

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 [1–6] 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 [7,8], 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 [1]. 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 μs [6], consistent with calculations [9]. 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) [10], 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 [11], coupled to two Fermi-liquid leads $l = 1, 2$ at chemical potentials μ_l . We consider the Hamiltonian $H = H_0 + H_T = H_{\text{lead}} + \tilde{H}_{\text{dot}} + H_T$, which describes leads, dot and the tunnel coupling between leads and dot, *resp.* Here, $\tilde{H}_{\text{dot}} = H_{\text{dot}} + H_{\text{ESR}}$, where H_{dot} includes charging and interaction energies of the electrons on the dot. H_{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 μ_B , and Pauli matrix σ_z . Coupling to

an oscillating magnetic ESR field in x -direction of frequency ω is included in $H_{\text{ESR}} = -\frac{1}{2}\Delta_x\cos(\omega t)\sigma_x$, with $\Delta_x = g\mu_B B_x^0$ and Pauli matrix σ_x . Such an oscillating field produces Rabi spin-flips at $\omega = \Delta_z$, as used in ESR. We assume Zeeman splitting of the leads $\Delta_z^{\text{leads}} \neq \Delta_z$ and $\Delta_z^{\text{leads}} \ll \varepsilon_F$, where ε_F is the Fermi energy, such that the field effects of B_z and $B_x(t)$, *resp.*, on the leads are negligible. Such a situation can be achieved by using materials of different g factors [4] and/or with local magnetic fields.

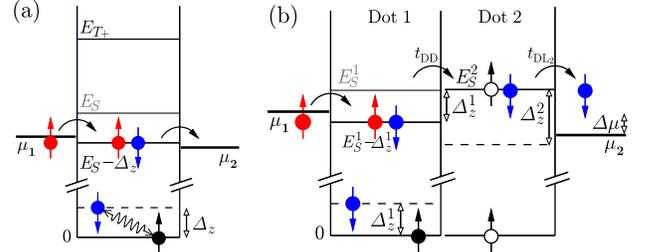


FIG. 1. (a) Dot coupled to unpolarized leads with chemical potentials $\mu_{1,2}$, in the sequential tunneling regime defined by $E_S > \mu_1 > E_S - \Delta_z > \mu_2$, with the singlet/triplet levels E_S/E_{T+} and Zeeman splitting $\Delta_z = g\mu_B B_z$. If initially the spin-state on the dot is $|\uparrow\rangle$, sequential tunneling is blocked by energy conservation. Exciting the dot-spin via ESR (Rabi flip) the dot becomes unblocked but only for spin up electrons from lead 1. Finally, from the singlet, spin up or down can tunnel into lead 2. (b) Extended setup where the additional dot 2 (with $|t_{\text{DD}}| < |t_{\text{DL}_2}|$ and tuned to resonance) acts as a spin filter in the regime $E_S^1 \approx E_S^2$; $E_S^1 > \mu_1 > E_S^1 - \Delta_z^1$; $\mu_2 > E_S^1 - \Delta_z^2$; $E_S^2 > \mu_2$, and $\Delta_z^1 \neq \Delta_z^2$ [12]. The allowed transition sequence is schematically given by $\uparrow(\uparrow)_1(\uparrow)_2 \xrightarrow{\text{ESR}} \uparrow(\downarrow)_1(\uparrow)_2 \rightarrow \uparrow(\downarrow)_1(\uparrow)_2 \rightarrow \uparrow(\uparrow)_1(\downarrow)_2 \leftrightarrow \uparrow(\uparrow)_1(\uparrow)_2\downarrow$ (see text).

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 $|S\rangle = (|\uparrow\downarrow\rangle - |\downarrow\uparrow\rangle)/\sqrt{2}$ (which can be achieved by tuning B_z [13]). 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 $\Delta_{S\uparrow(\downarrow)} = E_S - E_{S\uparrow(\downarrow)}$ (con-

taining charging energy U), which can then be compared with μ_l .

Master equation. We derive the master equation for the reduced density matrix of the dot, $\rho_D = \text{Tr}_L \rho$. Here, Tr_L is the trace taken over the leads and ρ 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 [14,15]. We obtain the following master equation

$$\begin{aligned} \dot{\rho}_\uparrow = & -(W_{\downarrow\uparrow} + W_{S\uparrow})\rho_\uparrow + W_{\uparrow\downarrow}\rho_\downarrow + W_{\uparrow S}\rho_S \\ & + \Delta_x \cos(\omega t) \text{Im}[\rho_{\downarrow\uparrow}], \end{aligned} \quad (1)$$

$$\begin{aligned} \dot{\rho}_\downarrow = & W_{\downarrow\uparrow}\rho_\uparrow - (W_{\uparrow\downarrow} + W_{S\downarrow})\rho_\downarrow + W_{\downarrow S}\rho_S \\ & - \Delta_x \cos(\omega t) \text{Im}[\rho_{\downarrow\uparrow}], \end{aligned} \quad (2)$$

$$\dot{\rho}_S = W_{S\uparrow}\rho_\uparrow + W_{S\downarrow}\rho_\downarrow - (W_{\uparrow S} + W_{\downarrow S})\rho_S, \quad (3)$$

$$\dot{\rho}_{\downarrow\uparrow} = -i\Delta_z\rho_{\downarrow\uparrow} - i\Delta_x\cos(\omega t)(\rho_\uparrow - \rho_\downarrow)/2 - V_{\downarrow\uparrow}\rho_{\downarrow\uparrow}, \quad (4)$$

$$\dot{\rho}_{S\uparrow} = -i\Delta_{S\uparrow}\rho_{S\uparrow} - V_{S\uparrow}\rho_{S\uparrow}, \quad (5)$$

$$\dot{\rho}_{S\downarrow} = -i\Delta_{S\downarrow}\rho_{S\downarrow} - V_{S\downarrow}\rho_{S\downarrow}, \quad (6)$$

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 ρ_D . For the sequential tunneling rates [16], $W_{nm} = \sum_{l=1,2} W_{nm}^l$, we find, $W_{S\downarrow}^l = \gamma_l^\uparrow f_l(\Delta_{S\downarrow})$ and $W_{\downarrow S}^l = \gamma_l^\uparrow [1 - f_l(\Delta_{S\downarrow})]$, with the Fermi function $f_l(\Delta_{S\downarrow}) = [1 + e^{(\Delta_{S\downarrow} - \mu_l)/kT}]^{-1}$. The transition rates are $\gamma_l^\uparrow = 2\pi\nu_\uparrow |\sum_p t_{lp} \langle \downarrow | d_{p\uparrow} | S \rangle|^2$ with (possibly spin-dependent, see below) density of states ν_\uparrow at the Fermi energy. The rates $W_{S\uparrow}^l$ and $W_{\uparrow S}^l$ 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 $W_{\uparrow\downarrow}$ and $W_{\downarrow\uparrow}$ were inserted in Eqs. (1) and (2), corresponding to the phenomenological rate $1/T_1 = W_{\uparrow\downarrow} + W_{\downarrow\uparrow}$ [14]. We assume $W_{\uparrow\downarrow} \gg W_{\downarrow\uparrow}$ for $\Delta_z > kT$ (consistent with detailed balance, $W_{\uparrow\downarrow}/W_{\downarrow\uparrow} = e^{\Delta_z/kT}$). 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 $V_{\downarrow\uparrow}$. The contribution of H_T to $V_{\downarrow\uparrow}$ is calculated as $(W_{S\uparrow} + W_{S\downarrow})/2$, i.e. electrons tunneling onto the dot destroy spin coherence on the dot. The total spin decoherence rate is $V_{\downarrow\uparrow} = (W_{S\uparrow} + W_{S\downarrow})/2 + 1/T_2$.

We calculate the stationary solution of Eqs. (1)–(6) in the rotating wave approximation [14], where only the leading frequency contributions of H_{ESR} are retained. We obtain effective spin-flip rates $\widetilde{W}_{\uparrow\downarrow} = W_{\uparrow\downarrow} + W_\omega$ and $\widetilde{W}_{\downarrow\uparrow} = W_{\downarrow\uparrow} + W_\omega$, where the Rabi-flips produced by the ESR field are described by the rate

$$W_\omega = \frac{\Delta_x^2}{8} \frac{V_{\downarrow\uparrow}}{(\omega - \Delta_z)^2 + V_{\downarrow\uparrow}^2}, \quad (7)$$

which is a Lorentzian in ω with maximum $W_\omega^{\text{max}} = \Delta_x^2/8V_{\downarrow\uparrow}$ at ESR resonance $\omega = \Delta_z$. 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.

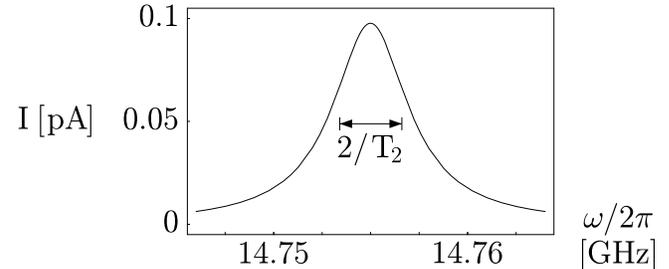


FIG. 2. The stationary current $I(\omega)$ [Eq. (8)] for $B_z = 0.5$ T, $B_x^0 = 0.45$ G, $T_1 = 1$ μ s, $T_2 = 100$ ns, $\gamma_1 = 5 \times 10^6$ s $^{-1}$, and $\gamma_2 = 5\gamma_1$, i.e. $W_\omega^{\text{max}} < \gamma_1 < 1/T_2$. Here, the linewidth gives a lower bound for the intrinsic spin decoherence time T_2 , while it becomes equal to $2/T_2$ for $B_x^0 = 0.08$ G and $\gamma_1 = 5 \times 10^5$ s $^{-1}$, where $I(\omega = \Delta_z) \approx 1.5$ fA.

Zeeman blockade. We consider a quantum dot as shown in Fig. 1, with $\Delta_z > \Delta\mu$, kT , and $\Delta_{S\uparrow} > \mu_1 > \Delta_{S\downarrow} > \mu_2$, and $f_l(\Delta_{S\uparrow}) = 0$, with $\Delta\mu = \mu_1 - \mu_2$ being the applied bias. Thus, $W_{S\uparrow} = 0$ and $W_{\uparrow S}^l = \gamma_l^\downarrow$. Without ESR field, the dot relaxes into its ground state $|\uparrow\rangle$ (since $W_{\downarrow\uparrow} \ll W_{\uparrow\downarrow}$), 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 $|\downarrow\rangle$. First we consider $\gamma_l^\downarrow = \gamma_l^\uparrow = \gamma_l$. For $kT > \Delta\mu$ and $W_\omega^{\text{max}} < \max\{W_{\uparrow\downarrow}, \gamma_1\}$ we obtain for the stationary current

$$I(\omega) = \frac{2e\gamma_1\gamma_2\widetilde{W}_{\downarrow\uparrow}}{W_{\uparrow\downarrow}(\gamma_1 + 2\gamma_2) + \gamma_1(\gamma_1 + \gamma_2)}, \quad (8)$$

while, for $W_\omega^{\text{max}} < \max\{W_{\uparrow\downarrow}, \gamma_1 f_l(\Delta_{S\downarrow} + \Delta\mu/2)\}$,

$$I(\omega) = \frac{e\gamma_1\gamma_2\widetilde{W}_{\downarrow\uparrow}}{\gamma_1 + \gamma_2} \frac{\Delta\mu}{2kTh(T)} \cosh^{-2}\left(\frac{\Delta_{S\downarrow} - \mu}{2kT}\right), \quad (9)$$

for $kT > \Delta\mu$ and with $\mu = (\mu_1 + \mu_2)/2$. The standard sequential-tunneling peak-shape in Eq. (9) is modified by $h(T) = 2W_{\uparrow\downarrow} + (\gamma_1 + \gamma_2 - W_{\uparrow\downarrow}) f_l(\Delta_{S\downarrow} + \Delta\mu/2)$, 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_ω . Thus, the current $I(\omega)$ as a function of the ESR frequency ω (or equivalently of B_z) has a resonant peak at $\omega = \Delta_z$ of width $2V_{\downarrow\uparrow}$. Since $V_{\downarrow\uparrow} \geq 1/T_2$, this width provides a lower bound on the intrinsic spin decoherence time T_2 of a single dot-spin [17]. For weak tunneling, $\gamma_1 < 2/T_2$, this bound saturates, i.e. the width $2V_{\downarrow\uparrow}$ 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 [8,18].

Pumping. Next we consider the case of zero bias, $\Delta\mu = 0$, and $f_1 = f_2$, but with $\gamma_i^\downarrow \neq \gamma_i^\uparrow$. Then, there is a finite current due to “pumping” [19] by the ESR source,

$$I(\omega) = e\widetilde{W}_{\downarrow\uparrow}(\gamma_1^\uparrow\gamma_2^\downarrow - \gamma_1^\downarrow\gamma_2^\uparrow)f_1(\Delta_{S\downarrow}) / \left\{ (\gamma_1^\downarrow + \gamma_2^\downarrow - \widetilde{W}_{\uparrow\downarrow}) + (\gamma_1^\uparrow + \gamma_2^\uparrow + \gamma_1^\downarrow + \gamma_2^\downarrow) \right\}, \quad (10)$$

where $\gamma_1^\uparrow\gamma_2^\downarrow - \gamma_1^\downarrow\gamma_2^\uparrow$ determines the direction of the current. As in Eqs. (8) and (9), Eq. (10) has resonant behavior and T_2 can also be measured.

In addition to setups using spin-polarized leads [20] or spin-dependent tunneling, we now propose an alternative for producing $\gamma_i^\downarrow \neq \gamma_i^\uparrow$. 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 $\Gamma_2 = 2\pi\nu|t_{\text{DL}_2}|^2$. We require $\Gamma_2 < \Delta_{S\uparrow}^2 - \mu_2$ to neglect electron-hole excitations in lead 2 [21]. We calculate the rates $\hat{\gamma}^\uparrow$ and $\hat{\gamma}^\downarrow$, for tunneling from dot 1 via dot 2 into lead 2 in a T -matrix approach, with tunnel Hamiltonian $H_T = H_{\text{DD}} + H_{\text{DL}_2}$. We evaluate the transition rates $W_{fi} = 2\pi|\langle f|H_T \sum_{n=0}^{\infty} [(\varepsilon_i + i\eta - H_0)^{-1}H_T]^n|i\rangle|^2 \delta(\varepsilon_f - \varepsilon_i)$ by summing up contributions from all orders in H_{DL_2} and taking $\eta \rightarrow +0$. The Zeeman splitting Δ_z^2 in dot 2 shall be such that $\Delta_z^1 \not\approx \Delta_z^2$ [12] and $\Delta_z^2 > \Delta_{S\downarrow}^1 - \mu_2$. 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,

$$\hat{\gamma}^\downarrow = \Gamma_2|t_{\text{DD}}|^2 / [(\Delta_{S\uparrow}^1 - \Delta_{S\uparrow}^2)^2 + (\Gamma_2/2)^2]. \quad (11)$$

Since dot 2 is always in state $|\uparrow\rangle$, tunneling of a spin \uparrow would involve the triplet level E_{T_+} on dot 2, and thus $\hat{\gamma}^\uparrow$ is suppressed to zero (up to cotunneling contributions [16]). The proposed setup is thus again described by Eqs. (1)–(6) with the tunneling rates $W_{S\downarrow}^2 = W_{\downarrow S}^2 = W_{S\uparrow}^2 = 0$ and $W_{\uparrow S}^2 = \hat{\gamma}^\downarrow$, and we can use all previous results (for one dot), but with $\gamma_2^\downarrow \rightarrow \hat{\gamma}^\downarrow$, $\gamma_2^\uparrow \rightarrow 0$, and $f_2(\Delta_{S\uparrow}) = 0$. 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 $\Delta_z^{\text{leads}} > \varepsilon_F > \Delta_z$ [20]. Since no electron with spin down can be provided or taken by the leads (since $\nu_\downarrow = 0$), the rates $W_{S\uparrow} = W_{\uparrow S}$ vanish (in contrast to the energy blocking $W_{S\uparrow} = 0$ described above). Thus, a current can only flow if initially the

state on the dot is $|\downarrow\rangle$ [22], which allows to detect the initial spin state of the dot (strong measurement). The goal is now to characterize a measurement time t_{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 $\rho_D(q, t)$ (now charge-dependent) is described by Eqs. (1)–(6), but with replacements $W_{\downarrow S}^2 \rho_S(q) \rightarrow W_{\downarrow S}^2 \rho_S(q - 1)$ in Eq. (2) and $W_{S\downarrow}^2 \rho_\downarrow(q) \rightarrow W_{S\downarrow}^2 \rho_\downarrow(q + 1)$ in Eq. (3). Next, we consider the distribution function $P_i(q, t) = \sum_n \rho_n(q, t)$ 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 $W_{\uparrow\downarrow}^{-1}$, $W_{\downarrow\uparrow}^{-1}$, and $1/\Delta_x$ must be smaller than t_{meas} . Eqs. (1)–(6) then decouple except Eqs. (2) and (3), which we solve for $\rho_\uparrow = 1$ and $\rho_\downarrow = 1$ 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 $P_\uparrow(q, t) = \delta_{q0}$. Second, for the initial state $|\downarrow\rangle$, we consider $kT < \Delta\mu$ and equal rates $W_{S\downarrow}^1 = W_{\downarrow S}^2 = W$. We relabel the density matrix $\rho_\downarrow(q) \rightarrow \rho_{m=2q}$ and $\rho_S(q) \rightarrow \rho_{m=2q+1}$, and Eqs. (2) and (3) become $\dot{\rho}_m = W(\rho_{m-1} - \rho_m)$, with solution $\rho_m(t) = (Wt)^m e^{-Wt}/m!$ (Poissonian distribution). $P_i(q, t)$ then becomes

$$P_\downarrow(q, t) = \frac{(Wt)^{2q} e^{-Wt}}{(2q)!} \left(1 + \frac{Wt}{2q+1} \right). \quad (12)$$

Experimentally, $P_\downarrow(q, t)$ can be determined by time series measurements. The (inverse) signal-to-noise ratio is defined as the Fano factor [24,25], which we calculate as $F_\downarrow(t) = \langle \delta q(t)^2 \rangle / \langle q(t) \rangle = 1/2 + [3 - 2e^{-2Wt}(4Wt + 1) - e^{-4Wt}]/4(2Wt - 1 + e^{-2Wt})$, with F_\downarrow decreasing monotonically from $F_\downarrow(0) = 1$ to $F_\downarrow(t \rightarrow \infty) = 1/2$ [26]. We can now quantify the measurement efficiency. If, after time t_{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 [25]). However, if no charges were detected ($q = 0$), the initial state of the spin memory was $|\uparrow\rangle$ with probability $1 - P_\downarrow(0, t) = 1 - (W_{S\downarrow}^1 e^{-W_{\downarrow S}^2 t} - W_{\downarrow S}^2 e^{-W_{S\downarrow}^1 t}) / (W_{S\downarrow}^1 - W_{\downarrow S}^2)$, which reduces to $1 - e^{-Wt}(1 + Wt)$, for equal rates. Thus, roughly speaking, we find that $t_{\text{meas}} \gtrsim 2W^{-1}$, 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 $\mu_1 > \Delta_{S\downarrow} > \mu_2$ and $kT < \Delta\mu$, the current in lead 1 is $I_1(t) = -eW_{S\downarrow}^1 \rho_\downarrow(t)$ and $I_2(t) = eW_{\downarrow S}^2 \rho_S(t)$ 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 ρ_\downarrow and ρ_S 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. (1)–(6), showing coherent Rabi oscillations of ρ_{\uparrow} , ρ_{\downarrow} and their decay to the stationary solution, dominated by the spin decoherence $V_{\downarrow\uparrow}$. Thus, $V_{\downarrow\uparrow}$ (and $1/T_2$) can be accessed here directly in the time domain [27].

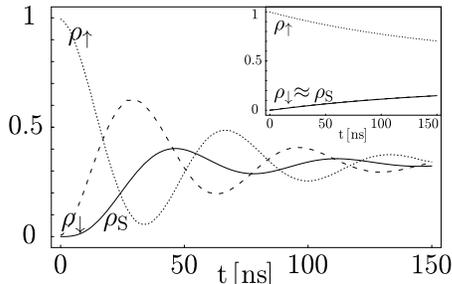


FIG. 3. Rabi oscillations visible in the time evolution of the density matrix ρ_{\uparrow} (dotted), ρ_{\downarrow} (dashed) and ρ_S (full line) for $W_{S\downarrow} = W_{\downarrow S} = 4 \times 10^7 \text{ s}^{-1}$, $T_1 = 1 \mu\text{s}$, $T_2 = 300 \text{ ns}$, and $\Delta_x = 5W_{S\downarrow}$ (corresponding to $B_x = 10 \text{ G}$ 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, $W_{S\downarrow} = W_{\downarrow S} = 10^9 \text{ s}^{-1}$. As a consequence of the Zeno effect (see text), the Rabi oscillations are suppressed. Further, ρ_{\downarrow} and ρ_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 ρ_{\uparrow} , ρ_{\downarrow} become suppressed (see Fig. 3) due to increased transfer of charges which perform a continuous strong measurement on the dot-spin. This suppression, known as Zeno effect [28], occurs in ρ_{\uparrow} , ρ_{\downarrow} 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 blockade 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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 [27] An alternative method for measuring T_2 is pulsed ESR. An initial spin state, say, $|\downarrow\rangle$, is rotated with a $\pi/2$ pulse into the xy -plane and starts to precess with frequency Δ_z . After a delay τ , the (remaining) coherent part of the spin is rotated by another $\pi/2$ pulse and finally ρ_{\downarrow} is measured via the current. Since the final spin direction, and thus ρ_{\downarrow} , depends on the precession angle, the measured current is an oscillating function of τ with frequency Δ_z , which decays on a timescale T_2 .
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