Detection of Single Spin Decoherence in a Quantum Dot via Charge Currents

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We consider a quantum dot attached to leads in the Coulomb blockade regime which has a spin \( 1/2 \) ground state. We show that by applying an ESR field to the dot-spin the stationary current in the sequential tunneling regime exhibits a resonance whose line width is determined by the single-spin decoherence time \( T_2 \). The Rabi oscillations of the dot-spin are shown to induce coherent current oscillations from which \( T_2 \) can be deduced in the time domain. We describe a spin-inverter which can be used to pump current through a double-dot via spin flips generated by ESR.

An increasing number of spin-related experiments show that the electron spin is a robust candidate for coherent quantum state engineering in solid state systems such as semiconductor nanostructures. Several techniques, most prominently electron spin resonance (ESR), can then be envisaged for manipulation of electron spins on quantum dots, where the coherence of the spin is limited by the intrinsic spin decoherence time \( T_2 \). In some related systems, time-resolved optical measurements were used to determine \( T_2 \), the decoherence time of many of spins, with \( T_2 \) exceeding 100 ns in bulk GaAs. More recently, the single spin relaxation time \( T_1 \) (generally \( T_1 \geq T_2 \)) of a single quantum dot attached to leads was measured via transport to be longer than a few \( \mu s \), consistent with calculations. In this work, we go one step further and propose a setup to extract the single spin decoherence time \( T_2 \) of an electron confined in a quantum dot from transport measurements. The dot, which is attached to leads, is operated in the Coulomb blockade regime, and the spin flips generated by an ESR source lead to a resonance in the stationary charge current with a line width determined by the spin decoherence time, see Figs. 1, 2. Making use of coupled master equations we analyze the time-dependence of the current and the spin-measurement process (read-out), and show that coherent Rabi oscillations of the dot-spin induce oscillations of the current, providing a measure of the spin decoherence directly in time space, see Fig. 3. In the absence of a bias, the current can be pumped through a double-dot with the ESR source (providing the necessary energy via spin flips on the dot) and by making use of a novel spin-inverter for producing spin-dependent tunneling.

Model. We study a quantum dot in Coulomb blockade regime, coupled to two Fermi-liquid leads \( l = 1, 2 \) at chemical potentials \( \mu \). We consider the Hamiltonian

\[
H = H_0 + H_T = H_{\text{lead}} + H_{\text{dot}} + H_T,
\]

which describes leads, dot and the tunnel coupling between leads and dot, resp. Here, \( H_{\text{dot}} = H_{\text{ESR}} + H_{\text{Zeeman}} \), where \( H_{\text{dot}} \) includes charging and interaction energies of the electrons on the dot. \( H_{\text{dot}} \) also contains a Zeeman coupling term to a constant magnetic field \( B_z \) in \( z \)-direction, \(-\frac{1}{2} \Delta_z \sigma_z \), with Zeeman splitting \( \Delta_z = g\mu_B B_z \), electron \( g \) factor \( g \), Bohr magneton \( \mu_B \), and Pauli matrix \( \sigma_z \). Coupling to an oscillating magnetic ESR field in \( x \)-direction of frequency \( \omega \) is included in \( H_{\text{ESR}} = -\frac{\omega}{2} \Delta_z \cos(\omega t) \sigma_x \), with \( \Delta_z = g\mu_B B_z \) and Pauli matrix \( \sigma_x \). Such an oscillating field produces Rabi spin-flips at \( \omega = \Delta_z \), as used in ESR. We assume Zeeman splitting of the leads \( \Delta_{\text{leads}} \neq \Delta_z \) and \( \Delta_{\text{leads}} \ll \varepsilon_F \), where \( \varepsilon_F \) is the Fermi energy, such that the field effects of \( B_z \) and \( B_z(t) \), resp., on the leads are negligible. Such a situation can be achieved by using materials of different \( g \) factors and/or with local magnetic fields.

![Detection of Single Spin Decoherence in a Quantum Dot via Charge Currents](image)

We now describe the electronic states of the dot. For an odd number of electrons on the dot with antiferromagnetic filling, the topmost (excess) electron can be either in the spin ground state, \( |\uparrow\rangle \), or in the excited state, \( |\downarrow\rangle \), see Fig. 1. For an additional electron on the dot, we assume the ground state to be the singlet \( |S\rangle = \frac{1}{\sqrt{2}}(|\uparrow\downarrow\rangle - |\downarrow\uparrow\rangle) \) (which can be achieved by tuning \( B_z \)). The energy of the dot is defined by \( H_{\text{dot}} |n\rangle = E_n |n\rangle \). It is convenient to use only one-particle energies \( \Delta_{S\uparrow(S\downarrow)} = E_S - E_{S\uparrow(S\downarrow)} \) (con-
taining charging energy $U$), which can then be compared with $\mu$.

**Master equation.** We derive the master equation for the reduced density matrix of the dot, $\rho_D = T_L \rho$. Here, $T_L$ is the trace taken over the leads and $\rho$ is the full density matrix. Using a superoperator formalism, we evaluate the von Neumann equation within Born approximation in $H_T$ while taking $H_{ESR}$ fully into account. Hereby we make the usual assumption that the correlations induced in the leads by $H_T$ decay rapidly (Markovian approximation) and that we can neglect non-secular terms [14]. We obtain the following master equation

$$\dot{\rho}_t = -(W_{\uparrow\downarrow} + W_{S\uparrow}) \rho_t + W_{\uparrow\downarrow} \rho_t + W_{\downarrow S} \rho_S + W_{\downarrow S} \rho_S + \Delta_x \cos(\omega t) |\langle \rho_{\downarrow S}|t\rangle|,$$

$$\dot{\rho}_l = W_{\uparrow l} \rho_t - (W_{\uparrow l} + W_{S l}) \rho_t + W_{\downarrow l} \rho_S - W_{\downarrow l} \rho_S + \Delta_x \cos(\omega t) |\langle \rho_{\downarrow S}|t\rangle|,$$

$$\dot{\rho}_S = W_{S\uparrow} \rho_t + W_{S l} \rho_l - (W_{\uparrow l} + W_{\downarrow l}) \rho_S,$$

$$\dot{\rho}_{\uparrow\downarrow} = -i\Delta_z \rho_{\uparrow\downarrow} - i\Delta_x \cos(\omega t)(\rho_{\downarrow\downarrow} - \rho_{\downarrow\downarrow})/2 - V_{\uparrow\downarrow} \rho_{\uparrow\downarrow},$$

$$\dot{\rho}_{\downarrow\downarrow} = -i\Delta_z \rho_{\downarrow\downarrow} - V_{\downarrow\downarrow} \rho_{\downarrow\downarrow},$$

where $\rho_n = \langle n |\rho_D| n\rangle$ and $\rho_{nm} = \langle n |\rho_D| m\rangle$. We now specify the transition rates $W_{nm}$ for the diagonal and the spin decoherence rates $V_{nm} = V_{mn}$ for the off-diagonal elements of $\rho_D$. For the sequential tunneling rates [14], $W_{nm} = \sum_{l=1,2} W_{nl}^l$, we find, $W_{S\uparrow} = \gamma_l^f(\Delta_{S\uparrow})$ and $W_{\downarrow l} = \gamma_l^f[1 - f(\Delta_{S\downarrow})]$, with the Fermi function $f(\Delta_{S\downarrow}) = [1 + e^{(\Delta_{S\downarrow} - \mu)/kT}]^{-1}$. The transition rates are $\gamma_l^f = 2\pi\nu_l \sum_p t_{lp}^2 |\langle |d_p|S\rangle|^2$ with (possibly spin-dependent, see below) density of states $\nu_l$ at the Fermi energy. The rates $W_{S\uparrow}$ and $W_{\downarrow l}$ are defined analogously. Further, we allow for additional coupling of the electron spin to the environment (e.g. hyperfine or spin-phonon coupling). First, the spin relaxation rates $W_{\uparrow l}$ and $W_{\downarrow l}$ were inserted in Eqs. (1) and (2), corresponding to the phenomenological rate $1/T_1 = W_{\uparrow l} + W_{\downarrow l}$ [14]. We assume $W_{\uparrow l} \gg W_{\downarrow l}$ for $\Delta_l > kT$ (consistent with detailed balance, $W_{\uparrow l}/W_{\downarrow l} = e^{\Delta_l/kT}$). Second, the rate $1/T_2$ describes the intrinsic decoherence of the spin on the dot (which persists even if the tunnel coupling is switched off), contributing to $V_{\uparrow l}$. The contribution of $H_T$ to $V_{\uparrow l}$ is calculated as $(W_{S\uparrow} + W_{S l})/2$, i.e. electrons tunneling onto the dot destroy spin coherence on the dot. The total spin decoherence rate is $V_{\uparrow l} = (W_{S\uparrow} + W_{S l})/2 + 1/T_2$.

We calculate the stationary solution of Eqs. (1) and (2) in the rotating wave approximation [14], where only the leading frequency contributions of $H_{ESR}$ are retained. We obtain effective spin-flip rates $W_{\uparrow l}$ and $W_{\uparrow S}$, where the Rabi-flips produced by the ESR field are described by the rate

$$\dot{W}_{\uparrow l} = \frac{\Delta_x^2}{8} \left( \frac{V_{\uparrow l}}{\Delta_x} \right)^2 + V_{\uparrow l}^2,$$

which is a Lorentzian in $\omega$ with maximum $W_{\uparrow l}^\text{max} = \Delta_x^2/8V_{\uparrow l}^2$ at ESR resonance $\omega = \Delta_l$. With the stationary solution, we can now calculate the current $I$ (too lengthy to be shown here [15]), which we shall discuss next in different regimes.

![Figure 2](image_url)

**FIG. 2.** The stationary current $I(\omega)$ [Eq. (8)] for $B_z = 0.5$ T, $B_{\rho}^0 = 0.45$ G, $T_i = 1$ ps, $T_2 = 100$ ns, $\gamma_1 = 5 \times 10^6$ s$^{-1}$, and $\gamma_2 = 5 \times 10^9$, i.e. $W_{\uparrow l}^\text{max} < \gamma_1 < 1/T_2$. Here, the linewidth gives a lower bound for the intrinsic spin decoherence time $T_2$, while it becomes equal to $2/T_2$ for $B_{\rho}^0 = 0.08$ G and $\gamma_1 = 5 \times 10^6$ s$^{-1}$, where $I(\omega = \Delta_l) \approx 1.5$ fA.

**Zeeman blockade.** We consider a quantum dot as shown in Fig. [1] with $\Delta_z > \Delta_l$, $kT$, and $\Delta_{S\uparrow} > \mu_1 > \Delta_{S\downarrow} > \mu_2$, and $f_l(\Delta_{S\uparrow}) = 0$, with $\Delta_{S\uparrow} = \mu_1 - \mu_2$ being the applied bias. Thus, $W_{S\uparrow} = 0$ and $W_{S\downarrow} = \gamma_l^f$. Without ESR field, the dot relaxes into its ground state $|\downarrow\rangle$ (since $W_{\uparrow\downarrow} \ll W_{\uparrow l}$), and the sequential tunneling current through the dot is blocked by energy conservation. However, if an ESR field is present, producing Rabi spin-flips (on the dot only), the current flows through the dot involving state $|\uparrow\rangle$. First we consider $\gamma_l^f = \gamma_\uparrow = \gamma_l$. For $kT > \Delta_{S\uparrow}$ and $W_{\uparrow l}^\text{max} < \max\{W_{\uparrow l}, \gamma_l\}$ we obtain for the stationary current

$$I(\omega) = \frac{2e\gamma_\uparrow\gamma \tilde{W}_{\uparrow l}}{W_{\uparrow l} + 2},$$

while, for $W_{\uparrow l}^\text{max} < \max\{W_{\uparrow l}, \gamma_l f_l(\Delta_{S\uparrow} + \Delta_{S\downarrow}/2)\}$,

$$I(\omega) = \frac{e\gamma_\uparrow\gamma \tilde{W}_{\uparrow l}}{\gamma_\downarrow + \gamma_2} \frac{\Delta_{S\downarrow} - \mu}{2kT h(T) \cosh\left(\frac{\Delta_{S\downarrow} - \mu}{2kT}\right)},$$

for $kT > \Delta_{S\uparrow}$ and with $\mu = (\mu_1 + \mu_2)/2$. The standard sequential-tunneling peak-shape in Eq. (9) is modified by $h(T) = 2W_{\uparrow l} + (\gamma_\uparrow + \gamma_\downarrow - \gamma_l) f_l(\Delta_{S\downarrow} + \Delta_{S\uparrow}/2)$, which can affect position and width of the peak. Most important, the current $I(\omega)$ [Eqs. (8) and (9)] is proportional to the resonant rate $W_{\uparrow l}$. Thus, the current $I(\omega)$ as a function of the ESR frequency $\omega$ (or equivalently of $B_z$) has a resonant peak at $\omega = \Delta_{S\downarrow}$ width $2V_{\uparrow l}$. Since $V_{\uparrow l} \geq 1/T_2$, this width provides a lower bound on the intrinsic spin decoherence time $T_2$ of the single dot spin-[17]. For weak tunneling, $\gamma_l < 2/T_2$, this bound saturates, i.e. the width $2V_{\uparrow l}$ becomes $2/T_2$. Further, we note that the $g$ factor of a single dot can be measured via the position of the peak in $I(\omega)$ [or in $I(B_z)$], which could provide a useful technique to study $g$ factor modulated materials [18].
Pumping. Next we consider the case of zero bias, $\Delta \mu = 0$, and $f_1 = f_2$, but with $\gamma_{\uparrow} \neq \gamma_{\downarrow}$. Then, there is a finite current due to "pumping" \[14\] by the ESR source,

$$I(\omega) = e\tilde{W}_1(\gamma_{\uparrow}^1\gamma_{\uparrow}^{-1} - \gamma_{\uparrow}^{-1}\gamma_{\downarrow}^1)f_1(\Delta S_1)\left/\{(\gamma_{\uparrow}^{-1} + \gamma_{\downarrow}^{-1} - \tilde{W}_1)\right\} \times (\gamma_{\uparrow}^1 + \gamma_{\downarrow}^1)f_1(\Delta S_2) + (\tilde{W}_1 + \tilde{W}_2)(\gamma_{\uparrow}^1 + \gamma_{\downarrow}^1 + \gamma_{\uparrow}^2 + \gamma_{\downarrow}^2),$$

where $\gamma_{\uparrow}^1, \gamma_{\downarrow}^1, \gamma_{\uparrow}^2, \gamma_{\downarrow}^2$ determines the direction of the current. As in Eqs. (9) and (10), Eq. (11) has resonant behavior and $T_2$ can also be measured.

In addition to setups using spin-polarized leads \[22\] or spin-dependent tunneling, we now propose an alternative for producing $\gamma_{\uparrow}^2 \neq \gamma_{\downarrow}^2$. A second dot \("dot 2"\), see Fig. 1(b), acting as a spin filter, is coupled to the previous dot \("dot 1"\) with tunneling amplitude $t_{DD}$. The coupling of dot 2 to the lead shall be strong, leading to resonant tunneling with resonance width $\Gamma_2 = 2\pi\nu_{DD}^2$. We require $\Gamma_2 < \Delta S_2^2 - \mu_2$ to neglect electron-hole excitations in lead 2 \[23\]. We calculate the rates $\dot{\gamma}_{\uparrow}^2$ and $\dot{\gamma}_{\downarrow}^2$, for tunneling from dot 1 via dot 2 into lead 2 in a $T$-matrix approach, with tunnel Hamiltonian $H_T = H_{DD} + H_{DLP}$. We evaluate the transition rates $W_{f_{\uparrow}} = 2\pi\int[H_T \sum_{\alpha=\uparrow,\downarrow} (\varepsilon_i + i\eta - H_0 - \mu_2)^2 \rho_{\uparrow}\rho_{\downarrow}\varepsilon \rangle \langle \varepsilon_f \varepsilon_i \rangle d\epsilon_f$ by summing up contributions from all orders in $H_{DLP}$ and taking $\eta \rightarrow +0$. The Zeeman splitting $\Delta S_2^2$ in dot 2 shall be such that $\Delta S_2^2 \neq \Delta S_2^2$, $\Delta S_2^2 \neq \Delta S_2^2 - \mu_2$. This ensures by energy conservation that dot 2 is always in state $\{|\uparrow\rangle\}$ after an electron has passed. We now integrate over the final states in lead 2 and obtain the Breit-Wigner transition rate of an electron with spin down to go from dot 1 to lead 2 via the resonant level $E_{S}^2$ of dot 2,

$$\dot{\gamma}_{\uparrow}^2 = \Gamma_2|t_{DD}|^2/[(\Delta S_2^2 - \Delta S_2^2)^2 + (\Gamma_2/2)^2].$$

Since dot 2 is always in state $|\uparrow\rangle$, tunneling of a spin $\uparrow$ would involve the triplet level $E_{S_1_D}$ on dot 2, and thus $\dot{\gamma}_{\uparrow}^2$ is suppressed to zero (up to cotunneling contributions \[10\]). The proposed setup is thus again described by Eqs. (9) (10) with the tunneling rates $W_{S_4\downarrow}^2 = W_{S_2\downarrow}^2 = W_{S_0\downarrow}^2 = 0$ and $W_{S_2\downarrow}^2 = \tilde{\gamma}_{\uparrow}^2$, and we can use all previous results (for one dot), but with $\gamma_{\uparrow}^2 \rightarrow \tilde{\gamma}_{\uparrow}^2, \gamma_{\downarrow}^2 \rightarrow 0$, and $f_2(\Delta S_1) = 0$. In particular, we see from Eq. (10) that for zero bias $\Delta \mu = 0$ a current flows from lead 1 via the dots 1 and 2 to lead 2. We emphasize that this setup \[see Fig. 1(b)\] acts as a "spin inverter" with spin up electrons as input and spin down electrons as output (and no transmission of spin down electrons).

Spin read-out. We analyze now the time-dynamics of the read-out of a dot-spin via spin-polarized currents. For this, we consider a dot coupled to fully spin polarized leads, such that $\lambda_{\rm ads}^2 > \varepsilon_F > \Delta_{\uparrow}$. Since no electron with spin down can be provided or taken by the leads (since $\nu_{\downarrow} = 0$), the rates $W_{S_2\uparrow} = W_{S_0\uparrow}$ vanish (in contrast to the energy blocking $W_{S_2\uparrow}$ described above). Thus, a current can only flow if initially the state on the dot is $|\downarrow\rangle$, which allows to detect the initial spin state of the dot (strong measurement). The goal is now to characterize a measurement time $t_{\rm meas}$ for the spin read-out. For this, we need to keep track of the number of electrons $q$ which have accumulated in lead 2 since $t = 0$ \[23\] (above we have only studied averaged currents), i.e. we now consider the states $|n\rangle \rightarrow |n, q\rangle$. The time evolution of $P_\mu(q, t)$ (now charge-dependent) is described by Eqs. (6) (8), but with replacements $W_{S\uparrow}^2 \rho_{\mu}(q) \rightarrow W_{S\uparrow}^2 \rho_{\mu}(q-1)$ in Eq. (6) and $W_{S\downarrow}^2 \rho_{\mu}(q) \rightarrow W_{S\downarrow}^2 \rho_{\mu}(q+1)$ in Eq. (8). Next, we consider the distribution function $P_s(q, t) = \sum_n \rho_n(q, t)$ that $q$ charges have accumulated in lead 2 after time $t$ when the dot was in state $|i\rangle$ at $t = 0$. For a meaningful measurement of the dot-spin, the spin flip times $W_{S_0\downarrow}^{-1}, W_{S\uparrow}^{-1}$, and $1/\Delta_s$ must be smaller than $t_{\rm meas}$. Eqs. (8)–(11) then decouple except Eqs. (6) and (8), which we solve for $P_{s0}$ and $P_{s1}$ at $t = 0$ (the general solution follows by superposition). First, if we start in state $|\uparrow\rangle$, no charges tunnel through the dot and thus $P_s(q, t) = \delta_{q,0}$. Second, for the initial state $|\downarrow\rangle$, we consider $kT < \Delta \mu$ and equal rates $W_{S}\downarrow = W_{S}\uparrow = W$. We relabel the density matrix $\rho_{\mu}(q) \rightarrow \rho_{m=2q}$ and $\rho_{\mu}(q) \rightarrow \rho_{m=q+1}$, and Eqs. (6) and (8) become $\dot{\rho}_{m} = W(\rho_{m-1} - \rho_{m})$, with solution $\rho_{m}(t) = (Wt)^m e^{-Wt}/m!$ (Poissonian distribution). $P_s(q, t)$ then becomes

$$P_s(q, t) = \frac{(Wt)^q e^{-Wt}}{(2q)!} \left(1 + \frac{Wt}{2q + 1}\right).$$

Experimentally, $P_s(q, t)$ can be determined by time series measurements. The (inverse) signal-to-noise ratio is defined as the Fano factor \[22\],\[23\], which we calculate as $F_s(t) = \langle q(t)^2 \rangle / \langle q(t) \rangle = 1/2 + (3 - 2e^{2Wt}(4Wt + 1) - e^{-2Wt}) / 4(2Wt + 1 + e^{-2Wt})$, with $F_s$ decreasing monotonically from $F_s(0) = 1$ to $F_s(t \rightarrow \infty) = 1/2$.\[22\]. We can now quantify the measurement efficiency. If, after time $t_{\rm meas}$, some charges $q > 0$ have tunneled through the dot, the initial state of the dot was $|\downarrow\rangle$ with probability 1 (assuming that single charges can be detected via an SET \[23\]). However, if no charges were detected ($q = 0$), the initial state of the spin memory was $|\uparrow\rangle$ with probability 1 $P_s(0, t) = 1 - (W_s^1 e^{-W_s^1 t} - W_s^2 e^{-W_s^2 t})/(W_s^1 - W_s^2)$, which reduces to $1 - e^{-W t}(1 + Wt)$, for equal rates. Thus, roughly speaking, we find that $t_{\rm meas} \gtrsim 2W^{-1}$, as expected.

Rabi oscillations and Zeno effect. We show that coherent oscillations of the dot-spin induced by ESR lead to coherent oscillation in the current, again for spin polarized leads. For $\mu_1 > \Delta_{S_1} > \mu_2$ and $kT < \Delta_{\mu}$, the current in lead 1 is $I_1(t) = -eW_{1\downarrow}^2 \rho_{\mu}(t)$ and $I_2(t) = eW_{1\uparrow}^2 \rho_{\mu}(t)$ in lead 2. (Note that in general $I_1(t) + I_2(t) \neq 0$, since charge can accumulate on the dot.) Thus, the time-dependence of $\rho_{\mu}$ and $\rho_{\mu}$ can be measured via the currents $I_{1,2}$. Note that the spin-polarized electrons from lead 1 perform a projective measurement, leaving the dot-spin in either up or down state. Thus, to obtain $I_{1,2}$
experimentally, an ensemble average is required, e.g.
by using an array of (independent) dots arranged in parallel
or by time-series measurement over a single dot. In Fig.
3 we plot the numerical solution of Eqs. (4)–(8), showing
coherent Rabi oscillations of $\rho_{\uparrow\downarrow}$, $\rho_{\downarrow\uparrow}$ and their decay to the
stationary solution, dominated by the spin decoherence
$V_{\uparrow\uparrow}$. Thus, $V_{\uparrow\uparrow}$ (and $1/T_2$) can be accessed here directly
in the time domain [27].

**FIG. 3.** Rabi oscillations visible in the time evolution of
the density matrix $\rho_{\uparrow\downarrow}$ (dotted), $\rho_{\downarrow\uparrow}$ (dashed) and $\rho_S$ (full line)
for $W_{S\downarrow} = W_{\downarrow S} = 4 \times 10^7$ s$^{-1}$, $T_1 = 1$ ms, $T_2 = 300$ ns,
and $\Delta_s = 5W_{S\downarrow}$ (corresponding to $B_s = 10$ G for $g = 2$).
During the time span shown here, on average 3 electrons
have tunneled through the dot. Here the spin decoherence
is dominated by the measurement process, $W_{S\downarrow} \gg 1/T_2$,
however, for weaker measurement, it will be determined by
$T_2$. In the inset we show the case of a strong measurement,
$W_{S\downarrow} = 10^9$ s$^{-1}$. As a consequence of the Zeno effect
(see text), the Rabi oscillations are suppressed. Further,$\rho_{\downarrow\uparrow}$ and $\rho_S$ are indistinguishable since $|\downarrow\rangle$ and $|S\rangle$ equilibrate
rapidly due to the increased tunneling.

Finally, increasing $W_{S\downarrow}$, the coherent oscillations of $\rho_{\uparrow\downarrow}$,$\rho_{\downarrow\uparrow}$ become suppressed (see Fig. 3) due to increased transfer
of charges which perform a continuous strong measure-
ment on the dot-spin. This suppression, known as
Zeno effect [28], occurs in $\rho_{\uparrow\downarrow}$, $\rho_{\downarrow\uparrow}$ and thus is observable
in the input current $I_1(t)$.

**Conclusions.** We have proposed a setup to measure
the single spin decoherence time $T_2$ of a dot in Coulomb blockage regime, coupled to leads, via the stationary and
time-dependent current by using ESR techniques. We
have discussed pumping and read-out processes.

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field, achieved e.g. by applying $B_s$ and/or $B_z$ locally or
with different $g$ factors; to account for this, we simply
assume $\Delta_1 \neq \Delta_2$.
[12] We use a time-series measurement over a single dot. In Fig.
3, we plot the numerical solution of Eqs. (4)–(8), showing
coherent Rabi oscillations of $\rho_{\uparrow\downarrow}$, $\rho_{\downarrow\uparrow}$ and their decay to the
stationary solution, dominated by the spin decoherence
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