Nonequilibrium current and noise in inelastic tunneling through a magnetic atom

Björn Sothmann\textsuperscript{1} and Jürgen König\textsuperscript{1}

Theoretische Physik, Universität Duisburg-Essen and CeNIDE, 47048 Duisburg, Germany
E-mail: bjoerns@thp.uni-due.de and koenig@thp.uni-due.de

New Journal of Physics 12 (2010) 083028 (9pp)
Received 8 June 2010
Published 11 August 2010
Online at http://www.njp.org/
doi:10.1088/1367-2630/12/8/083028

Abstract. In a recent experiment, Hirjibehedin et al (2007 Science 317 1199) performed inelastic tunneling spectroscopy of a single iron atom absorbed on a nonmagnetic substrate. The observed steps in the differential conductance marked the spin excitation energies. In this paper, we explain the observed nonmonotonicities in the differential conductance by a nonequilibrium population of the atom spin states. Furthermore, we predict super-Poissonian current noise due to this nonequilibrium situation. We argue that the remarkable absence of nonequilibrium features at certain conductance steps indicates the presence of an anisotropic relaxation channel.

Contents

1. Introduction 2
2. Model 2
3. Results 4
4. Conclusions 8
Acknowledgment 9
References 9

\textsuperscript{1} Authors to whom any correspondence should be addressed.
1. Introduction

Inelastic scattering processes provide a convenient tool to study the excitation spectra of various systems. By using inelastic electron tunneling spectroscopy, one can access vibrational excitations of ensembles of molecules in metallic tunnel junctions [1, 2] or single molecules in scanning tunneling microscope (STM) geometries [3]. The excitation energies reveal themselves as steps in the differential conductance whenever a new inelastic transport channel opens up. For the explanation of the measured signals, an equilibrium distribution of the molecule states was implicitly assumed. Further studies of molecular vibrations were performed using H₂ molecules [4], C₆₀ molecules [5] in mechanical break junctions and suspended carbon nanotubes [6]–[8].

Recently, the investigation of magnetic properties and interactions on an atomic level became possible due to the advent of spin inelastic electron spectroscopy [9]–[14]. Here, single magnetic atoms absorbed onto a nonmagnetic substrate were contacted using an STM tip. Describing the atom in terms of a localized spin, Hirjibehedin et al. [12] related the positions of the conductance steps to the energy associated with transitions between different eigenstates, while the relative step heights depend on the matrix elements of the spin operator. A more complete theoretical description based on perturbation theory in the tunnel coupling [15]–[19] still relies on the assumption of equilibrium occupations.

While the above studies could explain the conductance steps assuming the atom spin to be in thermal equilibrium with the substrate, nonmonotonic features clearly present in the experimental results of [12] were not addressed. Conductance overshoots due to nonequilibrium occupations together with their relaxation by spin–phonon interactions have been discussed in [20] for cotunneling through a quantum dot. A similar behavior was found in [21]–[23], where additionally the low-temperature nonequilibrium logarithmic Kondo enhancement of this overshoot was studied. While the Kondo effect is relevant for transport through a single Co atom studied in [13, 14], it is not important here. Nonequilibrium effects have been considered in [19] for spin-transfer torque on a single atom coupled to ferromagnetic substrates and tips. In [24], the nonequilibrium current and current noise through a single molecular magnet were analyzed in the charge fluctuation regime. In this paper, we explain the experimental results by calculating the nonequilibrium occupations together with a spin-dependent relaxation channel using a master-equation approach. We, furthermore, predict an enhanced Fano factor indicating super-Poissonian current noise as a clear sign of a nonequilibrium situation.

2. Model

We model the experimental setup of [12] as two reservoirs of noninteracting electrons coupled by a tunnel barrier with an embedded spin. Hence, the Hamiltonian describing the system is given by

\[ H = \sum_r H_r + H_{\text{spin}} + H_{\text{tun}}. \]  

Here, \( H_r = \sum_{k\sigma} \varepsilon_{rk} a_{r,k\sigma}^\dagger a_{r,k\sigma} \) models the two electrodes as reservoirs of noninteracting electrons with constant density of states \( \rho_r \) and electrochemical potential \( \mu_r \). The operator \( a_{r,k\sigma}^\dagger \) creates an
electron in lead \( r = \text{L}, \text{R} \) with momentum \( \mathbf{k} \) and spin \( \sigma \). The local spin is described by
\[
H_{\text{spin}} = -DS_z^2 + E(S_x^2 - S_y^2) + g \mu_B \mathbf{B} \cdot \mathbf{S},
\]
where the \( z \)-axis is the magnetic easy axis of the atom in its coordination environment. For an \( S = 2 \) iron atom on Cu$_2$N, the best fit to the experimental results in [12] gives a uniaxial anisotropy \( D = 1.55 \) meV, a transverse anisotropy \( E = 0.31 \) meV and a \( g \)-factor of \( g = 2.11 \).

In Table 1, we summarize the eigenenergies and eigenstates of the spin Hamiltonian for two different choices of the external magnetic field. Finally, the tunneling Hamiltonian is given by the Appelbaum Hamiltonian [25],
\[
H_{\text{tun}} = \sum_{rr'kk'\sigma\sigma'} j_{rr'} a_{r\sigma}^\dagger \sigma_{\sigma'\cdot} \mathbf{S} a_{r'k'\sigma'},
\]
with \( \sigma \) denoting the Pauli matrices, which describes an exchange interaction between the spin of the tunneling electron and the local spin. We neglect direct tunneling through the barrier not involving the localized spin as it only gives rise to a bias-independent elastic background to the differential conductance. Interference terms between direct and exchange tunneling do not appear in the total current and shot noise since the contributions from spin-up and -down electrons cancel each other out for nonmagnetic electrodes. The above model has been studied extensively to describe molecular magnets. The Kondo effect induced by transverse anisotropy [26]–[28], as well as Berry phase effects [29, 30] and the current-induced switching of the molecule spin [31], has been discussed.

We parametrize the couplings \( j_{rr'} \) through the sum \( J = j_{\text{LL}} + j_{\text{RR}} \) and the asymmetry \( a = (j_{\text{LL}} - j_{\text{RR}})/(j_{\text{LL}} + j_{\text{RR}}) \), i.e. \( j_{\text{LL}}^2 = (1 + a)^2 J^2/4 \), \( j_{\text{RR}}^2 = (1 - a)^2 J^2/4 \) and \( j_{\text{LR}}^2 = j_{\text{RL}}^2 = (1 - a^2) J^2/4 \). While the couplings \( j_{\text{LR}} \) and \( j_{\text{RL}} \) are responsible for the current through the atom, which may be accompanied by a spin excitation or disexcitation, the couplings \( j_{\text{LL}} \) and \( j_{\text{RR}} \) do not contribute to the current but give rise to a transport-induced relaxation mechanism for the local spin only.

Table 1. Eigenenergies \( E_m \) and eigenstates \( |m \rangle \) of the spin Hamiltonian (2) in the basis \( |S_z \rangle \) of the \( S_z \) eigenstates for a magnetic field applied in the \( z \) and the \( x \) direction, respectively.

| \( B_z \) (T) | \( E_m \) (meV) | \( |2 \rangle_z \) | \( |1 \rangle_z \) | \( |0 \rangle_z \) | \( |-1 \rangle_z \) | \( |-2 \rangle_z \) |
|------------|-------------|-------------|-------------|-------------|-------------|-------------|
| \( B_z = 7 \) T | \( B_z = 3 \) T |
| 0 | \(-7.982\) 0.021 0 \(-0.097\) 0 \(0.995\) |
| 1 | \(-4.612\) 0.987 0 \(-0.157\) 0 \(-0.036\) |
| 2 | \(-2.813\) 0 0.402 0 \(-0.916\) 0 |
| 3 | \(-0.287\) 0 0.916 0 0.402 0 |
| (4) | 0.194 0.159 0 0.983 0 0.092 |

New Journal of Physics 12 (2010) 083028 (http://www.njp.org/)
The dynamics of the system are governed by a generalized master equation for the probabilities $P_m$ to find the spin in one of its eigenstates $|m\rangle$ with energy $E_m$,\[ \frac{dP_m}{dt}(t) = \sum_{m'} \int_{-\infty}^{t} dt' L_{mm'}(t-t')P_{m'}(t'). \] (4)

In the stationary limit, $P_m \equiv P_m(t)$ is independent of time and only the time-integrated kernel $L_{mm'} \equiv \int_{-\infty}^{0} dt' L_{mm'}(-t')$ is needed. Its matrix elements are $L_{mm'} = W_{mm'} - \delta_{mm'} W_m$. Here, $W_{mm'}$ are the Fermi’s golden rule transition rates,

$$ W_{mm'} = \sum_{rr'\alpha} 2\pi |j_{rr'}|^2 \rho_r \rho_{r'} |\langle m | S_\alpha | m' \rangle|^2 \zeta(\mu_r - \mu_{r'} - \Delta_{mm'}) , \tag{5} $$

where $\zeta(x) = x/(1 - e^{-x/(k_BT)})$ and $\Delta_{mm'} = E_m - E_{m'}$. The sum runs over the lead indices $r, r' = L, R$ and the spin directions $\alpha = x, y, z$. The elements $W_m$ follow from $\sum_m L_{mm'} = 0$, which guarantees the conservation of probability.

To compute the current and current noise, we employ the formalism of full-counting statistics adopted to a system that can be described by rate equations [32, 33]. To this end, we introduce the matrix $W^\chi_{mm'}$, which is obtained from $W_{mm'}$ by multiplying each term in the sum (5) by a factor $e^{i\chi}$ if $r = L, r' = R$, $e^{-i\chi}$ if $r = R, r' = L$, and 0 otherwise, where $\chi$ is called a counting field. Furthermore, we define $L^\chi_{mm'} = W^\chi_{mm'} - \delta_{mm'} W_m$ (note that $W_m$ does not contain the counting field $\chi$). The smallest eigenvalue of $L^\chi_{mm'}$ defines the cumulant generating function $S(\chi)$, from which we can obtain the average current $I$ and the current noise $S$ by performing derivatives with respect to the counting field, $I = -ie(dS(\chi)/d\chi)|_{\chi=0}$ and $S = (-ie)^2(d^2S(\chi)/d\chi^2)|_{\chi=0}$.

Although the full-counting statistics formalism to compute the current and noise is very compact and elegant for the calculation, we introduce, in addition, an equivalent formulation for the average current, which offers a more transparent basis for distinguishing equilibrium from nonequilibrium effects. It is easy to show that the average current can also be written as

$$ I = -ie \sum_{m,m'} \frac{d}{d\chi} W^\chi_{mm'} \bigg|_{\chi=0} P_{m'} . \tag{6} $$

The derivatives $i(dW^\chi_{mm'}/d\chi)|_{\chi=0}$ are the current rates. Nonequilibrium effects of the current–voltage characteristics enter via the nonequilibrium probability distribution $P_m$, which is obtained by solving the master equation, equation (4). These nonequilibrium effects would be neglected if one replaced the $P_m$ with an equilibrium probability distribution, $P_m^{\text{eq}} = \exp(-E_m/k_BT)/\sum_{m'} \exp(-E_{m'}/k_BT)$, i.e. for low temperature $P_0 = 1$ for the ground state and $P_m = 0$ for the excited states $m \neq 0$.

3. Results

In the following, we discuss the influence of a nonequilibrium spin occupation on the transport properties for the system parameters of the experiment [12]. In figure 1, we show the differential conductance in the presence of a strong magnetic field $B_z = 7$ T along the easy axis for different values of the asymmetry parameter $a$ in the absence of the phenomenological relaxation (12) (see the discussion below). For very large asymmetries, $a \rightarrow 1$, there are flat plateaus between the conductance steps. In this limit, the coupling constant $j_{1R}$ for processes that drive the atom state population out of equilibrium is much smaller than $j_{1L}$ for processes that let the...
system relax to thermal equilibrium with the left electrode. Therefore, as in [1], nonequilibrium effects are absent, and the resulting conductance curve is identical to the one obtained in [15]. For smaller asymmetries, the situation is different. The height of the conductance steps at the excitation thresholds is increased. Beyond the threshold voltages, the differential conductance shows a slow power-law decay towards its value for the equilibrated system again. This overshooting behavior is observed for most of the steps in the experiments of [12]. While the coupling of the adatom to the substrate is fixed in experiment, the coupling to the STM tip can be controlled by changing the tip–atom distance. This corresponds to changing the total coupling $J$ and thereby the total tunnel current as well as the asymmetry $a$ and thereby the nonequilibrium effects. In a recent experiment using a magnetic STM tip [34], it was confirmed that, by decreasing the tip–atom distance and therefore increasing the current through the system, the nonequilibrium effects became more pronounced.

Before we discuss this for the system at hand, we illustrate the mechanism that leads to this conductance behavior explicitly for the simpler model of a local spin-$1/2$ with Zeeman energy $B_z$ symmetrically coupled to the electrodes, at zero temperature. Transport takes place by either spin-flip or spin-conserving transitions. The latter contribute to the current as

$$I_{sc} = \pi e |j_{LR}|^2 \rho_L \rho_R eV,$$

independent of the probabilities $P_\uparrow$ and $P_\downarrow$ to find the spin in state up and down, respectively. The differential conductance, measured in units of $G_0 = 4\pi e^2 S(S+1) |j_{LR}|^2 \rho_L \rho_R$, is $G_{sc} = G_0/3$. Therefore, the nonequilibrium population of the spin states is only probed by the spin-flip processes. They contribute for $eV \geq B$ as

$$I_{sf} = 2\pi e |j_{LR}|^2 \rho_L \rho_R \left[(eV - B) P_\uparrow + (eV + B) P_\downarrow\right].$$

In equilibrium, only the ground state is occupied, $P_\uparrow = 1$ and $P_\downarrow = 0$, such that only the first term in (8) contributes. Hence, the differential conductance $G_{sf}^{eq} = 2G_0/3$ remains constant.
above threshold. In the nonequilibrium situation, the occupation probabilities are obtained from the master equation (4) in the stationary state,

\[ 0 = \frac{d}{dt} \left( \frac{P_\uparrow}{P_\downarrow} \right) = 2\pi e|j_{LR}|^2 \rho_L \rho_R \left( \frac{-(eV - B)}{eV - B} \frac{eV + B}{-(eV + B)} \right) \left( \frac{P_\uparrow}{P_\downarrow} \right). \]  

(9)

The solution is \( P_\uparrow = 1 - P_\downarrow = 1 - \frac{eV - B}{2(eV + B)}. \) As a consequence, now both terms in (8) contribute, leading to the total conductance (above threshold)

\[ G = \frac{2}{3} G_0 \left( 1 + \frac{2B^2}{(eV + B)^2} \right). \]  

(10)

In the limit \( V \to \infty, \) both \( P_\uparrow, P_\downarrow \to 1/2 \) and the conductance approach the equilibrium value \( G_0 \) with a power law on voltage scale \( B \) (although the probability distribution remains highly nonequilibrium).

For the \( S = 2 \) spin of the iron atom with its more complicated spin Hamiltonian, the same mechanism as in the simpler spin-1/2 model gives rise to the enhanced conductance in the nonequilibrium situation. While in equilibrium only the ground state is occupied, \( P_0 = 1, \) leading to steps in the differential conductance, in the nonequilibrium case we obtain bias-dependent occupations by solving the master equation (4), which lead to an overshooting. Again, the conductance decreases above threshold to approach its equilibrium value \( G_0 \) in the limit of infinite bias voltage. From our analysis it is clear that the nonmonotonic differential conductance is due to an increase in transport enabled by the population of excited states above threshold but close to the step. It is not a signature of the excited spin states carrying less current than the ground state, i.e. a decrease in the conductance, as has been speculated in [19]. This conclusion can be checked experimentally by measuring the Fano factor, i.e. the ratio between current noise and average current, \( F = S/(eI), \) as we now explain.

The Fano factor \( F \) is shown in figure 2 for different values of the asymmetry parameter \( a, \) not taking into account the relaxation term (12) (see below). For \( a \to 1, \) we find \( F = 1, \) i.e. Poissonian behavior, as expected for transport through a normal tunnel barrier. When a nonequilibrium population of the atom spin states becomes important \((a < 1)\) and bias voltage exceeding the inelastic threshold, the Fano factor becomes super-Poissonian, reaches a maximum and then slowly drops towards the Poissonian limit for large bias (the transition \(|0\rangle \to |1\rangle \) hardly gives rise to super-Poissonian current noise as it is a very weak excitation; cf figure 1). The latter behavior is an indicator that the nonmonotonic differential conductance is not due to a smaller current contribution from the excited states. If this was the case, we would expect a random telegraph signal with a super-Poissonian Fano factor for \( V \to \infty. \)

The mechanism leading to the super-Poissonian noise for bias voltages above the inelastic thresholds can most easily be understood by considering the spin-1/2 model again. In this case, the Fano factor above threshold is

\[ F = 1 + \frac{2B^2}{(eV + B)^2} \cdot \frac{3(eV - B)^2 + 8(eV - B)B}{3(eV - B)^2 + 9(eV - B)B + 2B^2}. \]  

(11)

Spin-conserving tunneling processes are stochastically independent of each other and of the spin-flip transitions. They obey Poissonian statistics and can be ignored for the following discussion. Once the inelastic transport channel is open, spin-flip transitions set in. They lead to an alternating sequence of the spin being in the ground and the excited state. In the limit of large bias voltage, the rates for the spin-flip transitions \( \uparrow \to \downarrow \) and \( \downarrow \to \uparrow \) become equal, and the transport statistics become Poissonian. For voltages just above threshold, \( eV \geq B, \) however, the
Figure 2. Fano factor $F$ and $dF/dV$ as a function of bias voltage for different values of the asymmetry parameter $a$. For $eV/k_B T \to 0$ the Fano factor diverges due to thermal noise. Parameters as in figure 1.

two spin-flip rates differ from each other. As a consequence, we obtain an alternating sequence of a longer and a shorter waiting time, i.e. effectively there is a tendency for two electrons to bunch together, which yields super-Poissonian current noise.

While our theory predicts an overshooting of the differential conductance at all conductance steps, in the experiment of [12] this feature is absent for the steps associated with the transition between the ground state $|0\rangle$ and the first excited state $|1\rangle$ whenever these steps are pronounced, as is the case for a magnetic field along the $x$-axis. This indicates that some relaxation mechanism reduces the occupation of $|1\rangle$. We note that the transition matrix element of $S_z$ between the ground state and the first excited state is large compared to matrix elements of $S_x$ and $S_y$, as well as compared to matrix elements of $S_z$ between the ground state and any other excited state. This observation is not very sensitive to the direction and the strength of the applied magnetic field. Therefore, we make the ad hoc assumption that there is an additional spin relaxation channel that couples to the $z$-component of the local spin only.

We add to our master equation (4) the following phenomenological, spin-dependent relaxation rates,

$$W^\text{relax}_{mm'} = -\frac{|\langle m | S_z | m' \rangle|^2}{\tau} \Theta(\Delta_{m'm}),$$

for $m \neq m'$, where $\Theta(x)$ is the step function and $\tau$ is the time scale for relaxation. The energy dependence in (12) is not crucial for our conclusions. We therefore choose the simplest possible ansatz that allows relaxation only into states with lower energy. In contrast, the spin matrix elements are crucial as they suppress the nonequilibrium effects for the transition between ground and first excited state while leaving them unaffected for almost all the other transitions.

New Journal of Physics 12 (2010) 083028 (http://www.njp.org/)
Figure 3. Differential conductance taking into account a spin-dependent relaxation mechanism of the form (12). The relaxation time is given in units of $\tau_0$ with $\tau_0^{-1} = 2\pi DS^2 |j_{LR}|^2 \rho_L \rho_R$. Parameters are $B_x = 3$ T, $T = 0.5$ K and $a = 0$. The corresponding eigenenergies and eigenstates are summarized in table 1.

In figure 3 we plot the differential conductance in the presence of a magnetic field $B_x = 3$ T along the $x$-direction. The first transition is more pronounced than in figure 1, where $B_z = 7$ T along the $z$-direction. In the limit $\tau \to \infty$, we recover the situation discussed above where an overshooting effect can be observed for each conductance step. By choosing a finite value for the relaxation time comparable to the cotunneling rates exciting the system, we can, however, eliminate the overshooting at the first transition while leaving the remaining part of the conductance curve practically unaffected. In the limit $\tau = 0$, we recover the equilibrium value for the conductance at each step. These results are not sensitive to the choice of the size and direction of the magnetic field.

The anisotropic relaxation can also explain the absence of conductance steps due to transitions between excited states in the experiment. Such features should be present for a small magnetic field applied in the $z$-direction, as in this case the first excited state gets populated significantly at the first conductance step, the excitation energies satisfy $\Delta_{21} < \Delta_{20}$ such that the transition $|1\rangle \to |2\rangle$ occurs before the onset of the transition $|0\rangle \to |2\rangle$ and furthermore the transition matrix elements $\langle 2 | S_\alpha | 1 \rangle$ do not all vanish. However, as the relaxation prevents the system from populating the first excited state, these additional conductance features vanish together with the overshooting at the first step.

4. Conclusions

We have investigated the nonequilibrium effects in transport through a single iron atom. With our model, we were able to explain the nonmonotonic features of the differential conductance observed experimentally in [12]. Furthermore, we noted a striking absence of this nonmonotonicity at certain conductance steps, which can be explained by the presence of an anisotropic spin relaxation channel. The anisotropy [35] points at the importance of spin–orbit coupling in this process, in addition to the splitting of the multiplet. In addition, we predicted the occurrence of super-Poissonian current noise as a consequence of the nonequilibrium
spin occupations probabilities. In conclusion, for a full understanding of inelastic tunneling spectroscopy, it is crucial to account for nonequilibrium populations of the atom states established by the competition of transport and anisotropic relaxation.

Acknowledgment

We acknowledge financial support from DFG via SFB 491.

References

[1] Jaklevic R C and Lambe J 1966 Phys. Rev. Lett. 17 1139
[2] Scalapino D J and Marcus S M 1967 Phys. Rev. Lett. 18 459
[3] Stipe B C, Rezaei M A and Ho W 1998 Science 280 1732
[4] Smit R H M, Noat Y, Untiedt C, Lang N D, van Hemert M C and van Ruitenbeek J M 2002 Nature 419 906
[5] Park H, Park J, Lim A K L, Anderson E H, Alivisatos A P and McEuen P L 2000 Nature 407 57
[6] LeRoy B J, Lemay S G, Kong J and Dekker C 2004 Nature 432 371
[7] Leturcq R, Stampfer C, Inderbitzin K, Durrer L, Hierold C, Mariani E, Schultz M G, von Oppen F and Ensslin K 2009 Nat. Phys. 5 327
[8] Hüttel A K, Witkamp B, Leijnse M, Wegewijs M R and van der Zant H S J 2009 Phys. Rev. Lett. 102 225501
[9] Heinrich A J, Gupta J A, Lutz C P and Eigler D M 2004 Science 306 466
[10] Hirjibehedin C F, Lutz C P and Heinrich A J 2006 Science 312 1021
[11] Meier F, Zhou L, Wiebe J and Wiesendanger R 2008 Science 320 82
[12] Hirjibehedin C F, Lin C, Otte A F, Ternes M, Lutz C P, Jones B A and Heinrich A J 2007 Science 317 1199
[13] Otte A F, Ternes M, von Bergmann K, Loth S, Brune H, Lutz C P, Hirjibehedin C F and Heinrich A J 2008 Nat. Phys. 4 847
[14] Otte A F, Ternes M, Loth S, Lutz C P, Hirjibehedin C F and Heinrich A J 2009 Phys. Rev. Lett. 103 107203
[15] Fernández-Rossier J 2009 Phys. Rev. Lett. 102 256802
[16] Fransson J 2009 Nano Lett. 9 2414
[17] Persson M 2009 Phys. Rev. Lett. 103 050801
[18] Lorente N and Gauyaçaq J 2009 Phys. Rev. Lett. 103 176601
[19] Delgado F, Palacios J J and Fernández-Rossier J 2010 Phys. Rev. Lett. 104 026601
[20] Lehmann J and Loss D 2006 Phys. Rev. B 73 045328
[21] Rosch A, Paaske J, Kroha J and Wölfle P 2003 Phys. Rev. Lett. 90 076804
[22] Paaske J, Rosch A, Kroha J and Wölfle P 2004 Phys. Rev. B 70 155301
[23] Schoeller H and Reininghaus F 2009 Phys. Rev. B 80 045117
[24] Romeike C, Wegewijs M R and Schoeller H 2006 Phys. Rev. Lett. 96 196805
[25] Appelbaum J 1966 Phys. Rev. Lett. 17 91
[26] Romeike C, Wegewijs M R, Hofstetter W and Schoeller H 2006 Phys. Rev. Lett. 96 196601
[27] Romeike C, Wegewijs M R, Hofstetter W and Schoeller H 2006 Phys. Rev. Lett. 97 206601
[28] Wegewijs M R, Romeike C, Schoeller H and Hofstetter W 2007 New J. Phys. 9 344
[29] Leuenberger M N and Mucciolo E R 2006 Phys. Rev. Lett. 97 126601
[30] González G and Leuenberger M N 2007 Phys. Rev. Lett. 98 256804
[31] Misiryan M and Barnaś J 2007 Europhys. Lett. 78 27003
[32] Bagrets D A and Nazarov Y V 2003 Phys. Rev. B 67 085316
[33] Braggio A, König J and Fazio R 2006 Phys. Rev. Lett. 96 026805
[34] Loth S, von Bergmann K, Ternes M, Otte A F, Lutz C P and Heinrich A J 2010 Nat. Phys. 6 340
[35] Leuenberger M N and Loss D 2000 Phys. Rev. B 61 1286

New Journal of Physics 12 (2010) 083028 (http://www.njp.org/)