15], although no focused research on these effects has been carried out so far. Here we provide a systematic study. In our system, we can control the back-action strength with the QPC bias voltage $V_{dc}$ or the QD–QPC gap opening (modulated by voltage $V_T ≡ V_{LT} = V_{RT}$) [16]. Figure 2(a) shows the RTS statistics for the odd electron numbers $n = 5, 3$ and 1. $\mu_n − E_F$, the addition energy of the $n$th electron with respect to the Fermi level of the electron reservoir, is tuned with gate voltage $V_P$, whose lever arm is found to be $0.086\, meV\, mV^{-1}$ through transport experiments. The electron tunneling-out rate $\Gamma^{\text{out}}$ and tunneling-in rate $\Gamma^{\text{in}}$ describe the thermodynamics of RTS. All the odd electron numbers show the same saturation effect as we discussed earlier [16]. When $\mu_n$ is much lower than $E_F$, the $n$th electron is supposed to be trapped in the QD and $\Gamma^{\text{out}}$ should exponentially drop to zero. Instead, we see that $\Gamma^{\text{out}}$ saturates when $\mu_n − E_F ≪ 0$. With increasing $V_{dc}$ or $V_T$, a stronger saturation effect was observed. We attributed this to the back-action-assisted tunneling-out behavior and developed a phenomenological model to characterize this tunneling-out rate [16].

Supplementary figure S1 (available at stacks.iop.org/NJP/15/023021/mmedia) explains how we tune the tunneling rate for all the last six electrons to the same level while keeping the QD configuration unchanged.

We have carefully eliminated the leakage tunneling between the QD and QPC/left barrier, which is found to give rise to multi-step RTS or electron bunching due to different RTS amplitude/rate. We always observe an extra
Figure 2. RTS statistics of \((n - 1)e \leftrightarrow ne\) transitions. (a) Odd electron numbers \(n = 5, 3\) and \(1\). (b) Even electron numbers \(n = 6, 4\) and \(2\). The black closed dots are \(\Gamma^\text{out}\) and the red open circles are \(\Gamma^\text{in}\). Here \(V_T = -1.2\) V and \(V_{dc} = 1\) mV. For the even electron numbers, the red arrows point to additional features. The solid curves are the simulation result.

Analysis of the RTS statistics for even electron numbers reveals more information. Figure 2(b) shows the RTS statistics for \(n = 6, 4\) and \(2\). The saturation effect is also observed, just like the odd electron numbers. The difference is the dramatic additional features superimposed on the saturation tails, as indicated by red arrows. For instance, for \(n = 2\), \(\Gamma^\text{out}\) shows a strong side peak and \(\Gamma^\text{in}\) shows an elevated step at \(\mu_n - E_F = -0.80\) meV. For \(n = 4\), a pronounced extra feature is observed at \(-0.50\) meV. Another minor additional feature occurs at \(-0.90\) meV. For \(n = 6\), there are two likely extra features at \(-0.38\) and \(-0.62\) meV, although almost buried in noise.

These extra features grow stronger as we increase the back-action strength. Figure 3(a) shows the \(1e \leftrightarrow 2e\) RTS statistics under different back-action strengths. For the blue stars, red triangles and black circles, we increase \(V_{dc}\) or \(V_T\) and have seen increasingly higher back-action-induced saturation tails (best seen between \(-1.0\) and \(-1.2\) meV). This indicates a stronger back-action strength. What is more, the side peak in \(\Gamma^\text{out}\) (at \(-0.8\) meV) becomes stronger, implying that the observed extra features are relevant to the QPC back-action.

The dependence on the electron number and back-action strength strongly indicates that the extra features in the RTS statistics most likely arise from the back-action-induced excitation tunneling-out rate \(\Lambda^\text{out}\) and no extra tunneling-in rate \(\Lambda^\text{in}\) even if we switch the polarity or the connecting direction of \(V_{dc}\). This also excludes leakage as the origin of the non-equilibrium tunneling behavior.

New Journal of Physics 15 (2013) 023021 (http://www.njp.org/)
Figure 3. (a) RTS statistics of the $1e \leftrightarrow 2e$ transition under different back-action strengths. The closed shapes are $\Gamma_{\text{out}}$ and open shapes are $\Gamma_{\text{in}}$. Back-action from weak to strong: blue stars ($V_T = -1.5$ V and $V_{dc} = 1$ mV) $\rightarrow$ red triangles ($V_T = -1.5$ V and $V_{dc} = 2$ mV) $\rightarrow$ black circles ($V_T = -1.2$ V and $V_{dc} = 1$ mV). Solid curves are simulations. (b) Illustration of the back-action-driven excitation.

between the electron spin S–T states. At zero magnetic field, an even number of electrons form a spin S–T configuration with non-zero energy splitting, while an odd number of electrons leave a dangling spin whose two spin states are degenerate. Spin excited states also have a comparable relaxation rate ($1/T_1$ of the order of kHz for S–T [18]) with respect to the RTS tunneling rate here (about 1 kHz at the balance point) and should have a detectable effect. Orbital excited states, on the other hand, relax too fast ($1/T_1$ of the order of 100 MHz in GaAs QDs [17]) and we believe that they are invisible in this experiment. These can explain why we observe extra features for even electron numbers and not for odd numbers. Lastly, the energy scale of these extra features is comparable with the S–T energy splitting observed in biased transport experiments [19] and pulse-spectroscopy measurements [18] and also agrees with our own pulse-spectroscopy result on the same sample [20].

In this experiment, there is neither voltage bias on any of the QD contacts nor voltage pulse on any of the gates. Therefore, other than the QPC back-action, there seems to be no other source to excite the QD electrons to spin excited states. Certainly, the observed growth of this effect with back-action strength already serves as strong evidence.

3. Analysis

We use a master-equation model to quantitatively describe the influence of back-action on the excited energy states. In figure 3(b), we briefly illustrate the mechanism of back-action-driven excitation between two states. For $ne$, suppose that $\mu_n^1$ is the electron addition energy for the ground state and $\mu_n^2$ is that for an excited state. When $\mu_n^1 \ll E_F$, the $n$th electron is supposed to remain on $\mu_n^1$ and $\mu_n^2$ should be empty due to Coulomb repulsion. However, back-action drives the $n$th electron out of the QD (at a rate $\Lambda_{\text{out}}^n$) or up to $\mu_n^2$ (at a rate $\Lambda_{\text{in}}^n$), and initiates the tunneling between $\mu_n^2$ and the electron reservoir. Therefore an extra plateau in $\Gamma_{\text{in}}$ and a side peak in $\Gamma_{\text{out}}$ occur when $\mu_n^2$ is aligned with $E_F$. 
More generally, we consider the case of one \((n - 1)e\) state \(\mu_{n-1}\) and \(M\) \(ne\) states \(\mu^n_\alpha (\alpha = 1,\ldots,M)\). The occupancy of each state can be described by these rate equations [21]:

\[
\frac{d}{dt} P_i = \sum_j (\Gamma_{ij} - \delta_{ij} \sum_k \Gamma_{ki}) P_j.
\]

We always use Latin symbols \(i, j, k = 0, 1,\ldots,M\) and Greek symbols \(\alpha, \beta, \gamma = 1,\ldots,M\). \(P_0\) denotes the \((n - 1)e\) state occupancy and \(P_a\) denotes the occupancy of each \(ne\) state \(\mu^n_\alpha\). \(\Gamma_{ij}\) is the tunneling rate from state \(j\) to \(i\). Especially, \(\Gamma_{ii} = 0\) and \(\Gamma_{0\alpha} = \Gamma^\text{out}_\alpha (\Gamma_{\alpha0} = \Gamma^\text{in}_\alpha)\) is the tunneling-out (in) rate of state \(\mu^n_\alpha\). Because \(P_0 + \sum_a P_a = 1\), we can reduce the steady-state rate equations:

\[
\sum_{\beta} (X_{\alpha\beta} P_\beta^\prime) = Y_\alpha,
\]

\[
\begin{cases}
X_{\alpha\beta} = \Gamma^\text{in}_\alpha - \Gamma_{\alpha\beta} + \delta_{\alpha\beta}(\Gamma^\text{out}_\alpha + \sum_\gamma \Gamma_{\gamma\alpha}), \\
Y_\alpha = \Gamma^\text{in}_\alpha.
\end{cases}
\]

These equations can be easily solved if every rate \(\Gamma_{ij}\) is known. After considering the back-action-assisted tunneling, we set \(\Gamma_{ij}\) as the following [16]:

\[
\begin{align*}
\Gamma^\text{out}_\alpha &= g_{n-1}(\Lambda^\text{out}_\alpha + \Gamma^*_{\alpha}(1 - f(\mu^n_\alpha))), \\
\Gamma^\text{in}_\alpha &= g_n \Gamma^*_{\alpha} f(\mu^n_\alpha), \\
\Gamma_{\beta\alpha} &= \Lambda_{\beta\alpha}, \Gamma_{\alpha\beta} = \Lambda_{\beta\alpha} + 1/T^{\text{rel}}_{1\beta} (\beta > \alpha).
\end{align*}
\]

Here \(g_{n-1} = 2\) and \(g_n = 1\) are the spin degeneracy. For state \(\mu^n_\alpha\), \(\Gamma^*_{\alpha}\) is the thermal-equilibrium maximum tunneling rate and \(\Lambda^\text{out}_\alpha\) is the back-action-assisted tunneling-out rate. For \(\beta > \alpha\), \(\Lambda_{\beta\alpha}\) is the back-action-driven excitation rate from \(\mu^n_\alpha\) to \(\mu^n_\beta\), and \(1/T^{\text{rel}}_{1\beta}\) is the relaxation rate from \(\mu^n_\beta\) to \(\mu^n_\alpha\). We input these tunneling rates and relaxation rates \(\Gamma^*_{\alpha}, \Lambda^\text{out}_\alpha, \Lambda_{\beta\alpha}\) and \(1/T^{\text{rel}}_{1\beta}\) as free parameters to solve the steady-state rate equations. After obtaining each state occupancy \(P_\alpha\), we can simulate the RTS statistics according to

\[
\begin{align*}
\Gamma^\text{out}_\alpha &= \sum_\alpha \left( P_\alpha \Gamma^\text{out}_\alpha \right) / \sum_\alpha P_\alpha, \\
\Gamma^\text{in}_\alpha &= \sum_\alpha \Gamma^\text{in}_\alpha.
\end{align*}
\]

In figure 2(b), we show our simulation result as solid lines, which in general capture the basic features of the experimental data. For \(n = 2\), two spins form a ground state (singlet \(|S\rangle\)) and three energy-degenerate excited states (triplet \(|T^+\rangle, |T^0\rangle\) and \(|T^-\rangle\)). At zero magnetic field, we cannot distinguish the three triplet states and just denote the observed excited state as \(|T\rangle\) without losing generality. For \(n = 4\), we need to consider two spin excited states, since we observe two extra features. For \(n = 6\), the signal-to-noise ratio is too small for us to perform a faithful simulation due to the small S–T energy splitting and short relaxation time. The rapid decrease of the S–T energy splitting and relaxation time with increasing electron number, which is also seen in our pulse-spectroscopy measurements [20], might be an important factor to consider when choosing an appropriate electron number to implement spin S–T-based qubits.

We summarize the result for all even electron numbers in table 1. Here \(\Delta E_{1\alpha} = E^n_\alpha - E^n_1\) refers to the excited state energy with respect to the ground state. For comparison, we also list the charging energy \(E_C\) measured from figure 1(b) and the orbital level spacing \(\Delta\epsilon\) estimated through \(E_C\). If we assume \(\Delta E^{12}\) to be the spin S–T energy splitting \(\Delta E_{ST} = \Delta\epsilon - J\), with \(J\)
Table 1. Simulation results for $n = 2, 4$ and $6$.

| $n$ | $T_{12}$ (ms) | $T_{13}$ (ms) | $T_{23}$ (ms) | $\Delta E_{12}$ (meV) | $\Delta E_{13}$ (meV) | $E_C$ (meV) | $\Delta \epsilon$ (meV) |
|-----|---------------|---------------|---------------|------------------------|------------------------|-------------|------------------------|
| 6   | /             | /             | /             | 0.38                   | 0.62                   | 3.5         | 0.59                   |
| 4   | 1.2           | 0.025         | 1.0           | 0.50                   | 0.90                   | 4.5         | 0.98                   |
| 2   | 1.9           | /             | /             | 0.80                   | /                      | 6.6         | 2.1                    |

being the exchange energy), then $\Delta E_{13}$ should be the energy spacing between two neighboring orbital levels with the same spin configuration, i.e. $\Delta \epsilon$. This is verified by the good agreement between $\Delta E_{13}$ and the estimated $\Delta \epsilon$. This energy spectroscopy assignment is also consistent with the long relaxation times $T_{12}$ and $T_{23}$ which are between the spin S–T states, and the short relaxation time $T_{13}$ which is between orbital levels without spin flipping. We also note that the ratio of $\Delta E_{12} (= \Delta E_{ST})$ for $n = 6, 4$ and $2 \ (0.38 : 0.50 : 0.80 = 0.48 : 0.63 : 1)$ is in agreement with that of the charging energy $E_C \ (0.53:0.68:1)$. Because both of them originate from the Coulomb interaction whose potential energy linearly grows with the reciprocal of the QD size, their agreement implies that the observed increase of the S–T energy splitting with decreasing electron number is due to shrinking of the QD size.

Figure 3(a) presents a comparison between our simulation result and the controlled experimental study on the back-action strength. For the blue starts, red triangles and black circles, we increase the back-action strength by increasing either $V_{dc}$ or $V_T$. In this process, the thermal-equilibrium tunneling rate $\Gamma_{\alpha}^T$ remains at the same level: $\Gamma_{S}^T$ varies between 1.7 and 1.8 kHz, and $\Gamma_{T}^T$ changes from 12 to 14 kHz. The relaxation time $T_{ST}$ increases from 0.7 ms through 0.8 ms to 1.9 ms, showing a dependence on back-action that needs further study. What we are most concerned with are the back-action-induced tunneling rates. $\Lambda_{ST}^S$ is the back-action-induced tunneling-out rate and a characterization of the back-action strength [16]. It shows a substantial increase as we expected: from 0.5 Hz through 2 Hz to 15 Hz. $\Lambda_{TS}$, the back-action-driven excitation rate from spin singlet to triplet state, increases from 2 Hz through 10 Hz to 120 Hz. The dramatic growth of $\Lambda_{TS}$ with the back-action strength is evidence that the spin S–T excitation is driven by back-action.

4. Magnetic field dependence

The above experiments are performed at zero magnetic field. Now we apply a parallel magnetic field to lift the energy degeneracy between the two Zeeman states for the odd electron numbers, and between the three triplet states for the even electron numbers. Taking $1e$ and $2e$ as examples, we can distinguish the spin states and characterize the back-action-induced excitation rate to individual states. Figure 4(a) shows the total tunneling rate $\Gamma_{total} = \left[ (\Gamma_{in})^{-1} + (\Gamma_{out})^{-1} \right]^{-1}$ for $2e$ with different magnetic fields $B$. The black dashed line is at a low field, $B = 1 T$. Due to the strong back-action, we see both the ground state $|S\rangle$ and the excited states $|T\rangle$, located at $\mu_2 - E_F = 0$ and $-0.80$ meV, respectively. As we apply a large magnetic field, say, 8 T, the peak at $-0.80$ meV splits into two, with one shifting to the positive energy side (possibly $|T^+\rangle$) and the other shifting very little (possibly $|T^0\rangle$). Figure 4(b) is a color-scale plot of the numerical derivative of $\Gamma_{in}$. The splitting of the $|T\rangle$ state with increasing magnetic field is clearly seen.
Figure 4. (a) $\Gamma_{\text{total}}$ for $2e$ at 1 T (black dashed line) and 8 T (red straight line), respectively. (b) A color-scale plot of the numerical derivative of $\Gamma^{\text{in}}$ with respect to the electron energy, at different magnetic fields. The three green dashed lines indicate the suspected $|T^0\rangle$, $|T^+\rangle$, and $|S\rangle$ states. (c) Extracted energy between the suspected $|T^0\rangle$ and $|T^+\rangle$ states as a function of magnetic field. The solid line is the predicted Zeeman splitting assuming $g = -0.4$ for GaAs. (d) $\Gamma_{\text{total}}$ for $1e$ at 0 T (black dashed line) and 7 T (red straight line), respectively. (e) A color-scale plot of the numerical derivative of $\Gamma^{\text{in}}$ for $1e$. The two green dashed lines indicate the suspected spin-up and spin-down states. (f) Extracted energy between the suspected two spin states. The solid line is the predicted Zeeman splitting.
In figure 4(c), we explicitly extract the energy splitting as a function of magnetic field and see good agreement with the predicted Zeeman splitting \( g \mu_B B \) using a \( g \)-factor \(-0.4\) in GaAs. This quantitatively proves that the splitting is between the \(|T^0\rangle\) and \(|T^+\rangle\) states.

We note that the \(|T^0\rangle\) peak, in general, is much weaker than the \(|T^+\rangle\) peak. From figure 4(a), our simulation shows that \( \Lambda_{T^+S} = 600 \text{ Hz} \) and \( \Lambda_{T^0S} = 100 \text{ Hz} \) at 8 T. The excitation rate from \(|S\rangle\) to \(|T^+\rangle\) is much stronger than to \(|T^0\rangle\). It is usually assumed that due to the spin–phonon selection rules, the \(|T^0\rangle–|S\rangle\) interaction is suppressed in the lowest order, and as a result, the relaxation rate from \(|T^0\rangle\) to \(|S\rangle\) is slower than that from the other two triplet states [22]. This could, in turn, explain why the phonon-mediated back-action between \(|T^0\rangle\) and \(|S\rangle\) is weaker and why \( \Lambda_{T^0S} \ll \Lambda_{T^+S} \). Our experiment, therefore, establishes a direct comparison between the \(|T^0\rangle–|S\rangle\) and \(|T^+\rangle–|S\rangle\) coupling strengths.

Similarly, figures 4(d)–(f) show the effect of magnetic field for 1e. Apart from the major peak at \( \mu_2 - E_2 = 0 \text{ meV} \) (labeled as \(|\uparrow\rangle\)), a side peak in \( \Gamma_{\text{total}} \) or \( \Gamma_\text{in} \) appears at a large magnetic field (labeled as \(|\downarrow\rangle\)). The side peak is barely visible at a small field, possibly due to the large thermal energy (\( T = 240 \text{ mK} \)). Nonetheless, the extracted energy splitting between these two peaks at a large field (\( \geq 6.5 \text{ T} \)) agrees well with the predicted Zeeman splitting. Therefore, we believe that they are the spin-up and spin-down states. Using the hidden Markov method [14] to extract more information from the RTS statistics, we recently found that we can distinguish the two spin states of a single electron even at a magnetic field as low as 1 T, when the Zeeman splitting is much smaller than the thermal energy [23].

5. Conclusion

In conclusion, we observe real-time back-action-driven excitation between the spin states in a few-electron single QD. The spin excitation rate is obtained from the analysis of the charge counting statistics. Considering the potential application of few-electron spin states in quantum computation, this would be useful in evaluating the influence of back-action on the operation of qubits.

Acknowledgment

This work was supported by the NFRP (2011CBA00200 and 2011CB921200) and NNSF (10934006, 11074243, 11174267, 91121014 and 60921091).

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